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Burgeoning the Higgs mass to 125 GeV through messenger-matter interactions in GMSB models

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Abstract: A 125 GeV Higgs renders the simpler GMSB models unnatural, essentially pushing the soft spectrum beyond the LHC reach. A direct coupling of the matter and messenger fields, that facilitates an enhanced mixing in the squark sector, is a way to ameliorate this deficiency. We construct all possible messenger-matter interaction terms considering the messenger multiplets in 1, 5 and 10 dimensional representations of the SU(5). A Froggatt-Nielsen like flavor framework connected with the origin of fermion mass hierarchy is utilized to control the interaction terms and suppress FCNC. We perform a detailed comparative study of the efficiency of such interaction terms to boost the Higgs mass keeping the soft spectrum light. We identify the more promising models and comment on their status in present and future collider studies.

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1 Introduction

Identification of the recently observed [1, 2] scalar field at the LHC with the Higgs would cast a long shadow on Gauge Mediated Supersymmetry Breaking (GMSB) models [3, 4]. This is a direct consequence of the well known phenomenon that in pure gauge mediation models, mixing in the scalar sector is minimum. This implies that in order to raise the Higgs mass to $\sim 125\,\text{GeV}$ one needs stop masses at several TeV, castigating these models to an *unnatural* existence with bleak chances to be probed at collider experiments like the LHC.

A resolution of this predicament is to consider direct coupling between the messenger and matter fields or models with gauge messengers [5]. These scenarios can in principle generate a sizable trilinear coupling that can boost the Higgs mass given by,

$$m_h^2 = M_Z^2 \cos^2 2\beta + \frac{3}{4\pi^2}\frac{m_t^4}{v^2}\left[\log\frac{M_S^2}{m_t^2} + \frac{X_t^2}{M_S^2}\left(1 - \frac{X_t^2}{12M_S^2}\right)\right], \tag{1.1}$$

where, $X_t = A_t - \mu \cot \beta$ and $M_s = \sqrt{m_{\tilde{t}_1} m_{\tilde{t}_2}}$, while keeping the scalar spectrum within the range of interest for collider physics. In case of gauge messengers the trilinear couplings are proportional to gauge charges. However realistic models recently studied in [6] predict a relatively heavy spectrum. In this paper we will consider models of direct messenger-matter interactions that can lead to relatively large trilinear coupling and a considerably light soft spectrum making them more pleasing from a fine-tuning point of view and more interesting phenomenologically. These interaction terms generally lead to new contributions to the scalar masses, thus producing a correlated perturbation of the pure gauge mediation soft spectrum. The strength of the interaction and hence the size of the trilinear coupling is principally constrained from the considerations to keep the scalar masses non-tachyonic and achieve radiative electroweak symmetry breaking (REWSB). Also a cause for concern are the complications in the flavor sector [7–9]. These tend to push the messenger scale upward into regions where the Gravitino mass goes beyond $\sim$ sub-KeV range putting unfavorable upper bounds on the reheating temperature inviting strong constraints from BBN [10]. The flavor constraints are severe for the first two generations of fermions and can be minimized by considering that the messengers preferentially couple to the third generation, which is relevant for enhancing the Higgs mass. This can be ensured by imposing judiciously chosen flavor symmetries.

Several models of GMSB augmented with messenger-matter interactions to alleviate the problem with the heavy Higgs have been suggested in the literature [11–17]. The basic idea being the coupling of the messengers to one of the Q, U, H_u MSSM multiplets, generating a sizable A_t at the messenger scale. Possible interaction terms in a given scenario get determined by the content and quantum numbers of the messenger sector. In this paper we make a systematic study of messenger-matter interactions, possible within a well defined framework. As an organizing strategy we consider the messengers as vector pairs embedded in one of the simpler representations of the GUT group SU(5). We will restrict ourselves to messengers in the 1, 5 and 10 dimensional representations of SU(5) or an admixture. This ensures that the perturbative unification[1] of gauge coupling is not ruined due to introduction of these additional fields [18]. For the first time we make a quantitative comparison of various models, including some that have been discussed in the literature and some entirely new scenarios, in term of their effectiveness to raise the Higgs mass without the usual pitfall of large scalar masses. In order to make the comparison, we numerically scan all parameters of every given model over a suitable range and project the allowed regions on a common parameter space. This allows us to make precise statements about scenarios preferred by the recent data on the *supposed* Higgs mass.

For simplicity, in our study we will not attempt to model the hidden sector and simply assume supersymmetry is broken by the vev of a spurion field which couples to the messengers. The crucial point would be that the messengers other than having usual gauge couplings to the MSSM sector, now also couple directly through the superpotential. In this paper, we enumerate possible messenger-matter interaction terms allowed by a given mes-

[1]Note that it is possible to ensure perturbative gauge coupling unification without considering complete representations of the GUT groups [20]. These *magic* combinations can potentially lead to interesting phenomenological scenarios [21], we will not discuss them in this paper.

senger sector. For each scenario we compute the contributions of these new terms to the scalar masses at one and two loop order. We find that the one loop contributions are always tachyonic but they are suppressed by $x^2 \equiv (F/M^2)^2$ where M and F are messenger scale and supersymmetry breaking scale respectively. Thus for a given $\Lambda = F/M$ or soft scalar mass, this contribution decouples as the scale of supersymmetry breaking is increased. The sign of the two loop contribution is model dependent, however there is no suppression from the supersymmetry breaking scale. In most regions of the parameter space this becomes the dominant contribution from the new terms. One loop renormalization group equations are used to run these soft parameters down from the scale of supersymmetry breaking to the weak scale and the sparticle spectrum is generated.

The rest of the paper is organized as follows. In section 2 we describe our organizing principle and enumerate the possible interaction terms. In section 3 we compute the soft masses and trilinear couplings for the different models. In section 4 we describe our numerical procedure and present our results comparing the models. Finally we conclude.

2 Messenger-matter interactions

In this section, we collect the possible interactions between the messenger sector and the MSSM fields. A useful way to organize these is to consider the fields embedded in a representation of some GUT group like SU(5). As usual, the MSSM fields can be embedded into representations of the SU(5) group as follows,

$$\begin{aligned}
\bar{5} &= \left(\bar{3}, 1, \frac{1}{3}\right) \oplus \left(1, 2, -\frac{1}{2}\right) = D^c \oplus L, \\
10 &= \left(3, 2, \frac{1}{6}\right) \oplus \left(\bar{3}, 1, -\frac{2}{3}\right) \oplus (1,1,1) = Q \oplus U^c \oplus E^c \\
\bar{5}_H &= (\text{integrated out field}) \oplus H_d, \\
5_H &= (\text{integrated out field}) \oplus H_u
\end{aligned}$$

The messenger fields are in a vector like pair embedded in 1, 5, 10 and their conjugate representations of SU(5). For the rest of this paper we use the following nomenclature for the messenger sector,

$$\begin{aligned}
1_m &= S_m, & 5_m &= \tilde{D}^c_m \oplus H^m_u, & \bar{5}_m &= D^c_m \oplus H^m_d \\
10_m &= Q_m \oplus U^c_m \oplus E^c_m, & \bar{10}_m &= \tilde{Q}_m \oplus \tilde{U}^c_m \oplus \tilde{E}^c_m,
\end{aligned}$$

where the subscript m identifies the messenger fields.

Technically the singlet is not a valid gauge messenger in the usual sense. However, it can couple directly to the visible sector through superpotential terms and thus will be considered in the following discussion. Within this framework for the messenger and the matter sectors, it is simple to write down all the possible SU(5) invariants that can be constructed:

Only Singlets (i) $5_H\bar{5}_H 1_m$ (ii) $5_H\bar{5}\ 1_m$

Only $\mathbf{5}\oplus\bar{\mathbf{5}}$ (i) $10\ \bar{5}_H\bar{5}_m$ (ii) $10\ \bar{5}\ \bar{5}_m$ (iii) $10\ \bar{5}_m\bar{5}_m$ (iv) $10\ \bar{5}_{mi}\bar{5}_{mj}$ (v) $10\,10\ 5_m$

Only $\mathbf{10}\oplus\overline{\mathbf{10}}$ (i) $10\,10_m 5_H$ (ii) $10_m\,10_m 5_H$ (iii) $\overline{10}_m\,\overline{10}_m\bar{5}_H$ (iv) $\overline{10}_m\,\overline{10}_m\bar{5}$.

Note that the invariants $\overline{10}_m\,\overline{10}_m\bar{5}_H$ and $\overline{10}_m\,\overline{10}_m\bar{5}$ will not give A_t hence will not be considered further. The possibility that the messenger sector is composed of fields in different representation can also exist. In this case we expect the usual invariants listed above should reappear. The new possible invariants are given below,

Singlet + $\mathbf{5}\oplus\bar{\mathbf{5}}$ (i) $5_H\bar{5}_m 1_m$ (ii) $5_m\bar{5}1_m$ (iii) $5_m\bar{5}_H 1_m$

Singlet + $\mathbf{10}\oplus\overline{\mathbf{10}}$ (i) $10\,\overline{10}1_m$

$\mathbf{5}\oplus\bar{\mathbf{5}}+\mathbf{10}\oplus\overline{\mathbf{10}}$ (i) $10\,10_m\,5_m$ (ii) $\overline{10}_m\,5_m 5_H$

Again we find that terms like $5_m\bar{5}S$ and $5_m\bar{5}_H S$ will not give A_t. To keep the discussion tractable we will utilize prudently chosen flavor symmetries to suppress interaction terms other than the ones listed above.

As emphasized earlier a host of issues with flavor including FCNC can be controlled if we consider scenarios where the messenger sector preferentially couples to the third generation multiplet. An economical proposition is to connect this to the flavor symmetry that is responsible for the hierarchy in measured mass of the Standard Model fermions [12]. For example one can consider the Froggatt-Nielsen [22] framework that necessitates an $U(1)_F$ flavor symmetry group, under which the MSSM chiral multiplets may be charged. This flavor symmetry is spontaneously broken at some high scale M_{string} by the vev of the flavon field (Z), with a conventionally assigned flavor charge -1. For this choice, operators with a negative $U(1)_F$ charge mismatch (ΔF) are considered absent in the effective theory. However, if $\Delta F > 0$ for any superpotential term then it is suppressed by the usual Froggatt-Nielsen factor $\epsilon^{\Delta F}$ where $\epsilon = \langle Z\rangle/M_{string}$. Flavor symmetries of this kind can successfully explain observed fermion mass matrices within present experimental uncertainties. Interestingly the anomalous nature of the symmetry can motivate compensatory exotic particles that can be probed at present and future collider experiments [23–25], thus providing an handle to probe the existence of these flavor structures. Together with R-parity assignments this will be enough to determine the interaction terms uniquely for most of the scenarios that will be explored in this paper. The usual charges for the MSSM and the flavon multiplets are given in table 1.

Another symmetry that is usually useful in safely segregating the messenger and visible sectors in usual GMSB models is the messenger parity [26] under which all the messenger multiplets are assumed odd while the MSSM multiplets are considered even. One can classify the interaction terms introduced above into two categories with different implications for possible messenger parities. In models where the interaction term involves two messenger multiplets and one MSSM multiplet one can impose messenger parity consistently.

Multiplets	10_1	10_2	10_3	$\bar{5}_1$	$\bar{5}_{2,3}$	$5_H, \bar{5}_H$	Z
$\mathrm{U}(1)_F$	4	2	0	$p+1$	p	0	-1
Rp	-1			-1		1	1

Table 1. U(1) flavor symmetry and the R-parity charges of different MSSM multiplets. The subscripts denote the flavor index. Values $p = 0, 1, 2$ approximately explain the fermion mass hierarchy. The charge assignment for the MSSM multiplets will remain same throughout the paper.

This immediately forbids mass mixing terms between the messenger and matter multiplets at all orders of the perturbation theory. While for models where the interaction term involves one messenger multiplet and two MSSM multiplets, a consistent messenger parity cannot be constructed. This will lead to messenger-matter mixing either at tree level or at higher orders, leading to non-trivial contributions to the soft spectrum. We will discuss the consequences of this for specific scenarios in the next section.

3 The soft breaking masses

We assume that supersymmetry is broken due to some hidden sector dynamics that can be parametrized into a spurion (X) vev. The messengers which are charges under the MSSM gauge groups, couple to these spurion fields and communicate supersymmetry breaking to the visible sector through the usual gauge couplings. Usually the messengers are assumed to have some messenger parity that prevents these fields from mixing or interacting with the visible sector. In this paper we will relax this and consider all possible interaction terms between the visible and hidden sectors. The entire superpotential can be schematically written as,

$$W = W_{\mathrm{MSSM}}(\phi_{\mathrm{vis}}) + \lambda_\phi X \phi_m \tilde{\phi}_m + W_{\mathrm{int}}(\phi_m, \phi_{\mathrm{vis}}), \tag{3.1}$$

where the spurion gets a supersymmetry breaking vev $\langle X \rangle = M + \theta^2 F$ and $\{\phi_m, \tilde{\phi}_m\}$ are vector like pair of the messenger fields and ϕ_{vis} are the usual MSSM chiral supermultiplets. We define the messenger scale $M_{\mathrm{mess}} = \lambda_\phi M$ and $\Lambda = F/M$. The loop integrals can be expressed in terms of the dimensionless parameter $x = \Lambda/M_{\mathrm{mess}}$.

The most general messenger sector however can involve mass terms for the messenger fields of the form $m_{ij}\phi_i\tilde{\phi}_j$. This can introduce a new dimensionful parameter other than the messenger scale. The ensuing complication in determination of the soft spectrum through the wave-function renormalization technique lead to the so called $m_{H_u} - A_u$ problem [14] in models where H_u and/or H_d couple with messengers. In the models discussed in this paper we will find the these mass terms are either absent or suppressed by a factor ϵ^a, where a is the flavor charge of the spurion field X. We can in principle make a large to suppress these terms effectively. We will neglect these terms in our calculations below.

The soft masses get the usual contributions from the gauge interactions of the messenger fields at one loop for the gaugino masses and two loops for the scalar masses, given

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by [4],

$$M_r = d\frac{\alpha_r}{4\pi}g(x)\,\Lambda,$$
$$M_{\phi_i}^2 = 2d\sum_{r=1,2,3}\left[\left(\frac{\alpha_r}{4\pi}\right)^2 C_r(i)\right]f(x)\,\Lambda^2, \tag{3.2}$$

where d is the Dynkin index and is 1 for messengers in the $5\oplus\bar{5}$ and 3 for messengers in $10\oplus\overline{10}$. The $C_r(i)$'s are the usual Casimir invariants for the representation i and,

$$f(x) = \frac{1+x}{x^2}\left[\log(1+x) - 2\mathrm{Li}_2(x/[1+x]) + \frac{1}{2}\mathrm{Li}_2(2x/[1+x])\right] + (x\to -x),$$
$$g(x) = \frac{1}{x^2}\left[(1+x)\log(1+x) + (1-x)\log(1-x)\right]. \tag{3.3}$$

Note that two loop contributions to the gaugino masses were computed in [18, 19]. For messenger scale beyond 100 TeV the corrections are at a few percent level. We will neglect this small correction in our numerical calculations.

Now we turn to the contribution of the messenger-matter interaction term in eq. (3.1). In this paper we will only consider interaction terms of the form $W_{\rm int} = \lambda_{ijk}\phi_i\phi_j\phi_k$, where at least one chiral multiplet from both the messenger and the visible sector are present. The contribution at one loop level to a field ϕ_i belonging to the MSSM can be directly computed and is given by,

$$\delta M_i^2\big|_{1-\rm loop} = -C_{ijk}\frac{|\lambda_{ijk}|^2}{96\pi^2}x^2\Lambda^2 h(x), \tag{3.4}$$

where C_{ijk} is the multiplicity factor that measures the effective number of messenger fields that the MSSM supermultiplet ϕ_i couples to through the superpotential coupling λ_{ijk} including the appropriate group theoretic factors.[2] The expression of $h(x)$ is as follows [14, 16],

$$h(x) = 3\frac{(x-2)\log(1-x) - (x+2)\log(1+x)}{x^4}. \tag{3.5}$$

Let us make the following observations regarding the one loop contribution: (i) note that $h(x\to 0)\to 1$ and thus (ii) the one loop contribution decouples with the messenger scale $(\delta M_i^2(x\to 0)\big|_{1-\rm loop}\to 0)$, (iii) the contribution is always negative.

It is easier to use the wave function renormalization techniques [27, 28] in order to compute the the two loop contributions. They can be written in terms of the anomalous dimension and the β functions for the Yukawa coupling, below and above the messenger scales. Adaptations of the generic framework for messenger-matter interaction terms were computed in [11, 28], we quote them for the sake of completeness,

$$\delta M_i^2\big|_{2-\rm loop} = \frac{1}{2}\sum_\lambda\left[\beta_\lambda^+\frac{\partial(\Delta\gamma_i)}{\partial\lambda} - \Delta\beta_\lambda\frac{\partial(\gamma_i^-)}{\partial\lambda}\right]_{M_{\rm mess}}\Lambda^2, \tag{3.6}$$
$$A_{ijk}\big|_{1-\rm loop} = -\left(\Delta\gamma_i + \Delta\gamma_j + \Delta\gamma_k\right)\,h_A(x)\,\Lambda, \tag{3.7}$$

[2]For instance, in the interaction $\lambda Q U_m^c H_U$, C_{ijk} for Q is 1 whereas C_{ijk} for H_U is 3 (color factor). In contrast for $\lambda Q_m U_m^c H_U$, C_{ijk} for H_U is now 6.

where, $\Delta X|_{M_{\rm mess}} = [X^+ - X^-]|_{M_{\rm mess}} = [X(M_{\rm mess}+\delta) - X(M_{\rm mess}-\delta)]|_{\delta\to 0}$, $\beta_\lambda = d\lambda/dt$ and

$$h_A(x) = \frac{1}{2x}\log\left(\frac{1+x}{1-x}\right). \tag{3.8}$$

As has been pointed out recently in [29] one should be careful to interpret the above formula for models where a consistent messenger parity cannot be imposed and thus leads to the possibility of mixing between the messenger and MSSM multiplets. These kinetic mixing terms can be removed by an unitary rotation. The above formula gives correct results for this case [30] assuming a non-standard definition of the corresponding anomalous dimensions. We have checked that in all the cases where such mixing can arise, our results are in agreement with the treatment prescribed in [29].

We will now compute the mass spectrum of each of the models using the generic expressions for the soft masses introduced in this section. We will only indicate the new contributions arising from the interaction terms in addition to the usual contributions given in eq. (3.2). For each model we will indicate the R-parity and the flavor charges for the messenger sector. The corresponding charges for the MSSM multiplets are given in table 1. We will assume that wherever the interaction terms include two multiplets, a messenger parity is imposed under which the messenger sector is odd while the multiplets of the visible sector are even.

3.1 Only singlets

In these scenarios one has to assume that the messenger sector in addition to the singlets, also has the usual messenger fields charged under the MSSM gauge group. As a definite choice we will assume that along with the singlet the messenger sector consists of a single $5\oplus\bar{5}$ vector pair of chiral messengers (ϕ_m $\tilde{\phi}_m$). However a messenger parity prevents them from directly coupling with the visible sector. These spectator messengers will contribute to the soft masses through usual gauge interactions according to eq. (3.2).

3.1.1 Model 1

$\mathbf{5_H\bar{5}_H1_m}$: considering that the singlet (1_m) is even under R-parity, one obtains the following messenger-matter mixing superpotential term [11],

$$W_{\rm int} = \lambda H_u H_d S_m. \tag{3.9}$$

The new contributions to the soft masses can be read off from eqs. (3.4) and (3.6). They are given by,

$$\begin{aligned}
\delta M_Q^2 &= -\frac{\alpha_\lambda}{16\pi^2}(\alpha_t+\alpha_b)\Lambda^2, \qquad \delta M_{U^c}^2 = -\frac{\alpha_t\alpha_\lambda}{8\pi^2}\Lambda^2, \qquad \delta M_{D^c}^2 = -\frac{\alpha_b\alpha_\lambda}{8\pi^2}\Lambda^2,\\
\delta M_L^2 &= -\frac{\alpha_\lambda\alpha_\tau}{16\pi^2}\Lambda^2, \qquad \delta M_{E^c}^2 = -\frac{\alpha_\lambda\alpha_\tau}{8\pi^2}\Lambda^2,\\
\delta M_{H_u}^2 &= \left[-\frac{\alpha_\lambda}{24\pi}x_1^2 h(x_1) + \frac{\alpha_\lambda}{16\pi^2}\left(4\alpha_\lambda+\alpha_\tau+3\alpha_b-3\alpha_2-\frac{3}{5}\alpha_1\right)\right]\Lambda^2,\\
\delta M_{H_d}^2 &= \left[-\frac{\alpha_\lambda}{24\pi}x_1^2 h(x_1) + \frac{\alpha_\lambda}{16\pi^2}\left(4\alpha_\lambda+3\alpha_t-3\alpha_2-\frac{3}{5}\alpha_1\right)\right]\Lambda^2,\\
A_t &= A_b = A_\tau = -h_A(x)\frac{\alpha_\lambda}{4\pi}\Lambda,
\end{aligned} \tag{3.10}$$

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where $\alpha_{\lambda_i} = \lambda_i^2/4\pi$ and the subscripts have their usual meaning. We will follow this convention through out the paper. Unfortunately as commented in [11], one can anticipate an anomalous contribution to the $\mu - B_\mu$ parameters in this model.

3.1.2 Model 2

$\mathbf{5_H \bar{5} 1_m}$ **:** the other possibility here is to assume that the messenger field is odd under the R-parity. However in this case we get contributions to the neutrino mass arising through type I see-saw mechanism[3][31]. This puts a lower bound on the messenger scale at $\sim 10^{10} GeV$. A way to evade this to is consider another singlet field $\tilde{S}$ and impose non-trivial flavor charges:

$$\mathrm{U}(1)_F(1_m, \tilde{1}_m, 5_m, \bar{5}_m, X) = (-p, -a+p, -(a+q), q, a). \tag{3.11}$$

We obtain the following superpotential through which the singlet messenger field interact with the MSSM multiplets,

$$W_{\text{int}} = \lambda H_u L_3 S_m. \tag{3.12}$$

We made a field redefinition so that only the third family can interact with the messenger field. We considered a and q to be large positive numbers so that all the non-renormalizable terms in the messenger-matter mixing sector are highly suppressed. As is clear from the above expression, the field S_m must carry a unit lepton number. The new contributions to the soft masses from the given interaction term is given by,

$$\begin{aligned}
\delta M_{\tilde{Q}}^2 &= -\frac{\alpha_\lambda \alpha_t}{16\pi^2}\Lambda^2, \quad \delta M_{\tilde{U}^c}^2 = -\frac{\alpha_\lambda \alpha_t}{8\pi^2}\Lambda^2, \quad \delta M_{\tilde{E}^c}^2 = -\frac{\alpha_\lambda \alpha_\tau}{8\pi^2}\Lambda^2, \quad \delta M_{H_d}^2 = -\frac{\alpha_\lambda \alpha_\tau}{16\pi^2}\Lambda^2, \\
\delta M_{\tilde{L}}^2 &= \left[-\frac{\alpha_\lambda}{24\pi}x^2 h(x) + \frac{\alpha_\lambda}{16\pi^2}\left(4\alpha_\lambda + 3\alpha_t - 3\alpha_2 - \frac{3}{5}\alpha_1\right)\right]\Lambda^2, \\
\delta M_{H_u}^2 &= \left[-\frac{\alpha_\lambda}{24\pi}x^2 h(x) + \frac{\alpha_\lambda}{16\pi^2}\left(4\alpha_\lambda + \alpha_\tau - 3\alpha_2 - \frac{3}{5}\alpha_1\right)\right]\Lambda^2, \\
A_t &= A_\tau = -h_A(x)\frac{\alpha_\lambda}{4\pi}\Lambda.
\end{aligned} \tag{3.13}$$

3.2 Only $\mathbf{5 \oplus \bar{5}}$

In this section we look at models where the messengers are in the $5 \oplus \bar{5}$ representations. Depending on the choice of symmetries various invariants can be constructed. We will now study the possible terms in turn and compute the new contributions to the scalar masses. Again the usual contributions are given by eq. (3.2) where we set the Dynkin index $d = 1$. Unless mentioned we will assume that the number of generations of messenger is one.

3.2.1 Model 3

$\mathbf{0 \bar{5}_H \bar{5}_m}$ **:** let us consider that the messengers are odd under the R-parity and the following flavor charges are imposed,

$$\mathrm{U}(1)_F(5_m, \bar{5}_m, X) = (-a, 0, a). \tag{3.14}$$

[3] See [32, 33] for realistic models of neutrino mass within the GMSB framework that utilize messenger-matter interactions.

We obtain the following interaction term in the superpotential,

$$W_{\text{int}} = \lambda_q Q_3 D^c_m H_d + \lambda_e E^c_3 H^m_d H_d. \tag{3.15}$$

The absence of messenger parity allows the operator $\bar{5}_H\, 5_m\, X$ to be consistent with all other symmetries imposed. This can be absorbed in the superpotential above by a basis change. However there is still a loop level mixing between $\bar{5}_H$ and $\bar{5}_m$. This can be rotated away at the one loop order but contributes non trivially to two loop corrections of the soft spectrum. The contributions of these mixing terms to the trilinear coupling are numerically negligible compared to the unsuppressed one loop contributions and thus can be neglected. We extend the analysis in [12, 13] by including all the two loop contributions to the scalar masses and considering the effect of one loop mixing between the messenger and matter multiplets. The expressions are given by,

$$\begin{aligned}
\delta M^2_Q &= \Bigg[-\frac{\alpha_q}{24\pi}x^2 h(x) \\
&\qquad +\frac{\alpha_q}{16\pi^2}\left[6\alpha_q + 4\alpha_b + \alpha_\tau + \alpha_e - \frac{16}{3}\alpha_3 - 3\alpha_2 - \frac{7}{15}\alpha_1\right] - \frac{\alpha_b\alpha_e}{16\pi^2}\Bigg]\Lambda^2, \\
\delta M^2_{U^c} &= -\frac{\alpha_q\alpha_t}{8\pi^2}\Lambda^2, \\
\delta M^2_{D^c} &= -\frac{\alpha_b}{8\pi^2}(\alpha_e + 4\alpha_q)\Lambda^2, \\
\delta M^2_L &= -\frac{3\alpha_\tau}{16\pi^2}(\alpha_e + \alpha_q)\Lambda^2, \\
\delta M^2_{E^c} &= \left[-\frac{\alpha_e}{12\pi}x^2 h(x) + \frac{\alpha_e}{16\pi^2}\left[8\alpha_e + 6\alpha_q + 6\alpha_b + 4\alpha_\tau - 6\alpha_2 - \frac{18}{5}\alpha_1\right] - \frac{3\alpha_q}{8\pi^2}\alpha_\tau\right]\Lambda^2, \\
\delta M^2_{H_d} &= \Bigg[-\frac{\alpha_e + 3\alpha_q}{24\pi}x^2 h(x) + \frac{\alpha_e}{16\pi^2}\left[4\alpha_e + 2\alpha_\tau - 3\alpha_2 - \frac{9}{5}\alpha_1\right] \\
&\qquad +\frac{\alpha_q}{16\pi^2}\left[18\alpha_q + 3\alpha_t + 12\alpha_b + 6\alpha_e - 16\alpha_3 - 9\alpha_2 - \frac{7}{5}\alpha_1\right]\Bigg]\Lambda^2, \\
\delta M^2_{H_u} &= -\frac{3\alpha_q\alpha_t}{16\pi^2}\Lambda^2, \\
A_t &= -\frac{\alpha_q}{4\pi}\Lambda, \\
A_b &= -\frac{\alpha_e + 4\alpha_q}{4\pi}\Lambda, \\
A_\tau &= -\frac{3\,(\alpha_e + \alpha_q)}{4\pi}\Lambda.
\end{aligned} \tag{3.16}$$

3.2.2 Model 4

10105_m : if we consider the messengers (5_m and $\bar{5}_m$) are even under the R-parity we can have two different invariants $10\,10\,5_m$ and $10\,\bar{5}\,\bar{5}_m$ that are possible [13]. However depending on the assignment of the flavor charges, one or the other might may become more dominant. We will consider by turn the two extreme scenarios where only one of the invariants dominates. Considering the flavor charges,

$$\mathrm{U}(1)_F(5_m, \bar{5}_m, X) = (0, a, -a), \tag{3.17}$$

the dominating part of the superpotential is given by,

$$W_{\rm int} = \lambda_q Q_3 U_3^c H_u^m + \lambda_u U_3^c \tilde{D}_m^c E_3^c. \tag{3.18}$$

The contributions to the soft scalar masses are given by,

$$\begin{aligned}
\delta M_Q^2 &= \left[-\frac{\alpha_q}{24\pi}x^2 h(x) + \frac{\alpha_q}{16\pi^2}\left[6\alpha_q + 6\alpha_t + \alpha_u - \frac{16}{3}\alpha_3 - 3\alpha_2 - \frac{13}{15}\alpha_1\right] - \frac{\alpha_t\alpha_u}{16\pi^2}\right]\Lambda^2,\\
\delta M_{U^c}^2 &= \left[-\frac{2\alpha_q+\alpha_u}{24\pi}x^2 h(x) + \frac{\alpha_q}{8\pi^2}\left(6\alpha_q + \alpha_b + 6\alpha_t + 2\alpha_u - \frac{16}{3}\alpha_3 - 3\alpha_2 - \frac{13}{15}\alpha_1\right)\right.\\
&\qquad\left. + \frac{\alpha_u}{16\pi^2}\left(5\alpha_u + 2\alpha_\tau - \frac{16}{3}\alpha_3 - \frac{28}{15}\alpha_1\right)\right]\Lambda^2, \qquad (3.19)\\
\delta M_{D^c}^2 &= -\frac{\alpha_b\alpha_q}{8\pi^2}\Lambda^2,\\
\delta M_L^2 &= -\frac{3\alpha_u\alpha_\tau}{16\pi^2}\Lambda^2,\\
\delta M_{E^c}^2 &= \left[-\frac{\alpha_u}{8\pi}x^2 h(x) + \frac{3\lambda_u}{16\pi^2}\left(5\lambda_u + 2\lambda_t + 2\lambda_q - \frac{16}{3}\alpha_3 - \frac{28}{15}\alpha_1\right)\right]\Lambda^2,\\
\delta M_{H_u}^2 &= -3\alpha_t\frac{3\alpha_q+\alpha_u}{16\pi^2}\Lambda^2, \quad \delta M_{H_d}^2 = -3\frac{\alpha_b\alpha_q + \alpha_u\alpha_\tau}{16\pi^2}\Lambda^2,\\
A_t &= -\frac{3\alpha_q+\alpha_u}{4\pi}\Lambda,\\
A_b &= -\frac{\alpha_q}{4\pi}\Lambda,\\
A_\tau &= -\frac{3\alpha_u}{4\pi}\Lambda.
\end{aligned}$$

The lack of messenger parity in this case can lead to tree level mass terms of the form $m'\bar{5}_H\bar{5}_m$. It is expected that $m' \sim \mu$, where μ is the usual dimensionful parameter in the MSSM superpotential. This mass cannot be suppressed without suppressing the operators in eq. (3.18). The phenomenology of this model though not identical, closely resembles the scenario studied in [16].

3.2.3 Model 5

$10 5 \bar{5} \bar{5}_m$: this is the other possibility with even R-parity for the messengers. One can distinguish it from Model 4 by considering a different flavor symmetry given by,

$$\mathrm{U}(1)_F(5_m, \bar{5}_m, X) = (a+p, -p, -a). \tag{3.20}$$

However note that flavor charges assigned in table 1 do not distinguish between $\bar{5}_2$ and $\bar{5}_3$. So this is not enough to ensure that only third generation will couple strongly with the messenger sector. This leads to considerable contributions to FCNC that constraints the messenger scale. An additional complication is related to the mass terms of the form $m'5_H\bar{5}_m$ where $m' \sim \epsilon^p\mu$. Worse, the interaction can lead to rapid proton decay and suppression of the order of ϵ^5 is not strong enough to comply with present experimental bounds. However one can consider flavor symmetries which are uncorrelated to the origin of the fermion mass hierarchy that can enable the required suppression. We adopt the paradigm that this is possible, assuming this we can write down the superpotential as,

$$W_{\rm int} = \lambda_q Q_3 D_3^c H_d^m + \lambda_l Q_3 L_3 D_m^c + \lambda_u U_3^c D_3^c D_m^c + \lambda_e L_3 E_3^c H_d^m. \tag{3.21}$$

The contributions to the soft masses can be written as,

$$\begin{aligned}
\delta M_Q^2 &= \Bigg[-\frac{x^2 h(x)}{24\pi}(\alpha_q+\alpha_l) + \frac{\alpha_l}{16\pi^2}\left[6\alpha_l+\alpha_e+\alpha_u+\alpha_\tau-\frac{16}{3}\alpha_3-3\alpha_2-\frac{7}{15}\alpha_1\right] \\
&\quad +\frac{\alpha_q}{16\pi^2}\left[6\alpha_q+6\alpha_b+2\sqrt{\alpha_b\alpha_\tau\alpha_e/\alpha_q}+2\alpha_l+\alpha_e+\alpha_u-\frac{16}{3}\alpha_3-3\alpha_2-\frac{7}{15}\alpha_1\right] \\
&\quad -\frac{\alpha_u(\alpha_b+\alpha_t)}{16\pi^2}\Bigg]\Lambda^2, \\
\delta M_{U^c}^2 &= \left[-\frac{x^2 h(x)}{24\pi}\alpha_u + \frac{\alpha_u}{16\pi^2}\left[3\alpha_u+2\alpha_q+2\alpha_l+2\alpha_b-8\alpha_3-\frac{4}{5}\alpha_1\right]-\frac{\alpha_t(\alpha_q+\alpha_l)}{8\pi^2}\right]\Lambda^2, \\
\delta M_{D^c}^2 &= \Bigg[-\frac{x^2 h(x)}{24\pi}(2\alpha_q+\alpha_u)+\frac{\alpha_q}{16\pi^2}\Big[12\alpha_q+2\alpha_l+2\alpha_e+2\alpha_t+12\alpha_b \\
&\qquad +4\sqrt{\alpha_b\alpha_\tau\alpha_e/\alpha_q}-\frac{32}{3}\alpha_3-6\alpha_2-\frac{14}{15}\alpha_1\Big] \\
&\quad +\frac{\alpha_u}{16\pi^2}\left[3\alpha_u+2\alpha_t+4\alpha_q+2\alpha_l-8\alpha_3-\frac{4}{5}\alpha_1\right]-\frac{\alpha_b\alpha_l}{8\pi^2}\Bigg]\Lambda^2, \qquad (3.22)\\
\delta M_L^2 &= \Bigg[-\frac{x^2 h(x)}{24\pi}\alpha_e+\frac{\alpha_e}{16\pi^2}\left[4\alpha_e+3\alpha_q+2\alpha_\tau+6\sqrt{\alpha_b\alpha_\tau\alpha_q/\alpha_e}-3\alpha_2-\frac{9}{5}\alpha_1\right] \\
&\quad +\frac{\alpha_l}{16\pi^2}\left[18\alpha_l+6\alpha_e+3\alpha_b+3\alpha_q+3\alpha_t+3\alpha_u-16\alpha_3-9\alpha_2-\frac{7}{5}\alpha_1\right]\Bigg]\Lambda^2, \\
\delta M_{E^c}^2 &= \Bigg[-\frac{x^2 h(x)}{12\pi}\alpha_e+\frac{\alpha_e}{8\pi^2}\Big[4\alpha_e+3\alpha_l+3\alpha_q-3\alpha_2+2\alpha_\tau \\
&\qquad +6\sqrt{\alpha_b\alpha_\tau\alpha_q/\alpha_e}-\frac{9}{5}\alpha_1\Big]-\frac{3\alpha_l\alpha_\tau}{8\pi^2}\Bigg]\Lambda^2 \\
\delta M_{H_u}^2 &= -\frac{3\alpha_t}{16\pi^2}(\alpha_l+\alpha_q+\alpha_u)\Lambda^2, \\
\delta M_{H_d}^2 &= -\left[\frac{3\alpha_b}{16\pi^2}(\alpha_l+3\alpha_q+\alpha_u)+\frac{3\alpha_\tau}{16\pi^2}(\alpha_e+\alpha_l)\right]\Lambda^2, \\
A_t &= -\frac{\alpha_l+\alpha_q+\alpha_u}{4\pi}\Lambda, \quad A_b = -\frac{\alpha_l+3\alpha_q+\alpha_u}{4\pi}\Lambda, \quad A_\tau = -\frac{3\,(\alpha_e+\alpha_l)}{4\pi}\Lambda.
\end{aligned}$$

Note that in [13] a scenario that includes Model 4 and Model 5 was studied. Low energy flavor observables were utilized to constraint different cmbinations of the interaction Yukawa couplings. Considering that most flavor constraints are restrictive only for the first two fermion generations, the flavor symmetries imposed in our analysis will certainly relax some of these bounds. However a detailed study of these constraints, that warrants careful attention, is beyond the scope of this paper.

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3.2.4 Model 6

$\mathbf{10\bar{5}_m\bar{5}_m}$ **:** we expand the R-symmetry from Z_2 to Z_4. The parity of the MSSM multiplets are still given by table 1, whereas the messenger sector now has the following parity,

$$Rp(5_m, \bar{5}_m) = (i5_m, -i\bar{5}_m). \tag{3.23}$$

On top of this we impose the following flavor charges,

$$\mathrm{U}(1)_F(5_m, \bar{5}_m, X) = (-a, 0, a), \tag{3.24}$$

This ensures that we have the following dominant superpotential,

$$W_{\mathrm{int}} = \lambda Q_3 D_m^c H_d^m. \tag{3.25}$$

The new contributions are given by,

$$\begin{aligned}
\delta M_Q^2 &= \left[-\frac{\alpha_\lambda}{12\pi}x^2 h(x) + \frac{\alpha_\lambda}{16\pi^2}\left(6\alpha_\lambda - \frac{16}{3}\alpha_3 - 3\alpha_2 - \frac{7}{15}\alpha_1\right)\right]\Lambda^2, \\
\delta M_{U^c}^2 &= -\frac{\alpha_\lambda\alpha_t}{8\pi^2}\Lambda^2, \qquad \delta M_{D^c}^2 = -\frac{\alpha_\lambda\alpha_b}{8\pi^2}\Lambda^2, \\
\delta M_{H_u}^2 &= -\frac{3\alpha_\lambda\alpha_t}{16\pi^2}\Lambda^2, \qquad \delta M_{H_d}^2 = -\frac{3\alpha_\lambda\alpha_b}{16\pi^2}\Lambda^2, \\
A_t &= A_b = -\frac{\alpha_\lambda}{4\pi}\Lambda.
\end{aligned} \tag{3.26}$$

3.2.5 Models 7 & 8

$\mathbf{10\bar{5}_{jm}\bar{5}_{km}}$ **:** sticking to the same invariant as Model 6 an interesting scenario develops when we expand the number of messenger generations to two. The drastic change is more than a simple duplication of the results in Model 6. The form of the Lagrangian and thus the soft spectrum, depends on the symmetries we impose on the theory. In this regard we will discuss two slight variants:

- **Model 7:** We consider a Z_4 R-parity with,

  $$Rp(5_{jm}, \bar{5}_{jm}) = (i5_{jm}, -i\bar{5}_{jm}), \quad j = 1, 2, \tag{3.27}$$

  and flavor charges,

  $$\mathrm{U}(1)_F(5_{jm}, \bar{5}_{jm}, X) = (-a, 0, a). \tag{3.28}$$

- **Model 8:** We consider the following Z_2 R-parity,

  $$Rp(5_{1m}, \bar{5}_{1m}, 5_{2m}, \bar{5}_{2m}) = (-5_{1m}, -\bar{5}_{1m}, 5_{2m}, \bar{5}_{2m}), \tag{3.29}$$

  and the following flavor symmetry:

  $$\mathrm{U}(1)_F(5_{1m}, \bar{5}_{1m}, 5_{2m}, \bar{5}_{2m}, X) = (a - b, b, a + b, -b, -a). \tag{3.30}$$

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With these symmetries we obtain the following form of the superpotential for Model 7,

$$\begin{aligned} W_{\text{int}} &= (\lambda_{11} Q_3 D^c_{1m} H_d^{1m} + \lambda_{22} Q_3 D^c_{2m} H_d^{2m}) + \lambda_{12} Q_3 D^c_{1m} H_d^{2m} \\ &\quad + \lambda_{21} Q_3 D^c{}_2 m H_d^{1m} + \lambda_u U_3^c D^c_{1m} D^c_{2m} + \lambda_e E_3^c H_d^{1m} H_d^{2m}. \end{aligned} \tag{3.31}$$

The new contributions to the soft spectrum for Model 7 are given by,

$$\begin{aligned} \delta M_Q^2 &= \Bigg[-\frac{x^2 h(x)}{12\pi}(\alpha_{11}+\alpha_{12}+\alpha_{21}+\alpha_{22}) \\ &\quad + \frac{\alpha_{11}}{16\pi^2}\Bigg[\left(3\alpha_{11}+\alpha_e+\alpha_u-\frac{16}{3}\alpha_3-3\alpha_2-\frac{7}{15}\alpha_1\right) \\ &\qquad + (\alpha_{11}\to\alpha_{12})+(\alpha_{11}\to\alpha_{21})+(\alpha_{11}\to\alpha_{22})\Bigg] - \frac{\alpha_t\alpha_u}{16\pi^2} \\ &\quad + \frac{1}{8\pi^2}(3\alpha_{11}\alpha_{12}+4\alpha_{11}\alpha_{21}+\alpha_{11}\alpha_{22}+\alpha_{12}\alpha_{21}+4\alpha_{12}\alpha_{22}+3\alpha_{21}\alpha_{22})\Bigg]\Lambda^2, \\ \delta M_{U^c}^2 &= \Bigg[-\frac{x^2 h(x)}{12\pi}\alpha_u - \frac{\alpha_t}{8\pi^2}(\alpha_{11}+\alpha_{12}+\alpha_{21}+\alpha_{22}) \\ &\quad + \frac{\alpha_u}{16\pi^2}\left[3\alpha_u+2(\alpha_{11}+\alpha_{12}+\alpha_{21}+\alpha_{22})-8\alpha_3-\frac{4}{5}\alpha_1\right]\Bigg]\Lambda^2, \\ \delta M_{D^c}^2 &= \left[-\frac{\alpha_b}{8\pi^2}(\alpha_{11}+\alpha_{12}+\alpha_{21}+\alpha_{22})\right]\Lambda^2, \\ \delta M_L^2 &= \left[-\frac{\alpha_e\alpha_\tau}{8\pi^2}\right]\Lambda^2, \\ \delta M_{E^c}^2 &= \left[-\frac{x^2 h(x)}{6\pi}\alpha_e + \frac{\alpha_e}{8\pi^2}\left[4\alpha_e+3(\alpha_{11}+\alpha_{12}+\alpha_{21}+\alpha_{22})-3\alpha_2-\frac{9}{5}\alpha_1\right]\right]\Lambda^2, \\ \delta M_{H_u}^2 &= \left[-\frac{3\alpha_t}{16\pi^2}(\alpha_{11}+\alpha_{12}+\alpha_{21}+\alpha_{22}+\alpha_u)\right]\Lambda^2, \\ \delta M_{H_d}^2 &= \left[-\frac{3\alpha_t}{16\pi^2}(\alpha_{11}+\alpha_{12}+\alpha_{21}+\alpha_{22}+\alpha_u)-\frac{\alpha_e\alpha_\tau}{8\pi^2}\right]\Lambda^2, \\ A_t &= \left[-\frac{\alpha_{11}+\alpha_{12}+\alpha_{21}+\alpha_{22}+\alpha_u}{4\pi}\right]\Lambda, \\ A_b &= \left[-\frac{\alpha_{11}+\alpha_{12}+\alpha_{21}+\alpha_{22}}{4\pi}\right]\Lambda, \\ A_\tau &= \left[-\frac{\alpha_e}{2\pi}\right]\Lambda. \end{aligned} \tag{3.32}$$

The corresponding superpotential and the soft spectrum for Model 8 can be obtained by setting $\lambda_{11} = \lambda_{22} = 0$.

3.3 Only $\mathbf{10} \oplus \overline{\mathbf{10}}$

In this section we will collect the possible messenger-matter interaction terms possible assuming that the messengers are in a vector like representation of $10 \oplus \overline{10}$.

3.3.1 Model 9

$\mathbf{10 10_m 5_H + 10_m 10_m 5_H}$: consider that the messenger fields are odd under R-parity ($Rp(10_m) = -10_m$) and have the following flavor charges,

$$\mathrm{U}(1)_x(10_m, \bar{10}_m, X) = (0, -a, a) \tag{3.33}$$

The invariant $10_m\ \bar{5}_{2,3}\ \bar{5}_H$ has a coupling which is suppressed by the flavor factor and is at least as small as the λ_τ Yukawa. We thus neglect it from the discussion and obtain the following interaction superpotential [12],

$$W_{\rm int} = \lambda_q Q_3 U_m^c H_u + \lambda_u Q^m U_3^c H_u + \lambda_h Q_m U_m^c H_u. \tag{3.34}$$

The lack of messenger parity in this scenario again manifests into one loop mixing between messenger and matter multiplets. The one and two loop contributions to the soft masses including the mixing effects are given by,

$$\begin{aligned}
\delta M_Q^2 &= \Bigg[-\frac{\alpha_q}{24\pi} x^2 h(x) + \frac{\alpha_q}{16\pi^2}\left(6\alpha_q + 3\alpha_u + 5\alpha_h + 3\alpha_t - \frac{16}{3}\alpha_3 - 3\alpha_2 - \frac{13}{15}\alpha_1\right) \\
&\qquad - \frac{\alpha_t(5\alpha_u + 3\alpha_h)}{16\pi^2}\Bigg]\Lambda^2, \\
\delta M_{U^c}^2 &= \Bigg[-\frac{\alpha_u}{12\pi} x^2 h(x) + \frac{\alpha_u}{8\pi^2}\left(6\alpha_u + 3\alpha_q + 4\alpha_h + 3\alpha_t - \frac{16}{3}\alpha_3 - 3\alpha_2 - \frac{13}{15}\alpha_1\right) \\
&\qquad - \frac{\alpha_t(4\alpha_q + 3\alpha_h)}{8\pi^2}\Bigg]\Lambda^2, \\
\delta M_{D^c}^2 &= \left[-\frac{\alpha_b\alpha_q}{8\pi^2}\right]\Lambda^2, \\
\delta M_{H_d}^2 &= \left[-\frac{3\alpha_b\alpha_q}{16\pi^2}\right]\Lambda^2, \\
\delta M_{H_u}^2 &= \Bigg[-\frac{\alpha_q + \alpha_u + 2\alpha_h}{8\pi} x^2 h(x) - \frac{\alpha_q + \alpha_u + \alpha_h}{16\pi^2}\left(16\alpha_3 + 9\alpha_2 + \frac{13}{5}\alpha_1\right) \\
&\qquad + \frac{3\alpha_q}{16\pi^2}\left(6\alpha_q + 10\alpha_h + \alpha_b + 5\alpha_t\right) + \frac{9\alpha_u}{8\pi^2}\left(\alpha_u + \alpha_q + \frac{4}{3}\alpha_h + \frac{4}{3}\alpha_t\right) + \frac{9\alpha_h^2}{8\pi^2}\Bigg]\Lambda^2, \\
A_t &= -\frac{3\alpha_h + 4\alpha_q + 5\alpha_u}{4\pi}\Lambda, \qquad A_b = -\frac{\alpha_q}{4\pi}\Lambda.
\end{aligned} \tag{3.35}$$

3.3.2 Model 10

$\mathbf{10_m\, 10_m\, 5_H}$: the other alternative is to consider that the messengers are even under R-parity ($Rp(10_m) = 10_m$). In this case the messengers can only couple to the Higgs multiplets in the MSSM sector and thus we select the flavor charges of the messenger sector to be zero. The corresponding superpotential is given by [14],

$$W_{\rm int} = \lambda Q_m U_m^c H_u, \tag{3.36}$$

and the contributions to the soft masses are given by,

$$\begin{aligned}
\delta M_Q^2 &= -\frac{3\alpha_t\alpha_\lambda}{16\pi^2}\Lambda^2, \\
\delta M_{U^c}^2 &= -\frac{3\alpha_t\alpha_\lambda}{8\pi^2}\Lambda^2, && (3.37) \\
\delta M_{H_u}^2 &= \left(\frac{\alpha_\lambda}{8\pi^2}\left[9\alpha_\lambda - 8\alpha_3 - \frac{9}{2}\alpha_2 - \frac{13}{10}\alpha_1\right] - \frac{\alpha_\lambda}{4\pi}x^2h(x)\right)\Lambda^2, && (3.38) \\
A_t &= -\frac{3\alpha_\lambda}{4\pi}\Lambda.
\end{aligned}$$

3.4 Mixed messenger models

It is possible that the messenger sector is composed of messenger fields that are in different complete representations of the GUT group SU(5). Actually this is implicitly assumed in Models 1 and 2. In that case we can have scenarios where more than one of them simultaneously interact with the visible sector. There are a large number of possibilities in this class, mainly coming from a combination of two or more models already discussed, possibly augmented by some new terms. A study of all these models are beyond the scope of this study and we will restrict ourselves to models that lead to entirely new messenger-matter interaction terms.

3.4.1 Model 11

Let us consider a scenario where messengers in 1 and 5 representations interact simultaneously with the MSSM multiplets. The only new interaction term that can be envisaged in this case is $\mathbf{5_H}\,\mathbf{\bar{5}_m}\,\mathbf{1_m}$. In order to prevent other terms from showing up we impose the following symmetries. Again we conjecture the existence of a second singlet $\tilde{S}_m$. We impose the following assignment of R-parity,

$$Rp(5_m, \bar{5}_m, 1_m, \tilde{1}_m) = (-5_m, -\bar{5}_m, -1_m, -\tilde{1}_m), \tag{3.39}$$

and the following flavor charges,

$$\mathrm{U}(1)_F(5_m, \bar{5}_m, 1_m, \tilde{1}_m, X) = (-a-b, b, -b, -a+b, a) \text{ with } a > b. \tag{3.40}$$

We obtain the following unsuppressed terms in the superpotential [14],

$$W_{\text{int}} = H_u H_d^m S_m, \tag{3.41}$$

and the contributions to the soft masses are given by,

$$\begin{aligned}
\delta M_Q^2 &= -\frac{\alpha_t\alpha_\lambda}{16\pi^2}\Lambda^2, \quad \delta M_{U^c}^2 = -\frac{\alpha_t\alpha_\lambda}{8\pi^2}\Lambda^2, \\
\delta M_{H_u}^2 &= \left[\frac{\alpha_\lambda}{16\pi^2}\left[4\alpha_\lambda - 3\alpha_2 - \frac{3\alpha_1}{5}\right] - \frac{\alpha_\lambda}{12\pi}x^2h(x)\right]\Lambda^2, && (3.42) \\
A_t &= -\frac{\alpha_\lambda}{4\pi}\Lambda.
\end{aligned}$$

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3.4.2 Model 12

The messenger sector can very well be made up of singlets and $10 \oplus \overline{10}$ multiplets. If both are allowed to couple we can have the following new mixed invariant $\mathbf{10\overline{10}_m 1_m}$. In order for this to be possible we can assume that the messenger fields have the following R-parity,

$$Rp(10_m, \overline{10}_m, 1_m, \tilde{1}_m) = (-10_m, -\overline{10}_m, 1_m, \tilde{1}_m), \tag{3.43}$$

and the following flavor charges,

$$\mathrm{U}(1)_F(10_m, \overline{10}_m, 1_m, \tilde{1}_m, X) = (a-b, b, -b, a+b, -a). \tag{3.44}$$

The leading terms in the superpotential are given by,

$$W_{\text{int}} = \lambda_q Q_3 \tilde{Q}_m S_m + \lambda_u U_3^c \tilde{U}^c S_m + \lambda_e E_3^c \tilde{E}_m^c S_m, \tag{3.45}$$

the new contributions to the soft scalar masses are given by,

$$\begin{aligned}
\delta M_Q^2 &= \left[-\frac{\alpha_q x^2 h(x)}{12\pi} + \frac{\alpha_q}{16\pi^2}\left[8\alpha_q + 3\alpha_u + \alpha_e - \frac{16}{3}\alpha_3 - 3\alpha_2 - \frac{1}{15}\alpha_1\right] - \frac{\alpha_t \alpha_u}{16\pi^2}\right]\Lambda^2, \\
\delta M_{U^c}^2 &= \left[-\frac{\alpha_u x^2 h(x)}{12\pi} + \frac{\alpha_u}{16\pi^2}\left[6\alpha_q + 5\alpha_u + \alpha_e - \frac{16}{3}\alpha_3 - \frac{16}{15}\alpha_1\right] - \frac{\alpha_t \alpha_q}{8\pi^2}\right]\Lambda^2, \\
\delta M_{D^c}^2 &= -\frac{\alpha_b \alpha_q}{8\pi^2}\Lambda^2, \\
\delta M_L^2 &= -\frac{\alpha_e \alpha_\tau}{16\pi^2}\Lambda^2, \\
\delta M_{E^c}^2 &= \left[-\frac{\alpha_e x^2 h(x)}{12\pi} + \frac{3\alpha_e}{16\pi^2}\left[\alpha_e + 2\alpha_q + \alpha_u - \frac{4}{5}\alpha_1\right]\right]\Lambda^2, \\
\delta M_{H_u}^2 &= -\frac{3\alpha_t(\alpha_q + \alpha_u)}{16\pi^2}\Lambda^2, \\
\delta M_{H_d}^2 &= -\frac{3\alpha_b \alpha_q + \alpha_e \alpha_\tau}{16\pi^2}\Lambda^2, \\
A_t &= -\frac{\alpha_q + \alpha_u}{4\pi}\Lambda, \qquad A_b = -\frac{\alpha_q}{4\pi}\Lambda, \qquad A_\tau = -\frac{\alpha_e}{4\pi}\Lambda.
\end{aligned} \tag{3.46}$$

3.4.3 Models 13 & 14

Consider scenarios where the messenger sector is made up of vector pairs of $5 \oplus \bar{5}$ and $10 \oplus \overline{10}$. There are two distinct new interaction terms that can arise other than combinations of the models already studied earlier. We are going to consider these two model one by one,

$\mathbf{10 10_m 5_m}$ **:** one can motivate this by considering the following symmetry arrangements. Consider the following R-parity,

$$Rp(5_m, \bar{5}_m, 10_m, \overline{10}_m) = (-5_m, -\bar{5}_m, 10, \overline{10}_m), \tag{3.47}$$

and the following flavor charges,

$$\mathrm{U}(1)_F(5_m, \bar{5}_m, 10_m, \overline{10}_m, X) = (b, a-b, -b, a+b, -a), \tag{3.48}$$

with $a \gg 1$ as always. Thus the part of the superpotential that remains unsuppressed by the flavor factors is given by,

$$W_{\rm int} = \lambda_{q1} Q_3 U^c_m H^m_u + \lambda_{u1} Q_m U^c_3 H^m_u + \lambda_{q2} Q_3 Q_m \tilde{D}^c_m + \lambda_{u2} U^c_3 E^c_m \tilde{D}^c_m + \lambda_e U^c_m E^c \tilde{D}^c_m, \quad (3.49)$$

Note that proton decay can occur at one loop through the interactions in the superpotential. The suppression by the Froggatt-Nielsen factor notwithstanding, it requires a severe fine-tuning of the superpotential parameters to be consistent with proton decay constraints. Assuming additional discrete symmetry can ameliorate this problem, the new contributions to the soft masses are given by,

$$\begin{aligned}
\delta M^2_Q &= \Bigg[-\frac{\alpha_{q1}+\alpha_{q2}}{12\pi} x^2 h(x) + \frac{\alpha_{q1}}{16\pi^2}\left(6\alpha_{q1} + \alpha_e + 3\alpha_{u1} - \frac{16}{3}\alpha_3 - 3\alpha_2 - \frac{13}{15}\alpha_3\right) \\
&\quad + \frac{\alpha_{q2}}{16\pi^2}\left(4\alpha_{q2} + 2\alpha_{q1} + \alpha_{u1} + \alpha_{u2} + \alpha_e - 8\alpha_3 - 3\alpha_2 - \frac{1}{5}\alpha_1\right) \\
&\quad - \frac{\alpha_t}{16\pi^2}(2\alpha_{u1} + \alpha_{u2})\Bigg]\Lambda^2, \\
\delta M^2_{U^c} &= \Bigg[-\frac{2\alpha_{u1}+\alpha_{u2}}{12\pi} x^2 h(x) + \frac{\alpha_{u1}}{8\pi^2}\left(6\alpha_{u1} + 3\alpha_{q1} + \alpha_{q2} - \frac{16}{3}\alpha_3 - 3\alpha_2 - \frac{13}{15}\alpha_1\right) \\
&\quad + \frac{\alpha_{u2}}{16\pi^2}\left(5\alpha_{u2} + 4\alpha_{u1} + 2\alpha_{q2} + \alpha_e - \frac{16}{3}\alpha_3 - \frac{28}{15}\alpha_1\right) - \frac{\alpha_t}{8\pi^2}(\alpha_{q1} + \alpha_{q2})\Bigg]\Lambda^2, \\
\delta M^2_{D^c} &= -\frac{\alpha_b(\alpha_{q1}+\alpha_{q2})}{8\pi^2}\Lambda^2, \\
\delta M^2_L &= -\frac{3\alpha_e\alpha_\tau}{16\pi^2}\Lambda^2, \\
\delta M^2_{E^c} &= \left[-\frac{\alpha_e}{4\pi}x^2 h(x) + \frac{3\alpha_e}{16\pi^2}\left(5\alpha_e + 2(\alpha_{q1}+\alpha_{q2}) + \alpha_{u2} - \frac{16}{3}\alpha_3 - \frac{28}{15}\alpha_1\right)\right]\Lambda^2, \qquad (3.50) \\
\delta M^2_{H_u} &= -3\alpha_t \frac{\alpha_{q1}+\alpha_{q2}+2\alpha_{u1}+\alpha_{u2}}{16\pi^2}\Lambda^2, \\
\delta M^2_{H_d} &= -3\frac{\alpha_b(\alpha_{q1}+\alpha_{q2})+\alpha_e\alpha_\tau}{16\pi^2}\Lambda^2, \\
A_t &= -\frac{\alpha_{q1}+\alpha_{q2}+2\alpha_{u1}+\alpha_{u2}}{4\pi}\Lambda, \qquad A_b = -\frac{\alpha_{q1}+\alpha_{q2}}{4\pi}\Lambda, \qquad A_\tau = -\frac{3\alpha_e}{4\pi}\Lambda.
\end{aligned}$$

$\mathbf{\overline{10}_m 5_m 5_H}$: In this case consider the following R-parity,

$$Rp(5_m, \bar{5}_m, 10, \overline{10}_m) = (5_m, \bar{5}_m, 10_m, \overline{10}_m), \quad (3.51)$$

and the following flavor charges,

$$\mathrm{U}(1)_F(5_m, \bar{5}_m, 10_m, \overline{10}_m, X) = (b, a-b, a+b, -b, -a). \quad (3.52)$$

The corresponding superpotential is given by,

$$W_{\rm int} = \lambda_q \tilde{Q}_m \tilde{D}^c_m H_u + \lambda_e \tilde{E}^c_m H^m_u H_u, \quad (3.53)$$

which results in the following new contributions to the soft masses,

$$\begin{aligned}
\delta M_Q^2 &= -\frac{\alpha_t(\alpha_e + 3\alpha_q)}{16\pi^2}\Lambda^2, \\
\delta M_{U^c}^2 &= -\frac{\alpha_t(\alpha_e + 3\alpha_q)}{8\pi^2}\Lambda^2, \\
\delta M_{H_u}^2 &= \left[-\frac{3\alpha_q + \alpha_e}{12\pi}x^2 h(x) + \frac{\alpha_e}{16\pi^2}\left(4\alpha_e - 3\alpha_2 - \frac{9}{5}\alpha_1\right)\right. \\
&\qquad \left. + \frac{\alpha_q}{8\pi^2}\left(9\alpha_q + 3\alpha_e - 8\alpha_3 - \frac{9}{2}\alpha_2 - \frac{7}{10}\alpha_1\right)\right]\Lambda^2, \\
A_t &= -\frac{\alpha_e + 3\alpha_q}{4\pi}\Lambda.
\end{aligned} \tag{3.54}$$

4 Comparison and results

In this section we will compare the models introduced earlier. The ability of these models to raise the Higgs mass, without requiring a large stop mass, is through a sizable top trilinear coupling A_t. However it is clear from the expressions for the soft spectrum given in eqs. (3.4) and (3.6) that the terms that lead to a large A_t are also responsible for tachyonic contribution to the soft masses. The size of the trilinear coupling is thus constrained by the condition that all sfermion masses should remain positive and that proper radiative electroweak symmetry breaking (REWSB) should take place. In table 2 we list the models along with the sfermions that are most vulnerable to the tachyonic contributions. Assuming that the one loop tachyonic contributions are negligible owing to the extra suppression from the messenger breaking scale, it is possible to put upper limit on the size of the trilinear couplings. We exhibit these upper limits for models wherever they are analytically possible by relating them to the gluino mass, which can be considered as the representative of the soft masses. Some observations become apparent from the table,

1. The right handed sleptons get the smallest contribution from gauge interactions. Thus in models where they receive negative contributions from the messenger-matter interactions they become susceptible to turn tachyonic. This constrains the size of the trilinear couplings in these models and thus effects their efficiency in raising the Higgs masses.

2. Radiative electroweak symmetry breaking either requires a large stop mass that can turn one of the neutral Higgs eigenvalue negative radiatively or a contribution from the messenger-matter interactions that drives it to a negative value. An analytic study of each of these models in terms of electroweak symmetry breaking is difficult and we will rely on a numerical simulation of these models for this purpose.

The discussion above is confined to the spectrum at the supersymmetry breaking scale which already gives us an insight to models in terms of their ability to generate large trilinear couplings and thus ease the fine-tuning in the models. We point out that a generation of a large trilinear is not enough to alleviate the problem with the Higgs mass. In principle the subsequent renormalization group evolution to low energies can wash out

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Model	Main constraint on A_t	Analytic limit on $\frac{\lvert A_t\rvert}{M_{\tilde g}}\rvert_{\max}$	$m_{\tilde t_1}\rvert_{\min}$ TeV	NLSP	Ref.
1	$M^2_{\tilde E}$	$\frac{3}{5}\frac{\alpha_1^2}{\alpha_\tau\alpha_3}$	2.2	$\chi_1^0,\ \tilde\tau_1$	[11]
2	$M^2_{\tilde E}$	$\frac{3}{5}\frac{\alpha_1^2}{\alpha_\tau\alpha_3}$	1.9	$\chi_1^0,\ \tilde\tau_1$	New
3	$M^2_{\tilde L}$	$\frac{1}{\alpha_\tau\alpha_3}\left[\frac{\alpha_2^2}{2}+\frac{\alpha_1^2}{10}\right]$	0.6	$\chi_1^0,\ \tilde\tau_1,\ \tilde\mu_1$	[12, 13, 17]*
4	$M^2_{\tilde D^c},\ M^2_{\tilde L}$	$\frac{\alpha_1^2}{5\alpha_3}\left(\frac{1}{\alpha_b}+\frac{1}{2\alpha_\tau}\right)$ $+4\frac{\alpha_3}{\alpha_b}+\frac{\alpha_2^2}{2\alpha_\tau\alpha_3}$	>3	$\chi_1^0,\ \tilde\tau_1,\ \tilde\nu_e$	[13, 16]*
5	unconstrained	—	0.1	$\chi_1^0,\ \tilde\mu_1,\ \tilde\nu_e$	[13]*
6	$M^2_{\tilde U_3^c}$	$\frac{4}{3}\left[\frac{\alpha_3}{\alpha_t}+\frac{\alpha_1^2}{5\alpha_3\alpha_t}\right]$	0.7	$\tilde\tau_1,\ \tilde\mu_1$	New
7	unconstrained	—	0.5	$\tilde\chi_1^0,\ \tilde\mu_1$	New
8	unconstrained	—	0.5	$\tilde\chi_1^0,\ \tilde\mu_1$	New
9	unconstrained	—	0.5	$\tilde\tau_1,\ \tilde\chi_1^0$	[12, 17]*
10	$M^2_{\tilde U^c}$	$\frac{4}{3}\left[\frac{\alpha_3}{\alpha_t}+\frac{\alpha_1^2}{5\alpha_3\alpha_t}\right]$	0.6	$\tilde\tau_1$	[14]
11	$M^2_{\tilde U^c}$	$\frac{4}{3}\left[\frac{\alpha_3}{\alpha_t}+\frac{\alpha_1^2}{5\alpha_3\alpha_t}\right]$	1.9	$\tilde\tau_1$	[14]
12	REWSB	—	2.9	$\tilde\chi_1^0$	New
13	unconstrained	—	0.8	$\tilde\tau_1,\ \chi_1^0,\ \tilde\mu_1,\ \tilde\nu_1$	New
14	$M^2_{\tilde U^c}$	$\frac{4}{3}\left[\frac{\alpha_3}{\alpha_t}+\frac{\alpha_1^2}{5\alpha_3\alpha_t}\right]$	1.2	$\tilde\tau_1$	New

Table 2. We summarizer the main theoretical limit on the generated top trilinear coupling where they are analytically possible. We also list the smallest value of the lightest stop ($m_{\tilde t_1}\rvert_{\min}$) for which the Higgs mass could be boosted to the range $123-127$ GeV as obtained from our numerical scanning. Finally the NLSP for each model is listed. We also list the references for models which have been studied previously in the literature. * indicates that these models are slight variants of the references given or we have updated the older analysis to include all two loop corrections. Also note that the eventually the 'unconstrained' are limited by renormalization group running effects.

the trilinear coupling generated at the supersymmetry breaking scale due to the partial cancellation between the Yukawa and gaugino contributions to the β function [34]. In fact a large A_t at high scale is useful if it has a sign opposite to that of the gaugino masses. In the class of models studied here, this is naturally achieved as is evident from eq. (3.2) and eq. (3.4). The sign difference between the trilinear couplings and the gaugino masses can be traced back to the overall negative factor for fermionic loops in the leading contribution to the trilinear coupling. This generic feature of this class of models is aided by our choice of a low scale of supersymmetry breaking which reduces any potential harmful effect of the renormalization group running by shortening the range.

Next we perform an extensive numerical study of each of the above models to make a numerical comparison between them. The numerical procedure that is followed for each of these models is described below:

1. We consider $\Lambda = F/M$ and $x = \Lambda/M$ as the two independent parameters that define all the scales in the theory. In order to scan over the parameter space of the models we vary Λ between $4 \times 10^4 - 9 \times 10^5$ GeV. We consider two different values of $x \sim .01, .1$ to include scenarios where the one loop contributions become insignificant and comparable to the corresponding two loop contributions respectively.

2. We use the known Standard Model fermion masses at the weak scale and $\tan\beta$ as the boundary conditions and use one loop renormalization group equations to determine the known coupling constant at the messenger scale defined by $M_{\rm mess} = \Lambda/x$.

3. For given values of Λ, x, MSSM couplings and a given set of the new messenger-matter Yukawas we generate the soft spectrum at the messenger scale.

4. We then use one loop renormalization group equations [34] to determine the weak scale spectrum. Within the framework of pure gauge mediation the μ term in the MSSM superpotential cannot be generated. Note that messenger-matter interactions can in principle lead to generation of a μ term at the messenger scale, see [14, 35]. However all the models studied in this paper except Model 1 do not give rise to a μ term.[4] In our numerical study we will simply impose the matching condition at the weak scale for the potential minima and set the weak scale μ parameter from the following relation,

$$m_Z^2 \simeq -2(m_{H_u}^2 + |\mu|^2) + \frac{2}{\tan^2\beta}(m_{H_d}^2 - m_{H_u}^2). \tag{4.1}$$

5. Finally we use SuSpect [36] and micrOMEGAS [37, 38] to determine correct electroweak symmetry breaking, diagonalize the sparticle mass matrices and determine weak scale observables.

6. We scan over all the new messenger-matter couplings λ_i between $.05 - 1.5$. For every choice of the Yukawa couplings we repeat the above procedure to scan over the parameter space of the models.

In figure 1 we correlate the maximum attainable Higgs mass to the lightest stop mass for the various Models discussed in section 3. We take care to sum over all model parameters

[4] However generation of the μ term through gauge mediation can be incorporated in these models through the usual extension to a Z_3 NMSSM like scenario [14].

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and extract the outer envelop of the valid model points to obtain the displayed curves. The horizontal band corresponds to the presently preferred range of Higgs mass measured at the LHC. This allows a direct comparison between the models introduced in the previous section. To reduce clutter we have not displayed the plot for Model 8 which is a subset of the Model 7 in the figures. As anticipated from table 2, models where the size of the trilinear coupling is constrained at the messenger scale show marginal improvement over the pure GMSB models. Considerations of radiative electroweak symmetry breaking also enter the game. The table indicates the smallest stop mass for which we obtain a Higgs in the desired mass range for every model. A discussion about the correlation between the constraints on the trilinear couplings and the lightest stop mass is complicated by the effects of the renormalization group running. However a couple of comments regarding some of the apparent contradictions in table 2 are in order: (i) Note that though the limit on the trilinear coupling is more relaxed in Model 4 as compared to Model 3, the former receives large positive two loop contributions to $M^2_{\tilde{Q}}$ and $M^2_{\tilde{U}^c}$ thus making the stop spectrum heavier. This should be contrasted with the negative contribution to $M^2_{\tilde{U}^c}$ in Model 3. However the numerically dominant effect comes from the purely negative contribution to $M^2_{H_d}$ in Model 4 that inhibits successful REWSB in large region of the parameter space. (ii) While the trilinear couplings in Models 10 and 11 have the same limits they yield very different low energy spectrum. This is again related to the difference in the contributions to $M^2_{H_u}$ resulting a more stringent constraint from REWSB for Model 11 as compared to Model 10. Note that Model 14 effectively combines Models 10 and 11, however a larger Dynkin index makes the spectrum relatively heavier.

Clearly from this Models 5, 7 ,9, and 10 can be identified as most promising in terms of boosting the Higgs mass to experimentally favored range without admitting too much fine-tuning. In figure 2 we show the allowed region for the Models 5, 7 ,9 and 10 where the Higgs mass is between $123-127\,\mathrm{GeV}$, in the parameter space of the lightest stop and the gluino masses. These plots can be directly translated to a measure of naturalness of the models by relating them to the *Barbieri-Giudice* [39] fine-tuning parameter using the approximate relation [40], $\Delta \simeq \mathcal{O}(1)10t/33(\lambda_t M_S/650\,GeV)^2$, where $t = \log[M_{\rm mess}/M_Z]$, λ_t is the top Yukawa coupling and $M_S = \sqrt{m_{\tilde{t}_1} m_{\tilde{t}_2}}$. A recent study in pure mGMSB models [41] indicates that for a single messenger pair in $5\oplus\overline{5}$ of SU(5) requires a stop mass beyond 4 TeV to obtain a 123 GeV Higgs that tantamount to a considerable fine-tuning. We note that the benchmark points given in table 3 implies an improvement in the fine-tuning at $\Delta/\Delta_{\rm mGMSB} \simeq 0.03-0.11$. Thus an order of magnitude improvement in fine-tuning can be achieved through the messenger-matter interactions discussed in this paper.

Within the framework of GMSB models, the collider constraints on the sparticle masses are rather model sensitive, crucially depending on the NLSP [42]. Interestingly we observe that among the preferred models only one has a nutralino NLSP while the rest have slepton NLSPs.[5] The search strategies and hence the consequent constrains are different for the

[5]As a small digression it is curious to note that a light stau can lead to enhanced contribution to the $Br(h\to\gamma\gamma)$ rate through loop contributions in the large $\tan\beta$ regime [43] and can explain the slight excess over SM prediction indicated by the present experimental data.

Model	5	7	9	10
Messenger -matter Yukawa	$\lambda_q = 1.05$, $\lambda_u = 1.25$, $\lambda_l = 1.05, \lambda_e = .45$	$\lambda_{11} = \lambda_{22} = 0$, $\lambda_{12} = \lambda_{21} = 1.3$, $\lambda_u = 1.05, \lambda_e = .05$	$\lambda_q = .25$, $\lambda_u = .75$, $\lambda_h = 1.05$	$\lambda_q = 1.15$
Λ	1.4×10^5	8×10^4	4×10^4	7×10^4
M_h	127	126	125	125
M_H	2792	1147	257	829
M_A	2792	1147	257	829
M_{H^+}	2793	1150	271	733
NLSP	χ_1^0	$\tilde{\mu}_1$	$\tilde{\tau}_1$	$\tilde{\tau}_1$
$m_{3/2}$	29.8 eV	9.7 eV	2.4 eV	7.5 eV
$M_{\chi_1^0}, M_{\chi_2^0}$	190, 376	216, 429	123,171	283,550
$M_{\chi_3^0}, M_{\chi_4^0}$	2923, 2924	1954, 1956	185, 349	753, 770
$M_{\chi_1^\pm}, M_{\chi_2^\pm}$	376 , 2925	429 , 1956	149, 349	550, 770
$M_{\tilde{g}}$	1298	1491	1118	1947
$M_{\tilde{t}_1}, M_{\tilde{t}_2}$	766, 2374	879, 1741	498, 850	985, 1557
$M_{\tilde{b}_1}, M_{\tilde{b}_1}$	2363, 3142	1503,1726	732, 1041	985, 1557
$M_{\tilde{u}_R}, M_{\tilde{u}_L}$	1635, 1719	1557, 1600	1061,1087	1849,1895
$M_{\tilde{d}_R}, M_{\tilde{d}_L}$	1646, 1721	1541, 1602	1054, 1090	1832,1897
$M_{\tilde{\tau}_1}, M_{\tilde{\tau}_2}$	1342, 3347	128, 479	78, 294	56, 522
Fine-tuning $\left(\frac{\Delta}{\Delta_{\rm mGMSB}}\right)$	0.11	0.10	0.03	0.10

Table 3. Representative spectrum for the chosen models. All the masses except the gravitino mass are given in the GeV unit. Here $\tan\beta = 10$ and $x = 0.1$ The spectrum is chosen conservatively to be consistent with latest experimental bounds from direct collider studies and low energy observables.

two scenarios. In case of a bino NLSP one expects a $2\gamma + E_T^{miss}$ signal, the limits were discussed in a fairly model independent manner through *simplified models* in [44]. The limits presented for the natural SUSY (light stop) scenario indicate $m_{\tilde{g}} > 1.1\,\text{TeV}$ and $m_{\tilde{t}_1} > 700\,\text{GeV}$. Subsequent updates [45, 46] marginally enhance the exclusion on the gluino. The models with slepton NLSP are constrained from SS/OS *dileptons* $+ E_T^{miss}$ and $\tau +$ *leptons* $+$ *jets* $+ E_T^{miss}$ searches. Conservative limits around $m_{\tilde{g}} > 800\,\text{GeV}$ were suggested in [42]. Recent updates from the CMS, see for example [47–49] and the ATLAS, see [50] will certainly increase this bound. However we could not find a comprehensive model independent study of this scenario in terms of simplified models in the literature.

In table 3 we present indicative benchmark points for the four preferred models. Clearly the limits on the neutralino NLSP scenario already push the fine-tuning of the models beyond 10. Let us also point out that recent study on the gravitino phenomenology indicates that a cosmologically safe upper limit gravitino mass might be as low as 16 eV [52]. The gravitino mass is given by the usual relation $m_{3/2} = F/(\lambda_\phi \sqrt{3} M_{Pl})$ where F and λ_ϕ are defined in eq. (3.1) and M_{Pl} is the Planck mass. Comparing with the gaugino masses given in eq. (3.2) we see that,

$$\frac{m_{3/2}}{M_{\tilde{g}}} \propto \frac{M}{M_{Pl}} \frac{1}{N_{\rm mess} d}, \tag{4.2}$$

where M,$N_{\rm mess}$ and d are the messenger scale, messenger multiplicity and the Dynkin index respectively. It is clear that the restrictive limits from the gravitino mass can be accommodated in models with higher messenger multiplicity and/or high Dynkin index. It might be expected that the large gravitino mass (e.g. the benchmark point for Model 5 in table 3) for models with low Dynkin index can be handled by considering higher messenger multiplicity. A detailed study of the effect of varying the number of messengers is beyond the scope of this paper.

Note that most of the discussion in this paper corresponds to this light gravitino scenario. However as the gravitino mass crosses $\sim 50\, MeV$ the NLSP becomes increasingly stable leading to stronger constraints from searches for long living particles at the LHC. For example a quasi-stable stau has the constraints $m_{\tilde{\tau}_1} > 223\,\mathrm{GeV}$ on its mass [51]. Consequently this would rule out most of the low fine-tuned regions in our models providing an added motivation to consider low scale of supersymmetry breaking.

5 Conclusion

Supersymmetric scenarios with viable UV complete supersymmetry breaking sectors are severely challenged by the measured Higgs mass at the LHC. Several extension of these models that conjecture the existence of extra exotic particles have been proposed in the literature to address this issue [53–58] . Within gauge mediated models one can evade these constraints economically, by considering messenger-matter interactions. Some of these models have also been studied in the recent past, [11–17]. In the present paper we study a class of these models that can be embedded in SU(5) GUT group. We find R-parity and flavor symmetries provide an organized way to study these models. Interestingly the flavor symmetry can be tied to the origin of the observed SM fermion masses and mixing through a Froggatt-Nielsen like framework. We construct all the possible invariant terms with the messengers in the 1, 5 and 10 dimensional representations of the SU(5) group. Some models or their close variants studied in the literature show up naturally within this framework while a whole set of new models are predicted. However it should be emphasized that the choice of flavor symmetries presented here is not unique. They are presented mostly in the spirit of proof of principle. It might as well be that the form of the interaction superpotential discussed here arises from a completely different underlying flavor structure. Nevertheless the phenomenological features studied in the second half of the paper are independent of these assumptions.

A detailed numerical study of the relevant parameter space of the different scenarios is carried out to compare these models. We identify models that can effectively raise the Higgs mass to the favoured range without admitting too much fine-tuning. We find

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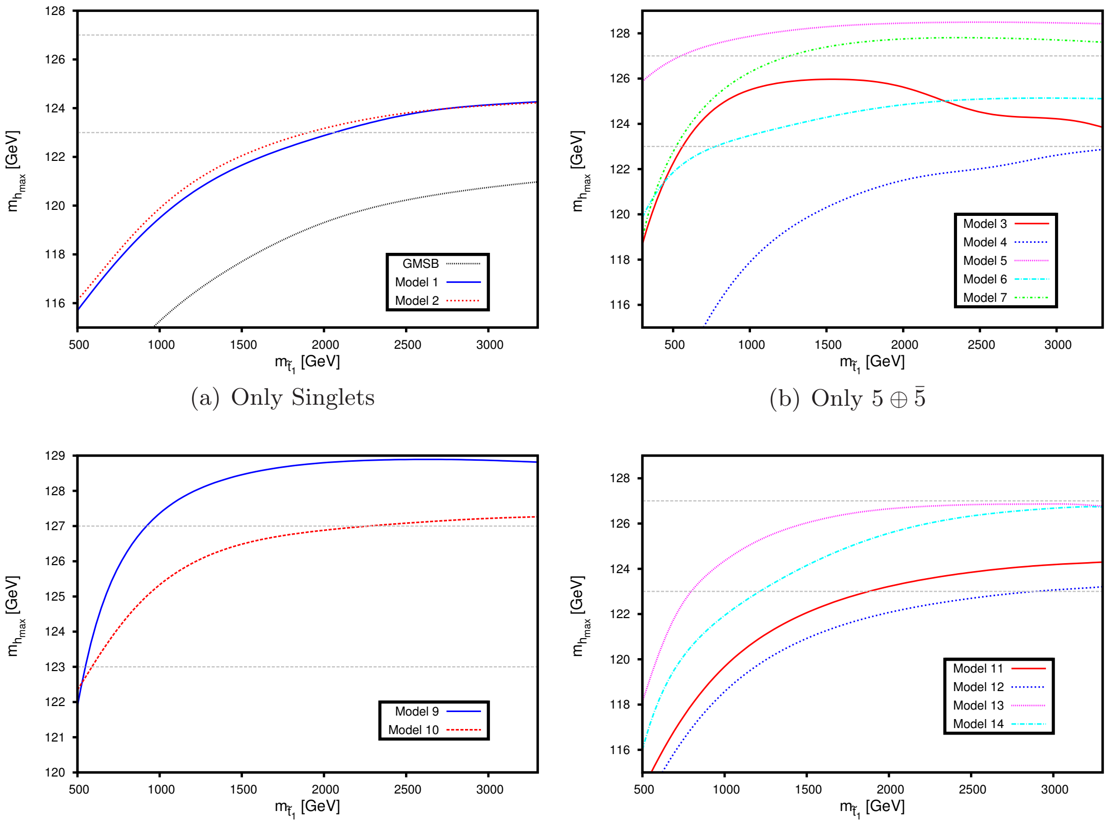

(a) Only Singlets (b) Only $5 \oplus \bar{5}$

(c) Only $10 \oplus \overline{10}$ (d) Mixed messenger models

Figure 1. The maximum Higgs mass for the different models is plotted against the lightest stop mass. In all the models the new Yukawa couplings are varied between $0.1 - 1.5$, $\tan\beta$ between $5 - 50$, Λ between $2 \times 10^4 - 9 \times 10^5$ and two different choices of $x = .01$ & $.1$ were considered. The region below each curve is accessible to the corresponding model. 'GMSB' in (a) represents the mGMSB model with a single messenger pair in $5 \oplus \bar{5}$ of SU(5).

many of the generic and specific SUSY searches at LHC imply considerable constraints on these models. Incidentally we observe that all but one of the preferred models have a slepton NLSP. However we could not locate a detailed model independent study of the LHC constraints on this scenario in the literature. Most studies are restricted to benchmark points or related to details of the entire spectrum. Considering that many other possible *natural* SUSY models within the context of GMSB could lead to a slepton NLSP, a detailed study of this scenario is highly anticipated.

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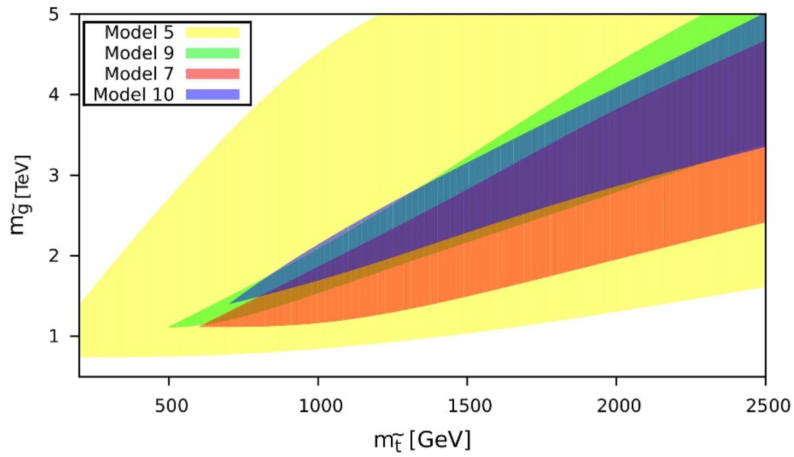

Figure 2. Allowed region plot of the preferred models which correspond to m_h between 123–127 GeV in the plane of the lightest stop mass and the gluino mass.

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