

$SO(10)$ inspired extended GMSB models

Tomasz Jeliński*

*Institute of Physics, University of Silesia,
Uniwersytecka 4, PL-40-007 Katowice, Poland
E-mail: tomasz.jelinski@us.edu.pl*

Influence of messenger-matter superpotential interactions on the renormalization of Yukawa couplings in the context of extended GMSB models is analysed. We present a convenient method for treating decoupling of messengers and related redefinition of MSSM fields. Discussed approach is used to study top-bottom-tau Yukawa unification within specific $SO(10)$ inspired GUT model.

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1. Introduction

Among supersymmetric GUT models those which involve the Gauge Mediated Supersymmetry Breaking (GMSB) mechanism [1] are of great interest for a long time. They exhibit many desired features, but they also require rather large masses of sparticles to accommodate for 125 GeV Higgs mass [2]. Recently, it was shown that this issue can be naturally addressed in so-called extended GMSB (EGMSB) models, in which messengers interact with MSSM fields through superpotential couplings [3, 4]. Such couplings are natural because messengers carry the same charges under $G_{SM} = SU(3) \times SU(2) \times U(1)$ as MSSM fields hence they can interact through superpotential in an analogous way. These couplings generate 1-loop A -terms (and 1- and 2-loop soft masses) at the messenger scale M [3, 4, 5, 6, 7, 8, 9]. In consequence, in many realizations of the EGMSB models stop mixing at the electroweak symmetry breaking (EWSB) scale is enhanced what in turn helps to get proper Higgs mass [10, 11, 12, 13]. The other motivation for considering these models is that they provide a reasonable framework to explore non-minimal flavour violation scenarios [13, 14, 15, 16, 17, 18], e.g. those which are inspired by the F-theory geometrical constructions [19, 20, 21] or Froggatt-Nielsen type models [22]. As a consequence of extra contributions to soft masses, EGMSB models have also quite rich and interesting phenomenology, where non-standard NLSP/NNLSP patterns like stop/bino or sneutrino/stau [11] can be realized.

It is clear that in the GMSB models the presence of messengers changes running of gauge couplings g_r above messenger scale M with respect to the standard RGE evolution of MSSM parameters. It is well known that when messengers are in full representations of GUT gauge group then unification of g_r is not spoiled by their presence [1]. Moreover, via RG equations, they also indirectly change running of top, bottom and tau Yukawa couplings $y_{t,b,\tau}$ [23]. Hence, it is natural to raise a question about Yukawa couplings unification when one allows for superpotential couplings between messengers and matter, as in EGMSB models. That issue can be rephrased also in the following form. Assuming unification of Yukawa couplings $y_{t,b,\tau}$ at the GUT scale $M_{GUT} \approx 2 \times 10^{16}$ GeV, one can ask what are their values at the EWSB scale and do they match values derived from fermion masses [24].

We show that the method [8] which was developed to properly derive soft terms in EGMSB models can be used to get information about the running of Yukawa couplings in those scenarios. It is especially useful in the top-down approach where one assumes precise Yukawa couplings unification at the GUT scale M_{GUT} and derives their values at the EWSB scale from RGE evolution. That method allows to keep track of messenger-matter kinetic mixing and study how messengers superpotential couplings contribute to values of Yukawas at the EWSB scale. We apply our method to analyse the issue of top-bottom-tau unification in the context of specific $SO(10)$ inspired GUT model with minimal matter content and only one messenger-matter superpotential coupling. Such $SO(10)$ -type GUT model, see e.g. [25, 26, 27], provides a natural setup to present the method and expose issues related to the kinetic mixing of messengers and MSSM matter triggered by superpotential couplings.

2. Extended GMSB models

Let us briefly recall the salient features of so-called extended GMSB models. For more details

see e.g. [3, 4, 5, 6, 8, 9, 28]. As in the standard GMSB models [1] supersymmetry is broken in the hidden sector by the F -term of the spurion superfield X . Messenger sector consists of chiral superfields $Y_A = (Y_A, \bar{Y}_A)$ in vectorial representations of G_{SM} . We assume that all the messenger fields couple in the same way to the spurion X through a superpotential term $XY_A\bar{Y}_A$ [29]. When X gets vev $\langle X \rangle = M + \theta^2 F_X$ that coupling gives mass M to the messengers and induces additional SUSY breaking masses $F_X Y \bar{Y}$ for their scalar components. For simplicity, let us consider messenger fields in the following representations of G_{SM} :

$$\begin{aligned} Y_{H_u} &: (1, 2)_{1/2}, & Y_Q &: (3, 2)_{1/6}, & Y_{\bar{U}} &: (\bar{3}, 1)_{-2/3}, \\ Y_{\bar{D}} &: (\bar{3}, 1)_{1/3}, & Y_L &: (1, 2)_{-1/2}, & Y_{\bar{E}} &: (1, 1)_1 \end{aligned} \quad (2.1)$$

and their partners in conjugate representations: $Y_{\bar{H}_u}(1, 2)_{-1/2}$, $Y_{\bar{Q}}(\bar{3}, 2)_{-1/6}$, $Y_U(3, 1)_{2/3}$ etc. We have not included Y_{H_d} in the list (2.1) as it would have the same quantum numbers as Y_L . Moreover, we also allow for a singlet messenger $Y_{N_R}(1, 1)_0$ under G_{SM} . Our choice of messengers is motivated by most common constructions of GUT models based on $SU(5)$ and $SO(10)$, where after breaking of G_{GUT} to G_{SM} all light fields are only in representations shown in (2.1). Let us note that more complicated setups are also possible e.g. when messengers are in higher-dimensional representations of G_{SM} . For example, when spectrum contains messenger Y_S in $(\bar{6}, 1)_{-1/3}$ of G_{SM} then the following coupling is possible $QY_S Q$.

In the EGMSB models messengers Y_A couple to MSSM fields Φ_a not only by gauge fields but also directly through the following superpotential:

$$\begin{aligned} W_3 &= \frac{1}{6} \lambda_{ijk} \Phi_i \Phi_j \Phi_k \\ &= \frac{1}{6} \gamma_{abc} \Phi_a \Phi_b \Phi_c + \frac{1}{2} h_{abc} \Phi_a \Phi_b Y_c + \frac{1}{2} h_{aBC} \Phi_a Y_B Y_C + \frac{1}{6} \eta_{ABC} Y_A Y_B Y_C, \end{aligned} \quad (2.2)$$

what leads to generating 1-loop A-terms and 1- and 2-loop soft masses [3, 4, 5, 8, 9]. We shall focus on the regime $F/M^2 \ll 1$ in which 2-loop soft masses dominate over 1-loop soft masses [3, 4]. Let us note that when λ_{ijk} couplings are zero then one gets the standard GMSB model. For discussion of the phenomenology of the EGMSB models see e.g. [10, 11, 12, 13].

3. Decoupling and redefinition of fields

It is clear that in the presence of superpotential couplings (2.2) between messengers and matter both fields are subject to kinetic mixing. Let us note that in the models under consideration there are two types of such mixing: (i) between two fields ($H_u \leftrightarrow Y_{H_u}$, $Q \leftrightarrow Y_Q$, $U \leftrightarrow Y_U$), and (ii) between three fields ($H_d \leftrightarrow L \leftrightarrow Y_L$) when down-type Higgs field H_d mixes with L and messenger field Y_L . For example, messenger field Y_Q with the same quantum numbers as quark field Q can mix with the latter when superpotential contains $H_u Y_Q U$ or $H_d Y_Q D$. As a consequence of kinetic mixing, wavefunction renormalization Z_{ij} receives non-diagonal loop corrections, even if it was chosen to be diagonal at the GUT scale M_{GUT} . That in turn, according to the non-renormalization theorem [30], influences running of Yukawa couplings $\tilde{\lambda}_{ijk}$ between canonically normalized fields $\tilde{\Phi}_i$. A method presented in [8], originally developed to properly compute 2-loop soft-masses induced by (2.2), is especially convenient to reveal influence of messenger-matter interactions on Yukawa couplings

\tilde{y}_{abc} renormalization. At the scale¹ $t = \ln \mu / Q$, lower than the GUT scale $t_{GUT} = \ln M_{GUT} / Q$, the model is characterized by the superpotential and the Kähler potential:

$$W = \frac{1}{6} \lambda_{ijk} \Phi_i \Phi_j \Phi_k + \frac{1}{2} M_{ij} \Phi_i \Phi_j, \quad K = \Phi_i^\dagger Z_{ij}(t) \Phi_j, \quad (3.1)$$

where Z_{ij} is positive-definite Hermitian wave-function renormalization matrix and $Z(t_{GUT})_{ij} = \delta_{ij}$, while M_{ij} is complex symmetric mass matrix. M_{ij} is chosen such that at scale t_{GUT} messenger fields Y_A are massive i.e. they have mass terms of the form $MY_A \bar{Y}_A$ while MSSM fields Φ_a are light.² In the so-called holomorphic scheme [8], in which Kähler potential is not canonical (3.1), couplings λ_{ijk} do not run. On the other hand, it is clear that the running couplings $\tilde{\lambda}_{ijk}(t)$ and masses $\tilde{M}_{ij}(t)$ of canonically normalized fields $\tilde{\Phi}_i = Z_{ij}^{-1/2} \Phi_j$ are related to their holomorphic counterparts, λ_{ijk} and M_{ij} , in the following way:

$$\tilde{\lambda}_{ijk}(t) = \lambda_{i'j'k'} Z_{i'i}^{-1/2} Z_{j'j}^{-1/2} Z_{k'k}^{-1/2}, \quad \tilde{M}_{ij}(t) = M_{i'j'} Z_{i'i}^{-1/2} Z_{j'j}^{-1/2}. \quad (3.2)$$

Hence non-diagonal Z_{ij} , generated by loop corrections, induce mixing mass terms between messenger fields and MSSM fields. In other words: RGE evolution of Z_{ij} reintroduces mixing mass terms between both fields. Hence to consistently define theory below messenger scale $\tilde{M}(t_M) = M$, one has to recognize light states (matter fields) and heavy states (messengers), properly decouple the latter at M , and appropriately derive couplings between light states (physical Yukawa couplings). To do that, one has to know which combination of original fields Φ_i are massive and which are light. One way to do this is to, first, diagonalize Z_{ij} and then find null vectors³ of mass matrix \tilde{M} (3.2) and orthogonal combinations (heavy states). The other method, which seems to be more convenient in the discussed case, is to use Cholesky decomposition [31] of Z_{ij} , which at once allows to find such a combination of holomorphic fields which stays heavy or massless under loop corrections. Here, instead of computing $Z^{-1/2}$ and then finding null vectors of \tilde{M} and their orthogonal combinations, one defines physical fields in the following way:

$$\tilde{\Phi} = \begin{pmatrix} \tilde{\phi} \\ \tilde{Y} \end{pmatrix} = V \begin{pmatrix} \phi \\ Y \end{pmatrix} = \begin{pmatrix} V_{11} & V_{12} \\ 0 & V_{22} \end{pmatrix} \begin{pmatrix} \phi \\ Y \end{pmatrix}, \quad (3.3)$$

where V is an upper triangular matrix which enters Cholesky decomposition of Z :

$$Z = V^\dagger V. \quad (3.4)$$

In (3.3) we displayed an example of such redefinition when only two fields, ϕ and Y , are subject to kinetic mixing $\phi \leftrightarrow Y$. Generalization to mixing of three or more fields is straightforward. As a consequence, the running physical superpotential couplings $\tilde{\lambda}_{ijk}(t)$ can be now written as:

$$\tilde{\lambda}_{ijk}(t) = \lambda_{i'j'k'} V_{i'i}^{-1} V_{j'j}^{-1} V_{k'k}^{-1}. \quad (3.5)$$

¹ μ is a renormalization scale, while Q is an arbitrary scale which we set to 1 GeV.

²'Massive' means that fields Y_A have mass of the order of M , while 'light' means that fields Φ_a are massless or have a supersymmetric mass of the order of EWSB scale.

³Or, more generally, light states.

Let us stress that the advantage of such approach is that, now, running couplings $\tilde{\lambda}_{abc}(t)$ between light fields $\tilde{\Phi}_a$, i.e. physical Yukawa couplings, are expressed only in terms of wave-function renormalization without the need for finding null vectors of the mass matrix \tilde{M} . Moreover, using (3.5) one can show that $\tilde{\lambda}_{abc}(t)$ are completely determined by the Z_{ij} running:

$$\tilde{\lambda}_{abc}(t) = \frac{\lambda_{abc}}{\sqrt{Z_{aa}(t)Z_{bb}(t)Z_{cc}(t)}}, \quad (3.6)$$

with no sum over repeated indices. It is known [8, 32] that Z_{ij} evolves according to the following one-loop renormalization group equation:

$$\frac{d}{dt}Z_{ij} = -\frac{1}{8\pi^2} \left(\frac{1}{2}d_{kl}\lambda_{ikl}^*Z_{km}^{-1*}Z_{ln}^{-1*}\lambda_{jmn} - 2C_{ij}^{(r)}Z_{ij}g_r^2 \right), \quad (3.7)$$

where d_{kl} and $C_{ij}^{(r)}$ are numerical factors appearing in the one-loop anomalous dimensions [8, 32], and g_r , ($r = 1, 2, 3$) are running gauge couplings of $U(1)$, $SU(2)$ and $SU(3)$, respectively. As mentioned above, the initial condition for (3.7) is chosen to be $Z_{ij}(t_{GUT}) = \delta_{ij}$. Beside the simplest case, it is hard to find an analytical solution of (3.7). One of the possible ways out is to expand Z_{ij} in power series:

$$Z_{ij}(t) = \sum_{n=0}^{\infty} Z_{ij}^{(n)}(t_{GUT}). \quad (3.8)$$

One can show that to compute $Z_{ij}^{(n)}(t_{GUT})$, i.e. the n -th derivative of Z_{ij} at t_{GUT} , it is enough to know all $Z_{ij}^{(k)}$ with $k < n$. As a result of such iterative procedure, $Z_{ij}^{(n)}$ can be expressed in terms of small parameter $\varepsilon = \ln 10/16\pi^2$, numerical factors d_{kl} and $C_{ij}^{(r)}$, and values of holomorphic couplings λ_{ijk} , which match values of running physical couplings at t_{GUT} , and gauge couplings $g_r(t_{GUT})$ at the GUT scale. Two comments are in order here. Let us note that this method is especially useful for studying unification of gauge and Yukawa couplings because values of physical Yukawa couplings at low scale (e.g. the EWSB scale) can be expressed via their unified values, y and g_{GUT} , numerical factors ε and value of messenger scale M . The advantage of this method is that it allows to keep track how messenger-matter couplings enter values of Yukawa couplings at lower scales. An explicit example of described method is given in the next section.

It is instructive to compare the running of top Yukawa coupling $\tilde{y}_t(t)$ when one neglects kinetic mixing above M and evolves $\tilde{y}_t(t)$ according to the standard RGE [32] with the running of $\tilde{y}_t(t)$ derived from (3.5). To this end, let us consider the following superpotential with only one coupling, $H_u Q Y_{\overline{U}}$, between messenger $Y_{\overline{U}}$ and MSSM fields:

$$W = y_t H_u Q \overline{U} + h_t H_u Q Y_{\overline{U}} + M Y_U Y_{\overline{U}}. \quad (3.9)$$

Clearly, $Y_{\overline{U}}$ has the same charges as \overline{U} and can mix with the latter.⁴ Here the messenger scale M is set to $M = 10^{10}$ GeV, while values of the superpotential couplings at the GUT scale are the

⁴To ensure gauge couplings unification, the spectrum has to contain also Y_Q , $Y_{\overline{E}}$ and their partners in conjugate representations. For simplicity, we assume that the superpotential couplings of those fields are suppressed with respect to the $H_u Q Y_{\overline{U}}$ coupling.

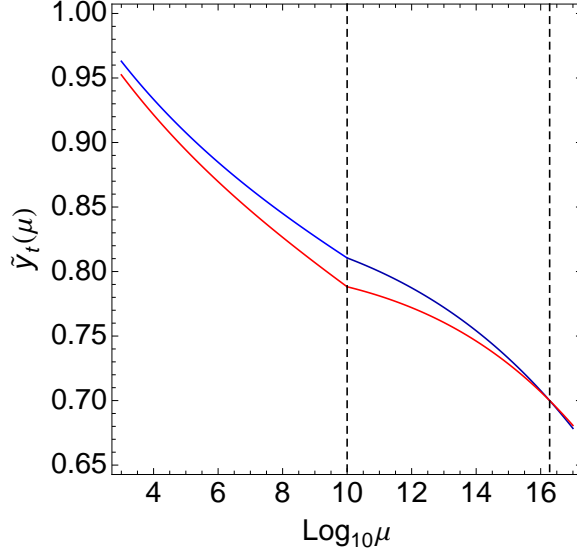


Figure 1: Running of physical top Yukawa coupling $\tilde{y}_t(\mu)$ when the mixing between messenger $Y_{\bar{U}}$ and MSSM field \bar{U} is taken into account (blue) and when that effect is neglected (red). μ is renormalization scale in units of GeV. Messenger scale M is set to 10^{10} GeV (left vertical dashed line), while messenger-matter coupling is $h = 0.4$. Value of top Yukawa coupling at GUT scale $M_{GUT} \approx 10^{16}$ GeV (right vertical dashed line) is set to 0.7. Below the scale M , the top Yukawa coupling \tilde{y}_t runs according to the standard RG equation of the MSSM.

following: $y_t(t_{GUT}) = y_t = 0.7$, $h_t(t_{GUT}) = h_t = 0.4$. Both solutions are presented on the Fig. 1, where the blue curve shows $\tilde{y}_t(\mu)$ derived from (3.5), while the red curve corresponds to the running obtained from standard RG equations without taking into account kinetic mixing. As one can see on Fig. 1, the difference between the two values of $\tilde{y}_t(\mu)$ grows when the renormalization scale μ decreases from M_{GUT} to M , and then it becomes smaller when $\tilde{y}_t(\mu)$ is evolved from M down to the EWSB scale. The discrepancy between both values of top Yukawa coupling at the messenger scale M is about 2%. Let us stress that below the messenger scale Yukawa coupling $\tilde{y}_t(\mu)$ runs according to the standard RGE of MSSM. If the kinetic mixing above M is neglected then to get right value of \tilde{y}_t at the EWSB scale one has to compensate the deficit at M by a small correction which is precisely a shift between the blue and red curve at the messenger scale M , see Fig. 1.

4. An example of SO(10) GUT model

Now, let us consider the following $SO(10)$ inspired GUT model. The spectrum of that model consists of one chiral field H_{10} in representation 10 of $SO(10)$, two chiral fields, ϕ_{16} and Y_{16} , in representation 16, and one field, $Y_{\bar{16}}$, in representation $\bar{16}$. We assume that below the GUT scale M_{GUT} , which is set by the condition $g_r(t_{GUT}) = g_{GUT}$, the $SO(10)$ gauge symmetry is broken down to $SU(5) \times U(1)_\chi$ and further to G_{SM} . Under $SO(10) \rightarrow SU(5) \times U(1)_\chi$ the chiral fields decompose as follows:

$$10 \rightarrow 5_2 + \bar{5}_{-2}, \quad 16 \rightarrow 10_{-1} + \bar{5}_3 + 1_{-5}, \quad (4.1)$$

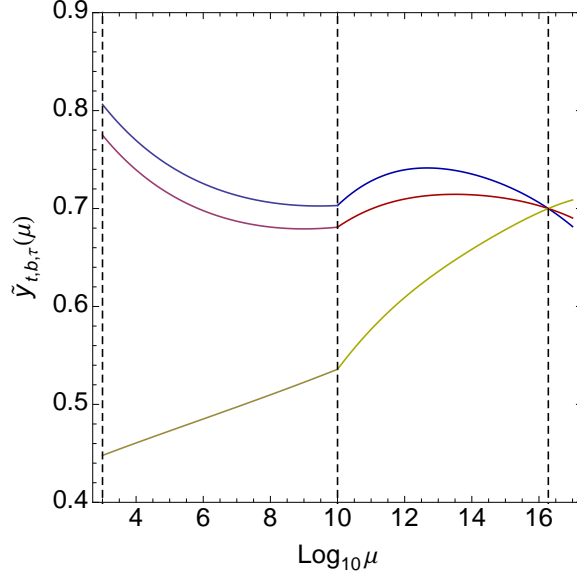


Figure 2: Running of top (blue curve), bottom (magenta curve) and tau (yellow curve) Yukawa couplings $\tilde{y}_{t,b,\tau}(\mu)$ in the $SO(10)$ inspired GUT model discussed in the Sec. 4. μ is the renormalization scale in GeV units. Messenger scale is set to 10^{10} GeV (middle vertical dashed line). Unified value of Yukawa couplings is set to $y_{GUT} = 0.7$ while messenger-matter coupling is $h = 0.4$. Left vertical dashed line corresponds to the EWSB scale, while the right vertical dashed line shows the GUT scale.

where subscripts denote $U(1)_\chi$ charges. Additionally, let us assume that after breaking of $SO(10)$ gauge symmetry, Higgs triplets and singlet components of 16 and $\overline{16}$, N_R and Y_1 , receive masses of the order of GUT scale, and they decouple from the spectrum. It is straightforward to extend the analysis to the case when N_R and Y_1 are lighter than M_{GUT} . The superpotential of the discussed model is of the form:

$$W = yH_{10}\phi_{16}\phi_{16} + hH_{10}\phi_{16}Y_{16} + MY_{16}Y_{\overline{16}}. \quad (4.2)$$

Let us note that two other messenger-matter couplings are also possible, $H_{10}Y_{16}Y_{16}$ and $H_{10}Y_{\overline{16}}Y_{\overline{16}}$, but, for the simplicity, we assume that they are suppressed with respect to those in (4.2). We applied the method presented in Sec. 3 to survey the issue of top-bottom-tau unification in the following way. First, the precise unification of Yukawa couplings was assumed: $y = y_t(t_{GUT}) = y_b(t_{GUT}) = y_\tau(t_{GUT})$. Then wave-function renormalization matrix $Z_{ij}(t)$ was derived from (3.7) and (3.8). As an example, we present an explicit expression for $Z_{H_u H_u}(t)$ at scale $t = \ln \mu / Q$ in the discussed model:

$$Z_{H_u H_u}(t) = 1 + \frac{6}{5}\varepsilon[3g_{GUT}^2 - 5(2h^2 - y^2)](t - t_{GUT}) + \frac{24}{25}\varepsilon^2[29g_{GUT}^2 + 35(2h^2 + y^2) - 25(2h^4 + 4h^2y^2 + y^4)](t - t_{GUT})^2 + \dots \quad (4.3)$$

As mentioned in Sec. 3, $Z_{ij}(t)$ can be expressed in terms of small parameter ε and GUT parameters of the model: t_{GUT} , y and h . Finally, the values of physical Yukawa couplings $\tilde{y}_{t,b,\tau}(t)$ at the messenger scale M were obtained from (3.6). In the next step, derived values of Yukawa couplings

$\tilde{y}_{t,b,\tau}(t)$ were appropriately implemented in the code of `SuSpect` [33]. Moreover, we computed 1- and 2-loop soft terms generated by (4.2), and also implemented them into the code of `SuSpect`. Then we used these numerical routines to derive physical spectrum at the EWSB scale and compared obtained values of $\tilde{y}_{t,b,\tau}(t)$ at the EWSB scale t_{EWSB} with the values obtained from t , b and τ pole masses. An example of the running of $\tilde{y}_{t,b,\tau}(t)$ is shown on Fig. 2. We scanned over the following range of parameters: $8 < t_M = \log_{10} M < 14$, $0.6 < y < 0.9$ and $0 < h < 1.2$ and checked if the simplest low-energy constraints are satisfied i.e.: (a) whether lightest neutral Higgs mass is about 125 GeV, (b) gluino \tilde{g} and 1st and 2nd generation squarks $\tilde{q}_{1,2}$ masses are bigger than 1.8 TeV and, finally, (c) whether EWSB vacuum is stable. We obtained the following results. When $\tilde{y}_t(t_{EWSB})$ and $\tilde{y}_\tau(t_{EWSB})$ match corresponding values derived from pole masses of t and τ , then one can get about 20% discrepancy between value of $\tilde{y}_b(t_{EWSB})$ and the related value $y_b^{(0)}$ derived from pole mass of b if: (i) $\tan \beta$ is about 45 and (ii) one allows for tachyonic $\tilde{\tau}$. On the other hand, when $\tan \beta$ is smaller, about 20, then there are no tachyons in the spectrum, but there is also large mismatch, of the order of factor 2, between $\tilde{y}_b(t_{EWSB})$ and $y_b^{(0)}$. Hence in the discussed minimal $SO(10)$ model, it is not possible to reconcile precise top-bottom-tau unification with the low-energy constraints on the spectrum. To avoid instability of the potential related to $\tilde{\tau}$ tachyonic mass, one could, for example, extend spectrum or allow additional messenger couplings.

5. Summary

We have discussed the issue of Yukawa couplings renormalization in the presence of messenger-matter interactions, which occur in the EGMSB models. Such messenger-matter couplings not only generate 1- and 2-loop soft terms but can also lead to kinetic mixing between heavy and light fields. It turns out that this mixing results in small shifts of the values of Yukawa couplings at the messenger scale, what in turn influence their values at the EWSB scale. We presented a method based on RGE for wave-function renormalization and showed that it is a handy tool to analyse RG flow of Yukawa coupling, especially useful in the top-bottom approach.

We applied proposed method to the specific $SO(10)$ inspired GUT model and showed that in this scenario it is not possible to reconcile precise top-bottom-tau unification with non-tachyonic spectrum, stable EWSB vacuum and right values of Yukawa couplings at the EWSB scale. It turned out that phenomenology of the discussed model was spoiled by tachyonic $\tilde{\tau}$. One of the possible ways out is to extend spectrum of the model or allow for additional messenger-matter couplings. Although the discussed example is quite specific, the presented method is general and can be used for studying Yukawa couplings renormalization in other, maybe less rigid, GUT models as $SU(5)$, or other $SO(10)$ scenarios which allow for small deviations from precise top-bottom-tau unification.

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