Gravitational rescue of minimal gauge mediation

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Abstract

Gravity mediation supersymmetry breaking become comparable to gauge mediated supersymmetry breaking contributions when messenger masses are close to the GUT scale. By suitably tuning the gravity contributions one can then modify the soft supersymmetry breaking sector to generate a large stop mixing parameter and a light Higgs mass of 125 GeV. In this kind of hybrid models, however the nice features of gauge mediation like flavour conservation etc, are lost. To preserve the nice features, gravitational contributions should become important for lighter messenger masses and should be important only for certain fields. This is possible when the hidden sector contains multiple (at least two) spurions with hierarchical vacuum expectation values. In this case, the gravitational contributions can be organised to be ‘just right’.

We present a complete model with two spurion hidden sector where the gravitational contribution is from a warped flavour model in a Randall-Sundrum setting. Along the way, we present simple expressions to handle renormalisation group equations when supersymmetry is broken by two different sectors at two different scales.

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I. INTRODUCTION AND MOTIVATION

The discovery of the Higgs boson with a mass $\sim 126$ GeV a couple of years ago \cite{1-4} has led to tremendous excitement in the field. While the discovery has validated the Higgs mechanism of the Standard Model, and is broadly consistent with expectations based on Minimal Supersymmetric Standard Model (MSSM) \footnote{We mean that the observed higgs is within theoretical upper bound of 135 GeV \cite{5-7}.}, minimal models of supersymmetry breaking like gauge mediation have been strongly constrained by this discovery.

Gauge Mediation has several nice features that make it an attractive mechanism for supersymmetry breaking: (i) no additional flavour violation in the soft sector (ii) minimal models have very few parameters, some times as low as one, making them very predictive (iii) different phenomenology compared to the traditional gravitational models. However, the discovery of the Higgs boson with mass $\sim 126$ GeV, has strongly constrained these models. The one loop corrected mass for the lightest CP even Higgs boson has the form \cite{5-10}

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

(1)

Where $M_S = \sqrt{m_{\tilde{t}_1}m_{\tilde{t}_2}}$ and $X_t = A_t - \mu \cot \beta$. For $M_S$ around 1 TeV, it can be seen that $X_t$ should be close to it’s maximal value ($\approx \sqrt{6}M_S$) to obtain a Higgs mass around 126 GeV \cite{11-14}. As is well known, in minimal gauge mediation, the trilinear A terms are highly suppressed at the mediation scale, which leads to much small values of $X_t$ than required at the weak scale. One can of course try to increase the mediation scale to increase the renormalisation group (RG) running from traditional values of $\sim 100$ TeV all the way up to the GUT scale. However for $\sim 1$ TeV stops it has been shown that a sufficient appreciation in the $X_t$ value is only possible with heavy gluinos and/or tachyonic stops at High scale \cite{15}. Since $X_t$ is very small in GMSB, $M_S \sim 4$ TeV is required to achieve the desired Higgs mass. This scenario is not very attractive as it comes at the expense of making the stop mass eigenvalues out of reach of the LHC. Thus minimal versions of GMSB models would fail to accommodate 125 GeV Higgs, if the SUSY particle masses lie below 2 TeV or so.

Various extensions have been proposed in the literature to make models of gauge mediation compatible with the 125 GeV Higgs mass \cite{16-31}. For instance, one could consider explicit Yukawa like messenger matter mixing in addition to gauge interactions\cite{32-34}. In some cases of this type it is possible to get solutions with minimal fine tuning, leading to a naturally light Higgs boson.
with mass 125 GeV\cite{25,35}. Another possibility is to consider, anomaly free, U(1) gauge extension of MSSM gauge group\cite{36–44}.

Alternatively, one can ask the question whether gravity mediated supersymmetry breaking contribution can help in generating large $A_t$, while keeping the spectrum within the reach of LHC ($\lesssim 1.5$ TeV). Note that these contributions are always present as Planck scale suppressed operators. Typically, this contribution is much suppressed due to the relatively "low" values of $F$ -terms required by gauge mediation. For example, the stop trilinear coupling at the reduced Planck would would be generated as

$$A_t \simeq \frac{\langle F_X \rangle}{M_{Pl}}$$

where $\langle F_X \rangle$, is the vacuum expectation value (VEV) of the F-term of the spurion multiplet and $M_{Pl} \sim 10^{18}$ GeV, is the reduced Planck scale. In gauge mediation models, typically $\langle F_X \rangle$ lies in the range $\sim 10^{10} – 10^{12}$ GeV$^2$. This $F_X$ value is too small to have any meaningful impact on the soft sector when contributions from gravity mediation are included; they are of the order $\sim (10^{-8} – 10^{-6})$ GeV. As will be discussed in the next section, if one increases the mass of the messengers the gravitational contribution become increasingly important. This result holds true as long as supersymmetry breaking by as single spurion parametrizes the hidden sector.

Consider the case where the hidden sector is parameterised by more than one spurion fields, the gravity contribution could be much larger as one can effectively ‘decouple’ both the gravity and gauge contributions. In such case, even if the messenger masses are smaller, gravitational mediation can play an important role. Models with multiple spurion fields are for example quite common in string based models \cite{45–54}, mirage mediation \cite{55–61} and models with multiple hidden sectors \cite{46,47,61–63}. In the recent times, supersymmetry breaking in multiple sectors has been receiving attention in literature. Most of the works are concentrated on Goldstini sector and their corresponding collider phenomenology \cite{64,67}.

The combination of gauge and gravity mediations, some times called “hybrid models” of supersymmetry breaking \cite{68,69} has been studied in literature with focus on specific issues in phenomenology like flavour or dark matter. In the present work, we show that this idea can be used to address the light CP even Higgs mass problem in minimal gauge mediation. However, it should be noted that the framework presented here is quite different compared to the ones in the literature.

In the present work, we focus on the higgs sector, especially concentrate on how to increase its mass through additional gravity mediation contributions to the minimal messenger model. It turns out that the spurions should have hierarchical F -terms, and furthermore the gravity contributions
should be ‘flavoured’ in the sense that the $A$ terms of the third generation should be larger compared to the first two generations. Similarly the soft terms should be such that the total spectrum should still be dominated minimal gauge mediation.

The paper is organized as follows. In the next section we discuss the character of the gravitational contributions required to fix the light Higgs mass in minimal messenger model. In section 3, we present an explicit model for gravity sector which has the essential features and also present explicit numerical examples. We close with a summary.

II. CHARACTER OF GRAVITATIONAL CONTRIBUTIONS

In minimal gauge mediation, soft terms are generated through 1-loop and 2-loop diagrams involving messenger fields coupled to the MSSM fields through gauge interactions. The messengers are connected through the hidden sector which is parameterised by a spurion field $X$ as

$$W_{\text{mess}} = \lambda X \Phi \Phi$$

where $\Phi, \Phi$ are messenger fields and $X = M_X + \theta^2 F$ parameterises the hidden sector. $M_X$ is the mass of the messenger fields and $F$ is the supersymmetry breaking F-term. The soft terms due to gauge mediation at the messenger scale $M_X$ are given by 70–73:

$$M_a' \approx \frac{\alpha_a}{4\pi} \left( \frac{F}{M_X} \right)$$
$$m'^2 \approx 2 \sum_{a=1}^3 \left( \frac{\alpha_a}{4\pi} \right)^2 C_a \left( \frac{F^2}{M_X^2} \right)$$
$$A' \approx 0$$

Typically $\Lambda \equiv F/M_X$ is chosen to be around 100 TeV to get low energy soft terms of $O(1)$ TeV. Generally $M_X$ is chosen to be close to $\Lambda$ (twice of $\Lambda$ is the usual choice) and it can be as large as the GUT scale. $F$ is scaled accordingly while keeping $\Lambda$ fixed. The low energy spectrum is computed by running of the renormalisation group equations (RGE) from $M_X$ to the weak scale. Since we are concerned mostly with the light CP even Higgs mass, we will concentrate on the parameters relevant for it. At the weak scale, the relevant parameters appearing in eq. 3 can be written as

$$m_{Q_3}^2(M_{\text{weak}}) \approx 1 \times 10^{-4} \Lambda^2 \text{ GeV}^2$$
$$m_{U_3}^2(M_{\text{weak}}) \approx 7.62 \times 10^{-5} \Lambda^2 \text{ GeV}^2$$
$$A_t(M_{\text{weak}}) \approx -0.002 \Lambda \text{ GeV}$$

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where \( \Lambda = \langle F \rangle / M_X \). The expressions are evaluated for \( M_X = 10^6 \). Keeping \( \Lambda \) fixed, if \( M_X \) is increased from \( 10^6 \) to \( 10^{14} \), then the magnitude of \( A_t \) only increases from 0.002\( \Lambda \) to 0.005\( \Lambda \). Increasing \( \Lambda \) will also increase \( m_{\text{susy}} \), the SUSY scale, thus spoiling naturalness. We can then use the Eq. (1) to compute the light Higgs mass. It is clear from the above that even by choosing a large \( M_X \) close to the GUT scale, so as to increase the logarithmic running, the Higgs mass barely manages to go beyond the LEP limit of 114.2 GeV. This is because \( X_t \) remains small, as long as \( M_S \) is fixed to be 1 TeV. This result is now well known [15].

Gravity mediation will always contribute to the soft terms at the high scale. However, for the full range of \( F \sim 10^{12} - 10^{21} \) GeV\(^2 \) (fixing \( \Lambda = 100 \) TeV), it’s clear that the gravitational contributions are only significant when the messenger scale reaches the GUT scale. The gravity contributions at the Planck scale are given in terms of effective operators as

\[
\begin{align*}
\tilde{m}_{ij}^2 (M_X) & \approx m_{3/2}^2 \left[ \frac{\alpha_a^2}{(4\pi)^2} C_a \frac{M_{Pl}^2}{M_X^2} \delta_{ij} + \tilde{k}_{ij} + \tilde{\beta}_{ij} \right] \\
\tilde{M}_{ab} (M_X) & \approx m_{3/2} \left[ \frac{\alpha_a}{(4\pi)} \frac{M_{Pl}}{M_X} + \gamma_a \right] \delta_{ab} \\
\tilde{A}_{ijk} (M_X) & \approx m_{3/2} \left[ \tilde{\eta}_{ijk} + \xi_{ijk} \right]
\end{align*}
\]

where \( k_{ij} \) and \( \eta_{ijk} \) are arbitrary \( O(1) \) parameters and \( f_{ab} \) is the gauge kinetic function. At the messenger scale \( M_X \), if the running is not significantly large from \( M_{Pl} \) to \( M_X \), the soft terms including both gauge and gravity contributions can be parameterised as

\[
\begin{align*}
\tilde{m}_{ij}^2 (M_X) & \approx m_{3/2}^2 \left[ \frac{\alpha_a^2}{(4\pi)^2} C_a \frac{M_{Pl}^2}{M_X^2} \delta_{ij} + \tilde{k}_{ij} + \tilde{\beta}_{ij} \right] \\
\tilde{M}_{ab} (M_X) & \approx m_{3/2} \left[ \frac{\alpha_a}{(4\pi)} \frac{M_{Pl}}{M_X} + \gamma_a \right] \delta_{ab} \\
\tilde{A}_{ijk} (M_X) & \approx m_{3/2} \left[ \tilde{\eta}_{ijk} + \xi_{ijk} \right]
\end{align*}
\]

where we have used \( \langle F \rangle = m_{3/2} M_{Pl} \) and \( \tilde{k} \) and \( \tilde{\eta} \) are the renormalisation group corrected \( \kappa \) and \( \eta \) couplings respectively and are also \( O(1) \). \( \beta, \xi \) and \( \gamma \) are corrections to the original couplings due to RG running. These parameters contain the Clebsch factors (which could be large in the presence of GUT group between \( M_{Pl} \) and \( M_X \)) and the logarithmic factors due to running. Numerically \( \beta = \xi \approx \frac{1}{16\pi^2} \log \left( \frac{M_{Pl}}{10^{16}} \right) \sim 0.02 \). At one loop, the co-efficients \( \gamma \) can be determined exactly and are given as \( \gamma_1 = 0.96 \quad \gamma_2 = 0.99 \quad \gamma_3 = 1.01 \). We note here that the gravitino mass in this case \( (m_{3/2} \sim 100 \) GeV\) is large as compared to low scale GMSB, where typically \( \langle F \rangle << M_{Pl} \). This is because, keeping \( \Lambda = 10^5 \) GeV fixed, a choice of a heavy messenger scale \( M_X \sim 10^{16} \) GeV also
pushes the vacuum expectation value of the spurion to $\langle F \rangle \sim 10^{21}$ GeV. This framework is similar to the one presented in Ref. [68, 74]. There it has been shown that an explicit realization requires certain conditions on the messenger scale.

From the above effective parameterisation, it is clear that for heavy messengers with mass $\sim M_{GUT}$ scale, the gravitational mediation contributions can become comparable and can significantly alter the spectrum. In the effective picture represented by Eq. (6), we can choose the gravity coefficients such that soft terms have the required form to give the correct Higgs mass at the weak scale. For example, using the analysis of [15], we can choose the gravity contributions to the stop sector and gluino sector such that they generate a large enough $A_t$ at the weak scale after RG running. Another possibility is to have a large $A_t$ at the $M_X$ scale which arises purely from gravity contributions. For instance, if large stop mixing with $A_t \sim -2500$ GeV can be generated at the Planck scale, then the corresponding weak scale value is sufficient to generate the right value for the stop mixing parameter with TeV scale stops. The $A$ terms at the messenger scale are effectively the RGE evolved contribution from the gravity sector. Of course, one has to then construct a gravity model which can lead to such soft terms, in this case, perhaps based on models presented in [45] and other string based models.

From the above discussion, one can conclude that gravitational contributions can indeed rescue minimal gauge mediation if the messenger masses are sufficiently high and they are of the right type. The question then arises what happens if the messenger masses are not so heavy and are only of the order of 100 TeV or so; Can gravitational contributions still play a role? This is the question we try to address now.

The advantage of such a ‘low’ scale for $M_X$ is that the nice features of gauge mediation are left intact. However, gravitational mediation contributions to the soft sector are small as discussed above. One simple way to increase the gravitational mediation would be to consider two different sectors of supersymmetry breaking. While only one of them couples to the messengers, both couple to gravity. In this set up, let us try to understand the character of the gravitational contributions.

We will consider two spurions $X_1$ and $X_2$ parameterising the two sectors where supersymmetry is broken. As before we denote $X_1 \equiv M_{X_1} + \theta^2 F_1$, which couples to the messengers at the scale $M_{X_1} \sim 100$ TeV. The other spurion $X_2$ which parameterises the sector which does not couple to messengers also contributes at $M_{Pl}$. In terms of the effective operators, we have just below $M_{Pl}$, the following contributions from the gravitational mediation:
\[ m_{ij}^2 = \int d^2 \theta d^2 \bar{\theta} \left( k_{ij} \frac{F_1^2 F_1}{M_{Pl}^2} + k'_{ij} \frac{F_2^2 F_2}{M_{Pl}^2} \right) \Phi^\dagger \Phi \]
\[ M_{ab} = f_{ab} \int d^2 \theta \left( \frac{F_1}{M_{Pl}} + \frac{F_2}{M_{Pl}} \right) W^a W^b \]
\[ A_{ijk} = \int d^2 \theta \left( \eta_{ijk} \frac{F_1}{M_{Pl}} + \eta'_{ijk} \frac{F_2}{M_{Pl}} \right) \Phi_i \Phi_j \Phi_k \]

(7)

where \( F_1 \ll F_2 \). At the scale \( M_{X_1} \) additional contributions from gauge mediation have the same form as in Eq.(3) with \( F \) replaced by \( F_1 \). It turns out that in this case, there is a nice way of organising the weak scale soft terms which are a result of the RGE evolution from \( M_{Pl} \) to \( M_{X_1} \) and then to \( M_{weak} \). Throughout this paper, we will use primed objects to denote the contributions from gauge mediation sector. Unprimed objects will denote those from gravity mediation. The total soft terms have a \( \tilde{} \) on them. At the weak scale, we have the following relations:

\[ \tilde{M}_a(weak) = M_a(weak) + M'_a(weak) \]
\[ \tilde{A}(weak) = A(weak) + A'(weak) \]

Thus both the gaugino masses and the trilinear terms just add up at the weak scale as though they have no knowledge about the presence of another supersymmetry breaking sector\(^2\).

The scalar masses however, as they depend quadratically on the gaugino masses do not add linearly. They have mixed contributions from both the supersymmetry breaking sectors. Schematically they can be written as (neglecting the Yukawa contributions)

\[ \tilde{m}_j^2(weak) = m_j^2(weak) + m'_j(weak) + \sum_a M_a M'_a \zeta_a \zeta'_a \]

(8)

\( \zeta(\zeta') \) are the RG factors from \( M_{X_2} = M_{Pl} \) to weak scale and \( M_{X_1} \) to weak scale respectively (for more details, see the Appendix A ).

The total solution can be schematically represented by the following two figures: Fig. (1) represents the total solutions for the gaugino masses \( M_i \) and the trilinears where as Fig. (2) represents the sfermion mass squared terms. The total gaugino masses and the A terms, at the weak scale, are just the linear sums of the weak scale solutions of the two independent sectors which is described in Fig. (1). The scalar mass terms at the weak scale, however, contain cross terms mixing

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*\(^2\) These results are true only up to one loop order. Issues about higher loops and multiple (greater than 2) SUSY breaking sectors will be considered elsewhere.*
FIG. 1: RGE running of gauginos and A terms in two scale SUSY breaking

FIG. 2: RGE running of scalar squared masses in two scale SUSY breaking

contributions from both of the sectors which is described in Fig. (2). The solution of the RGE presented here can be generalised easily for any scenario where supersymmetry is broken at two different scales and preserves the MSSM structure. Similarly it can be generalised to more than two sectors of supersymmetry breaking at different scales.

In Table [I], we give a numerical example of low energy spectrum due to supersymmetry breaking derived from two different hidden sectors i.e. gravity and gauge mediation. The first two rows are the low scale values of the soft masses from SUSY breaking due to gravity and gauge mediation respectively. The third row gives the low scale values when SUSY breaking due to both gravity and gauge mediation are present at two different scales. It is referred to as “Double SUSY breaking”. As expected, we find that the $A$ terms and the gaugino masses, at the weak scale, to be the linear sum of the gauge and gravity contributions individually\(^3\). The soft masses on the other hand do

\[^3\] The slight discrepancy in the $A$ terms due to individual sum of the two contributions and that due to double SUSY breaking can be attributed to the difference between numerical and analytic computations. It is independent of the boundary parameters.
not add up linearly. For illustration, we present the results for third generation squarks. The pure interference terms are listed in the fourth row. The interference terms could become important when the contributions to the soft terms due one of the sectors of SUSY breaking is sub-dominant.

\[
m_{\tilde{Q}^3}^2 \quad [GeV^2] \\
m_{\tilde{U}^3}^2 \quad [GeV^2] \\
m_{\tilde{D}^3}^2 \quad [GeV^2] \\
A_t [GeV] \\
M_1 [GeV] \\
M_2 [GeV] \\
M_3 [GeV]
\]

| Gravity mediation | 3.2×10^6 | 2.1×10^6 | 3.5×10^6 | -2091 | 456.5 | 846.9 | 2191.8 |
|-------------------|----------|----------|----------|-------|-------|-------|--------|
| Gauge mediation   | 1.04×10^6 | 0.84×10^6 | 1.01×10^6 | -280 | 139.0 | 263.7 | 722.94 |
| Double susy breaking | 4.67×10^6 | 3.43×10^6 | 5.04×10^6 | -2283.4 | 603.12 | 1110.09 | 2817.3 |
| Interference terms | 0.43×10^6 | 0.49×10^6 | 0.53×10^6 | \sim 0 | \sim 0 | \sim 0 | \sim 0 |

TABLE I: To illustrate the organizing principle discussed in the text, scalar squared masses, trilinear coupling \( A_t \) and gauginos are presented for different mediation schemes. Rest of the parameters are fixed as \( \tan \beta = 10 \) for all the schemes. For gravity mediation: \( m_0 = 500 \text{ GeV}, M_2 = 1000 \text{ GeV}, A_0 = -1000 \text{ GeV} \), For gauge mediation: \( \Lambda = 10^5 \text{ GeV} \) and \( M_{X_1} = 10^6 \text{ GeV} \)

Returning to the problem at hand, we demand that the gravitational contributions in this case should be "just right" to generate the right Higgs mass. This would mean that gravity should contribute significantly to \( A \) (mainly \( A_t \)) terms but not to anything else. Practically such a situation is hard to achieve unless one assumes a flavoured supersymmetry breaking in the gravity sector. Such examples are abound in literature either based on U(1) symmetries \[75, 76\] or extra dimensions \[77\]. In the following, we present an example based on Randall-Sundrum Models \[78\].

### III. A COMPLETE EXAMPLE WITH RANDALL SUNDRUM AT HIGH SCALE

The Randall-Sundrum (RS) setup consists of a single extra-dimension of radius \( R \) compactified on a \( S_1/Z_2 \) orbifold. Two branes, IR and UV are located at the \( y = \pi R \) and \( y = 0 \) orbifold fixed points. This model was originally proposed as solution to the hierarchy problem by means of a ‘tiny’ geometric warp factor defined as

\[
e = e^{-kR\pi} \sim 10^{-16}
\]

where \( k \) is the reduced Planck scale and \( kR \sim 11 \). The extent of warping is set by the radius \( R \) of the extra-dimension. Thus the radiative instability to the Higgs mass can ‘warped’ down to the electroweak scale by the relation \( M_{IR} = M_{ew} \sim e^{-kR\pi} M_{Pl} \).

In this example we consider a modification of the RS setup with a similar background geometry but with a radius \( R' \sim R/8 \). As a result the IR scale corresponds to the GUT \( M_{GUT} \) scale and hence the RS setup does not provide a solution to the gauge hierarchy problem. This can be
resolved by super-symmetrizing the bulk. Such a scenario can be explicitly realized in models with Randall-Sundrum (RS) like geometry and will be the focus of our attention in this section. Similar frameworks have been considered [79–84]. The spectrum of the effective effective 4D theory is that of MSSM which helps in stabilizing the radiative corrections to the Higgs mass.

We assume the two Higgs doublet $H_u$ and $H_d$ to be localized on the low energy brane (IR brane). Matter and gauge fields are in the bulk. Supersymmetry breaking spurion $X = \theta^2 F$ is introduced on the IR brane. SUSY breaking terms are generated by the brane local interactions of the spurion with the bulk fields. The $F$ term of the spurion $X$ develops a vacuum expectation value and generates the soft masses. In the canonical basis, the soft breaking terms are of the form [79–82]:

\[
m_{H_u, H_d}^2 = \hat{\beta}_{u,d} m_3^{3/2}
\]
\[
(m^2_f)_{ij} = m_3^{3/2} \beta_{ij} e^{(1-c_i-c_j)kR \pi} \xi(c_i)\xi(c_j)
\]
\[
A^u_{ij} = m_3^{3/2} A'_{ij} e^{(1-c_i-c'_j)kR \pi} \xi(c_i)\xi(c'_j)
\]
\[
m_{1/2} = k_h m_3^{3/2}
\]

(10)

where the gravitino mass is defined as

\[
m_{3/2}^2 = \frac{(F)^2}{k^2} = \frac{<F>^2}{M_{Pl}^2}
\]

(11)

$\hat{\beta}_{ij}, A', k_h$ are dimensionless $O(1)$ parameters. $\xi(c_i)$ is defined as

\[
\xi(c_i) = \sqrt{(0.5 - c_i) e^{-2(1-2c_i)kR}}
\]

(12)

We point here that the the $c$ parameters which determine the soft masses in Eq.(10) are constrained by their requirement to fit fermion mass and mixing angles at $M_{GUT}$. Since the two Higgs doublets are localized at the IR brane, the bulk mass parameters for the light fields are $> 0.5$ such that their zero modes are localized near the Planck brane. Alternatively a configuration where the doublet and the singlets are localized away from each other in the bulk will also lead to small fermion masses. For instance the doublet can be localized close to the IR brane while the singlets are close to the UV brane and vice versa. Since the Higgs doublets and the SUSY breaking spurion are localized on the same brane, a choice of $c$ values which fits small masses to lighter generations also ensures that the corresponding sfermion masses are also small at $M_{GUT}$. On the other hand the $c$ parameters $c_{Q_3,U_3}$ are typically less than or close to 0.5 which is required to accommodate a large top quark mass. Noting that the structure of the trilinear coupling in Eq.(10) is similar to
Hadron | Parameter | Value | Hadron | Parameter | Value
------ | -------- | ----- | ------ | -------- | -----
| \(c_{Q_1}\) | 1.494 | \(c_{D_1}\) | 1.524 |
| \(c_{Q_2}\) | 1.099 | \(c_{D_2}\) | 1.686 |
| \(c_{Q_3}\) | -1.593 | \(c_{D_3}\) | 1.287 |
| \(c_{U_1}\) | 2.005 | \(c_{U_2}\) | 1.896 |
| \(c_{U_3}\) | 0.710 | - | - |

Lepton | Parameter | Value | Lepton | Parameter | Value
------- | -------- | ----- | ------- | -------- | -----
| \(c_{L_1}\) | 0.507 | \(c_{E_1}\) | 2.250 |
| \(c_{L_2}\) | -0.630 | \(c_{E_2}\) | 2.109 |
| \(c_{L_3}\) | -0.997 | \(c_{E_3}\) | 1.329 |
| \(-\) | - | \(-\) | - |

TABLE II: Example point of the bulk mass parameters which satisfy both fermion mass fits.

the mass matrix, we find that fitting a large top quark mass will also ensure large \(A_t\) necessary to generate Higgs mass for low \(m_{susy}\).

We now provide a numerical example of this case where we use a set of \(c\) parameters to determine the boundary conditions for the soft masses at \(M_{GUT}\). The results of the analysis of [84] are used where the technique of \(\chi^2\) minimization was used to determine the \(c\) parameters. A particular choice of \(c\) parameters which we use for our analysis is given in Table II.

We use a slightly modified version of SUSEFLAV [85] to determine the low energy spectrum. Table III gives the low energy spectrum corresponding to gravity contribution at \(M_{Pl}\) and gauge contribution at the messenger scale \(M_{X_1} = 10^6\). The gravitino mass is chosen to be \(m_{3/2} = \langle F_1 \rangle \frac{M_{X_1}}{M_{X_2}} = 1000\) GeV. At the messenger scale \(M_{X_1}\), we choose \(\Lambda = \langle F_1 \rangle \frac{M_{X_1}}{M_{X_1}} = 10^5\). The boundary conditions for the gaugino masses from the gravity sector are chosen to be 400 GeV to limit the contribution to the soft masses due to gravity mediation. This corresponds to a choice of \(k_h = 0.4\) in Eq.(10). Additionally we choose \(\beta_{ij} = 1 \forall i,j\) for the soft masses. \(A'_{ij} = Y'_{ij}\) except for \(i = j = 3\) where \(Y'_{ij}\) is the \(O(1)\) Yukawa coupling which fits the fermion masses. \(A'_{33} = 3.4Y'_{33}\) is chosen to further enhance the trilinear couplings. From the table we see that it is possible to accommodate a heavy Higgs mass of \(\sim 125\) GeV in a low scale gauge gauge mediation scenario with fairly light stop masses. The rest of the spectrum including the gluinos are reasonably light owing to the suppressed contribution from the gravity sector. Thus the role played by flavourful gravity mediation is to only generate large trilinear coupling by RG evolution at the gauge messenger scale.

The origin of flavourful soft masses at the GUT scale could potentially lead to large flavour violation at the low scale \(^4\). The most stringent constraints are due to transitions between the first two generations in both the leptonic and the hadronic sector. Since the lighter generations are localized away from source of SUSY breaking spurion \(X\) (and the Higgs) the off-diagonal elements

\(^4\) Since the lowest KK masses are \(O(M_{GUT})\) they do not contribute to the flavour processes.
TABLE III: Soft spectrum for Dirac case: $m_{susy} = 1.13$ TeV, $m_g = 1.67$ TeV, $\tan\beta = 10$

| Parameter | Mass (TeV) | Parameter | Mass (TeV) | Parameter | Mass (TeV) | Parameter | Mass (TeV) |
|-----------|------------|-----------|------------|-----------|------------|-----------|------------|
| $\tilde{t}_1$ | 0.739 | $\tilde{b}_1$ | 1.51 | $\tilde{t}_1$ | 0.508 | $\tilde{\nu}_\tau$ | 2.58 | $N_1$ | 0.304 |
| $\tilde{t}_2$ | 1.93 | $\tilde{b}_2$ | 1.92 | $\tilde{\tau}_2$ | 2.58 | $\tilde{\nu}_\mu$ | 0.339 | $N_2$ | 0.600 |
| $\tilde{c}_R$ | 1.43 | $\tilde{s}_R$ | 1.52 | $\tilde{\mu}_R$ | 0.351 | $\tilde{\nu}_e$ | 0.327 | $N_3$ | 1.04 |
| $\tilde{c}_L$ | 1.57 | $\tilde{s}_L$ | 1.58 | $\tilde{\mu}_L$ | 0.548 | - | - | $N_4$ | 1.05 |
| $\tilde{u}_R$ | 1.43 | $\tilde{d}_R$ | 1.52 | $\tilde{e}_R$ | 0.347 | - | - | $C_1$ | 0.577 |
| $\tilde{u}_L$ | 1.57 | $\tilde{d}_L$ | 1.58 | $\tilde{e}_L$ | 0.548 | - | - | $C_2$ | 1.04 |
| $m_{A_0}$ | 1.46 | $m_H^\pm$ | 1.46 | $m_h$ | 0.1259 | $m_H$ | 1.46 | - | - |

as well the diagonal elements are very small at the GUT scale. Consider the flavour violating parameter $\delta^f_{ij}$ defined as

$$\delta^f_{ij}(i \neq j) = \frac{(U^*\tilde{m}^2_{susy}U)_{ij}}{\tilde{m}^2_{susy}(i \neq j)}$$

$U$ is rotation matrix which rotates from flavour basis to the mass basis and $\tilde{m}_{susy}$ is the geometric mean of diagonal elements $\tilde{m}^2_{ii}$ and $\tilde{m}^2_{jj}$, $i \neq j$. Since the diagonal elements $\tilde{m}^2_{ii}$ and $\tilde{m}^2_{jj}$ are small at $M_{GUT}$, $\delta^f_{ij}$ is large and $\mathcal{O}(1)$ at the GUT scale.

The diagonal elements however, receive significant contributions due to RGE effects and hence are large at $m_{susy}$. The off-diagonal elements on the other hand do not evolve much and are small even at $m_{susy}$. Thus, while there could be large flavour violation at $M_{GUT}$, $\delta^f_{ij}$ reduces from $\mathcal{O}(1)$ values at the high scale to values consistent with the experimental bounds at the low scale.

To conclude this section, we point out that it is also possible to generate flavourful soft terms in 4D models. One example in 4D is to consider Froggatt-Nielsen (FN) type of models which have an additional horizontal family symmetry $U(1)_X$ [86]. They were introduced to explain the observed hierarchy in the fermion mass and the mixing angles by assigning different charges to different generations of fermion multiplets under $U(1)_X$. Requiring the Kahler potential be canonical, rendered a non-trivial flavour structure to the soft mass sector. One drawback of these models is that $U(1)_X$ group is anomalous and they are only cancelled by the Green-Schwarz anomaly cancellation mechanism [87]. This severely restricts the parameter space of the FN charges of the multiplets which satisfy the fermion mass fits as well as the anomaly cancellation requirement. Further these solutions are also constrained by FCNC processes. The Randall-Sundrum framework is a more generalised setup which is equivalent to the setup of $U(1)_{FN}$ but not as constraining.
IV. CONCLUSIONS

Gravitational contributions to supersymmetry breaking are typically suppressed in cases where gauge mediated supersymmetry breaking contributions are dominant. In the present work, we show that however this is not the case if the supersymmetry breaking happens in multiple sectors which are parameterised in terms different spurion fields. If they are hierarchal, one can have a situation where gravity contributions can play a supporting role to the gauge mediated contributions and help in generating the right Higgs mass. Of course, this would require that the gravitational contribution should be flavourful which we have achieved by using a Randall Sundrum set up. In the course of this work, we have presented simple “rules of thumb” to RGE equations when supersymmetry is broken by multiple sectors at multiple scales.

Appendix A: Analytic expressions for soft masses in models with two sectors of SUSY breaking

In this section we give the analytic expressions for the soft masses at the weak scale due to SUSY breaking from two different hidden sectors. To begin with, consider the gauginos and the A terms. The net contribution at any scale due to the presence of both gauge and gravity sectors is just the linear sum of individual contributions due to pure gauge and pure gravity mediation. We define the following useful functions:

\[ z_i(t) = \frac{1}{(1 + b_i \tilde{\alpha}_i(0) t)} \quad z_i'(t) = \frac{1}{(1 + b_i \tilde{\alpha}_i(g) t)} \]
\[ k_i(t) = \frac{1}{2b_i} (1 - z_i(t)^2) \quad k_i'(t) = \frac{1}{2b_i} (1 - z_i'(t)^2) \quad (A1) \]

where

\[ t_g = 2 \log \left( \frac{M_{X_2}}{M_{X_1}} \right) \quad t_z = 2 \log \left( \frac{M_{X_2}}{m_{\text{susy}}} \right) \]

and \( \tilde{\alpha}_i(t) = z_i(t) \tilde{\alpha}_i(0) \) and \( \tilde{\alpha}_i(0) \) is the value at \( M_{X_2} \). Without loss of generality we can assume \( m_{\text{susy}} < M_{X_1} < M_{X_2} \).

The solutions to the one loop renormalization group evolution equations (RGE) for the gauginos...
can be solved exactly and are given as

\[ \tilde{M}_i(m_{\text{susy}}) = M_i z_i(t_z) + M_i' z_i'(t_z) \]

where \( M_o(M_o') \) are the boundary values for the gauginos at scale \( M_{X_2}(M_{X_1}) \). In our notation quantities with \( \sim \) represent the expressions at the \( m_{\text{susy}} \) scale due two sectors of SUSY breaking.

Similarly the \( A \) terms also add up linearly and are given as \(^5\)

\[
\begin{align*}
\tilde{A} &= A(t_z) + A'(t_z) \\
A(t_z) &= 0.33 A_0 - 0.03 M_1 - 0.26 M_2 - 1.8 M_3 \\
A'(t_z) &= -0.002 \Lambda
\end{align*}
\]

where \( A_0 \) is the boundary value for the trilinear coupling due to gravity mediation.

Next we consider the lighter generations. In the limit the Yukawa coupling \( Y_{b,\tau} \rightarrow 0 \), the solutions to the RGE equations for the first two generation squarks and the sleptons take a very simple analytic form. The solutions for the squarks and \( \tilde{m}_{H_u}^2 \) except \( \tilde{m}_{Q_3, U_3, H_u}^2 \) is given as

\[
\begin{align*}
\tilde{m}_{Q_{1,2}}^2 &= m_{Q_{1,2}}^2(t_z) + m_{Q_{1,2},1}(t_z) + \frac{32}{3} M_3^2 M_3 z_3(t_g) k_3'(t_z) + 6 M_2^2 M_2 z_2(t_g) k_2'(t_z) + \frac{2}{15} M_1^4 M_1 z_1(t_g) k_1'(t_z) \\
\tilde{m}_{U_{1,2}}^2 &= m_{U_{1,2}}^2(t_z) + m_{U_{1,2},1}(t_z) + \frac{32}{3} M_3^2 M_3 z_3(t_g) k_3'(t_z) + \frac{32}{15} M_1^4 M_1 z_1(t_g) k_1'(t_z) \\
\tilde{m}_{D_{1,2,3}}^2 &= m_{D_{1,2,3}}^2(t_z) + m_{D_{1,2,3},1}(t_g z) + \frac{32}{3} M_3^2 M_3 z_3(t_g) k_3'(t) + \frac{8}{15} M_1^4 M_1 z_1(t_g) k_1'(t) \\
\tilde{m}_{H_d}^2 &= m_{H_d}^2(t_z) + m_{H_d,1}(t_z) + 6 M_2^2 M_2 z_2(t_g) k_3'(t_z) + \frac{6}{5} M_1^4 M_1 z_1(t_g) k_1'(t_z)
\end{align*}
\]

The corresponding equations for the sleptons can be similarly written as

\[
\begin{align*}
\tilde{m}_L^2 &= m_L^2(t_z) + m_{L,1}(t_z) + 6 M_2^2 M_2 z_2(t_g) k_2'(t_z) + \frac{6}{5} M_1^4 M_1 z_1(t_g) k_1'(t_z) \\
\tilde{m}_E^2 &= m_E^2(t_z) + m_{E,1}(t_z) + \frac{24}{15} M_1^4 M_1 z_1(t_g) k_1'(t_z)
\end{align*}
\]

1. Semi-analytic expression for \( m_{Q_3, U_3, H_u}^2 \)

The expressions for \( m_{Q_3, U_3, H_u}^2 \) are more complicated as it involves the top quark Yukawa coupling which cannot be set equal to zero. However they follow the generic form of Eq. (8). We present

\(^5\) The semi-analytic expressions are for \( M_{X_1} = 10^6 \) and \( M_{X_2} = M_{Pl} \)
semi-analytic expressions for them as follows:

\[
\begin{align*}
\tilde{m}_{Q_3}^2 &= m_{Q_3}^2(t_z) + m_{Q_3}^2(t_z) + 0.0001A_0\Lambda + 0.0054M_3\Lambda \\
\tilde{m}_{U_3}^2 &= m_{U_3}^2(t_z) + m_{U_3}^2(t_z) + 0.0004A_0\Lambda + 0.0045M_3\Lambda \\
\tilde{m}_{H_u}^2 &= m_{H_u}^2(t_z) + m_{H_u}^2(t_z) + 0.0006A_0\Lambda - 0.0041M_3\Lambda
\end{align*}
\] (A7)

where \( \Lambda = \frac{F_1}{M_{X_1}} \). In the above expression we have neglected the terms proportional to \( M_1 \) and \( M_2 \) as their co-efficients are small compared to that of \( M_3 \). The semi-analytic expressions for the soft terms at \( m_{\text{susy}} \) scale due to pure gravity mediation is given as

\[
\begin{align*}
m_{Q_3}^2 &= -0.0369233A_0^2 + 0.00443031A_0M_1 + 0.00154387M_1^2 + 0.0266445A_0M_2 - 0.00205633M_1M_2 \\
&+ 0.233435M_2^2 + 0.1096A_0M_3 - 0.00916469M_1M_3 - 0.0610636M_2M_3 + 3.31617M_3^2 \\
&+ 0.3312m_{Q_3}^{(0)} - 0.11145(m_{H_u}^{(0)} - 5m_{Q_3}^{(0)} + m_{U_3}^{(0)}) \\
m_{U_3}^2 &= -0.0738465A_0^2 + 0.00886063A_0M_1 + 0.0287372M_1^2 + 0.0532891A_0M_2 - 0.00411266M_1M_2 \\
&- 0.0134655M_2^2 + 0.2192A_0M_3 - 0.0183294M_1M_3 - 0.1221727M_2M_3 + 3.16522M_3^2 \\
&- 0.222901(m_{H_u}^{(0)} + m_{Q_3}^{(0)} - 2m_{U_3}^{(0)}) + 0.3312m_{U_3}^{(0)} \\
m_{H_u}^2 &= -0.11077A_0^2 + 0.0132909A_0M_1 + 0.0156242M_1^2 + 0.0799336A_0M_2 - 0.00616898M_1M_2 \\
&+ 0.21997M_2^2 + 0.3288A_0M_3 - 0.0274941M_1M_3 - 0.183191M_2M_3 - 0.452999M_3^2 \\
&+ 0.3312m_{H_u}^{(0)} - 0.334351(-m_{H_u}^{(0)} + m_{Q_3}^{(0)} + m_{U_3}^{(0)})
\end{align*}
\] (A8)

where \( m_{Q_3, U_3, H_u}^{(0)} \) represents the boundary value of the soft masses due to gravity mediation. The corresponding expressions due to gauge mediation with messenger scale at \( M_{X_1} = 10^6 \) is given as

\[
\begin{align*}
m_{Q_3}^2 &\approx 1 \times 10^{-4}\Lambda^2 \\
m_{U_3}^2 &\approx 7.62 \times 10^{-5}\Lambda^2 \\
m_{H_u}^2 &\approx -4.65 \times 10^{-5}\Lambda^2
\end{align*}
\] (A9)

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