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Constraints of the B_μ/μ solution due to the hidden sector renormalization

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Abstract: We investigate the validity of an idea that the B_μ problem is solvable via the renormalization effect in the strongly interacting hidden sector within the gauge mediated supersymmetry breaking scenario. Our analysis starts with a naive boundary condition, which is that the squared scalar masses experience $16\pi^2$ suppression. We use softsusy to get the low energy spectra of superparticles with the boundary condition at the scale ($\Lambda_{\rm CFT}$) where the hidden sector is integrated out. We visit the low energy spectra and return to $\Lambda_{\rm CFT}$ where the boundary conditions are given. We find that there is a sign problem, which seems to be generic.

Keywords: Supersymmetric Standard Model, Higgs Physics, Supersymmetry Phenomenology.

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Contents

1. Introduction

Some attractive outcomes of TeV scale SUSY are a solution of the gauge hierarchy problem [1], the existence of dark matter candidate with the R-parity [2], and the radiative breaking of the electroweak symmetry [3]. The TeV scale soft SUSY breaking parameters are derived from the source of SUSY breaking, presumably at the hidden sector, and the messenger sector which couples to both the observable and the hidden sectors. In the visible sector, the standard model (SM) degrees of freedom are accommodated; in the messenger sector, we introduce the carriers of the SUSY breaking information to the observable sector. Thus, the phenomenology of the minimal supersymmetric standard model (MSSM) depends on where the hidden sector breaks SUSY and how this information is transmitted to the visible sector. The most popular mediation scheme is the minimal supergravity (mSUGRA) because the gravity can couple to both the hidden and the observable sectors. In mSUGRA, the information is revealed as non-renormalizable operators suppressed by powers of the Planck mass $M_{\rm Pl}$ [4]. There is, however, the serious flavor changing neutral current (FCNC) problem in mSUGRA. The FCNC problem is improved if one introduces very small squark mass differences $\Delta\tilde{m}_i{}^2$ or very large squark masses $\tilde{m}_i{}^2$ [6]. Since TeV scale SUSY does not permit very large $\tilde{m}_i{}^2$, there should be unnatural conditions such as the universal soft breaking terms for the scalar masses. Another scenario is the anomaly mediation scenario (AMSB), where the FCNC amplitude can be suppressed by the scale of the distance between the hidden sector and the visible sector branes [7]. AMSB, however, has an intrinsic serious problem that the scalar partners of the lepton are tachyonic at the low energy scale. To remedy this problem, one has to introduce a baroque structure [8].

The gauge mediated SUSY breaking (GMSB) scenario was introduced as another alternative and seems to remain as the simplest solution of the SUSY FCNC problem [5]. Recently, the concept of the metastable vacua in the hidden sectors brought renaissance of the GMSB scenario [9]. Especially, this concept made the GMSB scheme accommodated

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in the string theory [10 − 12] so that the interest on GMSB has increased. In GMSB, all SUSY breaking dimensionful parameters can be obtained by gauge interactions except B_μ.

On the other hand, μ is very strange in the MSSM because it is a unique dimensionful parameter in the supersymmetric part of the MSSM. So, if we consider the unbroken SUSY at a considerably high scale, then it is natural to consider μ of that scale, for example the order of the Planck mass $M_{\rm Pl}$. In addition, if one considers the $\mathrm{U}(1)_{\rm PQ}$ symmetry at the electroweak scale, the μ term is forbidden. If a non-trivial Kalher potential is considered, it is known that the SUSY breaking sector and the Higgs fields can be coupled. Therefore, the existence of SUSY breaking can give a rise to a correct order for μ and B_μ terms [14]. However, this supergravity generation of μ assumes negligible tree level μ contribution in the superpotential. The plausible reason for forbidding the μ term in the superpotential is originated from symmetries such as $\mathrm{U}(1)_{\rm PQ}$ and/or $\mathrm{U}(1)_R$. In general, it is difficult to obtain the μ term in the GMSB scenario though it is easily implemented in mSUGRA [13, 14], because $\mathrm{U}(1)_{\rm PQ}$ cannot be broken by gauge interactions. Thus, it is required for $\mathrm{U}(1)_{\rm PQ}$ breaking terms to enter the superpotential. To generate μ and B_μ, we introduce the direct interaction between Higgs and messengers as follows

$$W_{H_1 H_2} = \xi_1 H_1 \psi_1 \bar{\psi}_2 + \xi_2 H_2 \bar{\psi}_1 \psi_2, \tag{1.1}$$

where ψ_1 and ψ_2 are the messenger fields carrying appropriate weak and hyper charges to couple to $H_{1,2}$ at the tree level. After integrating out rather massive messengers, we get the appropriate operators for μ and B_μ. There is, however, another problem in GMSB: if we use the superpotential in (1.1), μ and B_μ are generated at a same loop level so that it is hard to satisfy low energy phenomenology. Several studies on this topic [24] exist already. Recently, the role of the hidden sectors has been raised [15]. Using this idea, an alternative solution for the B_μ/μ problem is suggested in the GMSB setup [16, 17], and we will discuss how this works in section 2. In this work, we investigate what conditions are required for the B_μ/μ problem in the setup of refs. [16, 17]. We set the boundary condition for this study that the squared scalar masses are $16\pi^2$ suppressed compared to the gaugino mass squared. This is similar to the usual gaugino mediation [25] in the ratio of the gaugino mass and the scalar mass. However, the messenger scales are quite different in the two cases so that the mass spectra at low energy may be totally distinguishable. As a result, we find that the idea suggested in [16, 17] has a tachyonic sector at low energy. We pursue the study on the region where low energy spectra satisfy the experimental result. We trace back to the 'effective' messenger scale, where the boundary conditions, which contain the hidden sector RG effects, are given. We find that B_μ carries opposite sign to μ at the 'effective' messenger scale. If the visible sector running effects do not give a significant contribution, this relation holds to the scale where the operators for B_μ and μ are generated, and is not compatible with the original relation between μ and B_μ.

In section 2, we will briefly review on the mechanism and the menace of tachyonic stau in the low energy spectra. In section 3, we will obtain the low energy spectra and postulate the valid parameter region in the sense of the low energy spectra. In section 4, we will discuss on the consistency of this mechanism, and make a conclusion.

2. The basic scheme

One of the most important success of SUSY is that it can explain how the electroweak symmetry breaking occurs. This can be achieved by the stop loop. The potential for Higgs in MSSM is given as

$$\begin{aligned} V =& |\mu|^2(|H_1|^2 + |H_2|^2) + \frac{1}{8}(g_1^2 + g_2^2)(|H_1|^2 - |H_2|^2)^2 \\ &+ m_{H_1}^2|H_1|^2 + m_{H_2}^2|H_2|^2 - (B_\mu H_1 H_2 + c.c.). \end{aligned} \tag{2.1}$$

From this we can derive the condition for EWSB, using the hessian for the Higgs mass matrix at the origin

$$B_\mu^2 > (|\mu|^2 + m_{H_1}^2)(|\mu|^2 + m_{H_2}^2). \tag{2.2}$$

Moreover, we require that the Higgs potential is bounded from below. There is a possible dangerous direction in (2.1). Therefore, this implies

$$2B_\mu < 2|\mu|^2 + m_{H_1}^2 + m_{H_2}^2. \tag{2.3}$$

For the CP even Higgs fields the mass matrix is given as

$$\begin{pmatrix} h \\ H \end{pmatrix} = \sqrt{2} \begin{pmatrix} \cos\alpha & -\sin\alpha \\ \sin\alpha & \cos\alpha \end{pmatrix} \begin{pmatrix} ReH_2 - v_2 \\ ReH_1 - v_1 \end{pmatrix}. \tag{2.4}$$

With the quantum correction, the mass of the lightest Higgs field saturates this inequality

$$m_h^2 \lesssim \cos^2 2\beta M_Z^2 + \frac{3\alpha_2}{2\pi}\frac{m_t^4}{M_Z^2}\ln\frac{\tilde{m}_t{}^2}{M_Z^2}. \tag{2.5}$$

From this if $\tan\beta > 4$, we can consider the lightest CP even Higgs field as the SM Higgs field. From the result of the LEP we know that the lower bound of the SM Higgs mass is 114.4GeV [19].[1]

At the intermediate scale, which is between the electroweak (M_Z) and the messenger ($M_{\rm mess}$) scales, the MSSM couplings are not so large. Therefore, once B_μ and μ are generated, the ratio between B_μ and μ^2 does not suffer a significant change. This is undesirable at the electroweak scale. For successful electroweak symmetry breaking, we require that both of these are order of the gaugino masses,

$$B_\mu \sim \mu^2 \qquad \mu \sim m_{\frac{1}{2}}. \tag{2.6}$$

Let us consider how the hidden sector strong dynamics works toward the electroweak symmetry breaking. For a concrete discussion of the messenger effect toward the MSSM physics and the μ generation, we adopt the simple superpotential in (1.1). Integrating out the heavy messenger fields ($M_{\rm mess} \gg \Lambda_{\rm CFT} \gg M_Z$), we obtain non-renormalizable

[1] If the R-parity is broken, the lightest Higgs mass can be lower than the LEP bound [20].

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interaction terms between the hidden and the visible sector fields. In this way, let us consider the following operators relevant for the dimensionful parameters in the MSSM,

$$
\begin{aligned}
O_\phi : &\int d^4\theta\; c_\phi^q\; \frac{q^\dagger q}{M^2}\phi^\dagger\phi, \\
&\int d^4\theta\; c_\phi^s\; \frac{S^\dagger S}{M^2}\phi^\dagger\phi, \\
O_{B_\mu} : &\int d^4\theta\; c_{B_\mu}^q\; \frac{q^\dagger q}{M^2}H_1H_2 + h.c., \\
&\int d^4\theta\; c_{B_\mu}^s\; \frac{S^\dagger S}{M^2}H_1H_2 + h.c., \\
O_\lambda : &\int d^4\theta\; c_\lambda^s\; \frac{S}{M}W^{a\alpha}W_\alpha^a + h.c, \\
O_A : &\int d^4\theta\; c_A^s\; \frac{S}{M}\phi^\dagger\phi + h.c., \\
O_\mu : &\int d^4\theta\; c_\mu^s\; \frac{S^\dagger}{M}H_1H_2 + h.c.,
\end{aligned}
\tag{2.7}
$$

where $H_{1,2}$ and ϕ are the MSSM fields and the rest are the intermediate scale fields which constitute the ingredients for SUSY breaking. Here, S, q and W^α are spurion, quark, and gaugino fields, respectively, in the intermediate scale, and cs are the couplings. Refs. [16, 17] consider a hidden conformal sector at the intermediate scale, which is guaranteed by Seiberg's duality [26]. In Seiberg's conformal window, the electric and magnetic descriptions are the same. At this window the gauge coupling is asymptotically free, and hence it is meaningless to use the perturbation method in the low energy limit. Thus, the theory naturally has a low energy cutoff, which is usually represented as a mass parameter Λ. In QCD, for example, it is denoted as $\Lambda_{\rm QCD}$. In this vein, we will define the intermediate mass scale as $\Lambda_{\rm CFT}$. Next, integrating out the fields at $\Lambda_{\rm CFT}$, the effective operators for the soft terms are obtained. That is to say, the renormalization effect below the scale $\Lambda_{\rm CFT}$ is nothing but that of MSSM with the boundary conditions fixed at $\Lambda_{\rm CFT}$. The MSSM RG has been widely studied, in the literature such as [18].

From (2.7) we note that the soft scalar mass has the same property as B_μ. On the other hand, the trilinear coupling A behaves the same as the gaugino mass or μ. From (1.1) we also note that μ and the gaugino mass are generated at one loop level. So the relative size between the gaugino mass and μ can be easily fitted to the phenomenological expectation. However, not only μ but also B_μ are generated at one loop level. It turns out that the ratio between B_μ and μ at the messenger scale is too large to fit the phenomenological requirement. The renormalization in the SUSY gauge theory is revealed as the wave function renormalization. Considering the 1PI renormalization for $S^\dagger S$ in addition to the wave function renormalization below $\Lambda_{\rm CFT}$, then the effective operators in (2.7) should be substituted by

$$
\begin{aligned}
O_\phi : & \int d^4\theta \left(\frac{\Lambda_{\rm CFT}}{M_{\rm mess}}\right)^{\alpha_q} Z_q^{-1}\, c_\phi^q\, \frac{q^\dagger q}{M^2}\phi^\dagger\phi, \\
& \int d^4\theta \left(\frac{\Lambda_{\rm CFT}}{M_{\rm mess}}\right)^{\alpha_s} Z_s^{-1}\, c_\phi^s\, \frac{S^\dagger S}{M^2}\phi^\dagger\phi, \\
O_{B_\mu} : & \int d^4\theta \left(\frac{\Lambda_{\rm CFT}}{M_{\rm mess}}\right)^{\alpha_q} Z_q^{-1}\, c_{B_\mu}^q\, \frac{q^\dagger q}{M^2}H_1H_2 + h.c., \\
& \int d^4\theta \left(\frac{\Lambda_{\rm CFT}}{M_{\rm mess}}\right)^{\alpha_s} Z_s^{-1}\, c_{B_\mu}^s\, \frac{S^\dagger S}{M^2}H_1H_2 + h.c., \\
O_\lambda : & \int d^4\theta Z_s^{-1/2}\, c_\lambda^s\, \frac{S}{M}W^{a\alpha}W_\alpha^a + h.c, \\
O_A : & \int d^4\theta Z_s^{-1/2}\, c_A^s\, \frac{S}{M}\phi^\dagger\phi + h.c., \\
O_\mu : & \int d^4\theta Z_s^{-1/2}\, c_\mu^s\, \frac{S^\dagger}{M}H_1H_2 + h.c.,
\end{aligned}
\tag{2.8}
$$

where α_q and α_s are anomalous dimensions and the wave function renormalization factors Z_s and Z_q are defined as

$$
Z_{s,q} = \left(\frac{\Lambda_{\rm CFT}}{M_{\rm mess}}\right)^{3R(S,q)-2}. \tag{2.9}
$$

There exist some subtle points in this mechanism which are caused by the mixing between operators given in (2.8). There can be mixing between the quadratic operator and the linear operators as well as mixing between the quadratic operators. Each mixing is induced by direct interactions between matter and messenger fields, and the strongly interacting hidden sector respectively. The latter appears as the anomalous dimension, although we can not get the exact value. Here we assume that we obtain the hidden sector effect as

$$
\text{suppression factor} = \left(\frac{\Lambda_{\rm CFT}}{M_{\rm mess}}\right)^{\alpha}, \tag{2.10}
$$

where α is the smallest eigenvalue of matrix for the anomalous dimensions, which reflects the mixing between the quadratic operators. And the other affects the boundary conditions too. In ref. [27], the effects on the soft parameters under the presence of such a superpotential (1.1) are shown. Now we will turn to very subtle points of mixing between operators. In (2.8), we do not consider the operator mixing, however, within the sense of effective field theory there is mixing, which affects on the boundary condition which will be shown in the next section. Generally, the soft parameters of mass dimension two are some combinations of the quadratic operators and the linear operators. We, therefore, provide the terms which appear as the boundary conditions. This is well explained in ref. [16], and we follow its description. For the ordinary scalar fields, it is $c_\phi^s - |c_A^s|^2$. For the Higgs field, there is another contribution from μ so that it is $c_{H_{1,2}}^s - |c_A^s|^2 - |c_\mu^s|^2$. Finally, it is $c_{B_\mu}^s - c_\mu^s(c_{A_{H_1}}^s + c_{A_{H_2}}^s)$ for B_μ. These all experience the hidden sector RG effect, and it should be realized in the boundary conditions. Via these effects at the scale of CFT breaking, the ratio of B_μ and μ can be made to satisfy the relation (2.6). Then we can see that B_μ suffers the renormalization effects through the strongly interacting CFT sector. However, there is an effect which we should not ignore. It is that the squared masses of scalar also suffer the same kind of renormalization as B_μ. In the operator sense, scalar masses and gaugino masses are generated properly at the the messenger scale. However, when we reach at the scale $\Lambda_{\rm CFT}$, the scalar masses would suffer $16\pi^2$ suppression by the same mechanism that reduces B_μ. On the other hand, the trilinear A terms do not undergo such a suppression.

Keeping these in mind, let us consider the mixing matrices for the $\tilde{m}_t^2, \tilde{m}_b^2$ and $\tilde{m}_\tau^2$:

$$\begin{pmatrix} \tilde{m}_{tL}^2 & m_t(A_t - \mu\cot\beta) \\ m_t(A_t - \mu\cot\beta) & \tilde{m}_{tR}^2 \end{pmatrix}, \tag{2.11}$$

$$\begin{pmatrix} \tilde{m}_{bL}^2 & m_b(A_b - \mu\tan\beta) \\ m_b(A_b - \mu\tan\beta) & \tilde{m}_{bR}^2 \end{pmatrix}, \tag{2.12}$$

$$\begin{pmatrix} \tilde{m}_{\tau L}^2 & m_\tau(A_\tau - \mu\tan\beta) \\ m_\tau(A_\tau - \mu\tan\beta) & \tilde{m}_{\tau R}^2 \end{pmatrix} \tag{2.13}$$

with

$$\begin{aligned}
\tilde{m}_{tL}^2 &= \tilde{m}_Q^2 + m_t^2 + \frac{1}{6}(4M_W^2 - M_Z^2)\cos 2\beta, \\
\tilde{m}_{tR}^2 &= \tilde{m}_U^2 + m_t^2 - \frac{2}{3}(M_W^2 - M_Z^2)\cos 2\beta, \\
\tilde{m}_{bL}^2 &= \tilde{m}_Q^2 + m_b^2 - \frac{1}{6}(2M_W^2 + M_Z^2)\cos 2\beta, \\
\tilde{m}_{bR}^2 &= \tilde{m}_D^2 + m_b^2 + \frac{1}{3}(M_W^2 - M_Z^2)\cos 2\beta, \\
\tilde{m}_{\tau L}^2 &= \tilde{m}_L^2 + m_\tau^2 - \frac{1}{2}(2M_W^2 - M_Z^2)\cos 2\beta, \\
\tilde{m}_{\tau R}^2 &= \tilde{m}_E^2 + m_\tau^2 + (M_W^2 - M_Z^2)\cos 2\beta.
\end{aligned}$$

Now, we can obtain the masses at the electroweak scale. As denoted above, below $\Lambda_{\rm CFT}$ the renormalization equations are that of the MSSM. Thus, the renormalization property of each diemensionful parameter is given as

$$\begin{aligned}
\frac{dM_i}{dt} &= b_i\alpha_i M_i. \\
\frac{dA_U}{dt} &= \frac{16}{3}\alpha_3 M_3 + 3\alpha_2 M_2 + \frac{13}{15}\alpha_1 M_1 + 6Y_U A_U + Y_D A_D, \\
\frac{dA_D}{dt} &= \frac{16}{3}\alpha_3 M_3 + 3\alpha_2 M_2 + \frac{7}{15}\alpha_1 M_1 + 6Y_D A_D + Y_U A_U + Y_L A_L, \\
\frac{dA_L}{dt} &= 3\alpha_2 M_2 + \frac{9}{5}\alpha_1 M_1 + 3Y_D A_D + 4Y_L A_L, \\
\frac{dB}{dt} &= 3\alpha_2 M_2 + \frac{3}{5}\alpha_1 M_1 + 3Y_U A_U + 3Y_D A_D + Y_L A_L, \\
\frac{d\tilde{m}_Q^2}{dt} &= -\left[\left(\frac{16}{3}\alpha_3 M_3^2 + 3\alpha_2 M_2^2 + \frac{1}{15}\alpha_1 M_1^2\right) - Y_U(\tilde{m}_Q^2 + \tilde{m}_U^2 + m_{H_2}^2 + A_U^2)\right. \\
&\qquad\qquad\qquad\qquad\qquad \left. - Y_D(\tilde{m}_Q^2 + \tilde{m}_D^2 + m_{H_1}^2 + A_D^2)\right],
\end{aligned}$$

$$
\begin{aligned}
\frac{d\tilde{m}_U^2}{dt} &= -\left[\left(\frac{16}{3}\alpha_3 M_3^2 + \frac{16}{15}\alpha_1 M_1^2\right) - 2Y_U(\tilde{m}_Q^2 + \tilde{m}_U^2 + m_{H_2}^2 + A_U^2)\right], \\
\frac{d\tilde{m}_D^2}{dt} &= -\left[\left(\frac{16}{3}\alpha_3 M_3^2 + \frac{4}{15}\alpha_1 M_1^2\right) - 2Y_D(\tilde{m}_Q^2 + \tilde{m}_D^2 + m_{H_1}^2 + A_D^2)\right], \\
\frac{d\tilde{m}_L^2}{dt} &= -\left[3\left(\alpha_2 M_2^2 + \frac{1}{5}\alpha_1 M_1^2\right) - Y_L(\tilde{m}_L^2 + \tilde{m}_E^2 + m_{H_1}^2 + A_L^2)\right], \\
\frac{d\tilde{m}_E^2}{dt} &= -\left[\left(\frac{12}{5}\alpha_1 M_1^2\right) - 2Y_L(\tilde{m}_L^2 + \tilde{m}_E^2 + m_{H_1}^2 + A_L^2)\right], \\
\frac{d\mu^2}{dt} &= -\mu^2\left[3\left(\alpha_2 + \frac{1}{5}\alpha_1) - (3Y_U + 3Y_D + Y_L\right)\right], \\
\frac{dm_{H_1}^2}{dt} &= -\left[3\left(\alpha_2 M_2^2 + \frac{1}{5}\alpha_1 M_1^2\right) - 3Y_D(\tilde{m}_Q^2 + \tilde{m}_D^2 + m_{H_1}^2 + A_D^2)\right. \\
&\qquad\qquad \left. - Y_L(\tilde{m}_L^2 + \tilde{m}_E^2 + m_{H_1}^2 + A_L^2)\right], \\
\frac{dm_{H_2}^2}{dt} &= -\left[3\left(\alpha_2 M_2^2 + \frac{1}{5}\alpha_1 M_1^2\right) - 3Y_U(\tilde{m}_Q^2 + \tilde{m}_U^2 + m_{H_2}^2 + A_U^2)\right],
\end{aligned}
$$

where $\alpha_i = \frac{g_i^2}{4\pi}$ and $t = \ln\frac{Q}{\mu_r}$. Here, we observe that the right-handed stau can be dangerous. Because it carries only U(1) hypercharge it receives the contribution from bino mass. Thus, the gaugino mass contribution cannot be large via the MSSM renormalization. On the other hand, A_τ can not be neglected in general at the $\Lambda_{\rm CFT}$ scale. From the mixing matrix for the stau (2.13), we can expect that stau becomes tachyonic in a certain range of parameters. To show it explicitly, we use `softsusy` for the scalar masses running [21]. Even if the trilinear couplings are made zero at the messenger scale, the scalar masses are not free from the menace of the tachyonic states, because there are also contributions from μ as can be seen from (2.11), (2.12), (2.13). We should be careful about this effect too.

3. Numerical analysis: visiting the low energy spectra

To check the discussion in the previous section, we use `softsusy`. In `softsusy`, the input parameters are $\tan\beta$, $M_{\rm mess}$, number of messengers, gravity contribution and $\Lambda = \frac{F}{M_{\rm mess}}$. We will use $\Lambda_{\rm CFT}$ as an 'effective' messenger scale. At this scale, the squared scalar masses suffer $16\pi^2$ suppression and the trilinear term can be obtained. It is a good approximation to use the basic setup provided by `softsusy` except the suppression of scalar masses and non-zero trilinear coupling. The others such as μ and B_μ are obtained in the range where the low energy phenomenology allows. In the previous section, we discuss that the mixing can exist, and we will apply the argument to the boundary conditions. Since we consider the minimal case, i.e. all MSSM fields do not have Yukawa interaction with the messenger fields except Higgs, there is no significant contribution to A_ϕ. On the other hand, the soft masses of the Higgs obtain these contributions of $\delta m_{H_{1,2}}^2 \sim -\mu^2$. In addition, there is a contribution of order of $-\mu$ to the trilinear coupling $A_{U,D,L}$ respectively. Therefore, we set the universal trilinear coupling $A = -\mu$. The relation between trilinear couplings is as follows

$$A_{U,D,L} = Y_{U,D,L}A,$$
$$A = A_{H_{1,2}},$$

where $Y_{U,D,L}$ are Yukawa matrices.

We will do our calculations as follows:

- Set the scale, where the hidden sectors are integrated out, as $\Lambda_{\rm CFT}$.
- Set the messenger scale as 10^{14}GeV.
- At $\Lambda_{\rm CFT}$, the scalar masses are suppressed by $16\pi^2$.
- Set $\Lambda_{\rm CFT}$, i.e. the 'effective' messenger scale, as 10^8GeV.
- The sign of μ is positive.
- Set the gravity contribution as zero.
- Set $m_t = 170.9$GeV.
- The trilinear couplings of Higgs are generated as $\delta A_{H_{1,2}} \sim -\mu$. By them, the universal trilinear coupling satisfies $A \sim -\mu$.
- The soft Higgs masses receive the contribution of $\delta m^2_{H_{1,2}} \sim -\mu^2$.
- Set $\tan\beta$ and $\frac{F}{M_{\rm mess}}$ as free parameters.
- Scan $\tan\beta$ from 4 to 50 and $\frac{F}{M_{\rm mess}}$ from 5.0×10^4GeV to 2.0×10^5GeV for the case of 1 messenger.
- No consideration about other low energy constraints such as $B_s\gamma$.

In figure (1) the blue section represents the tachyonic region and the green part does the stau direct search bound. Since there can be theoretical errors in calculating the mass spectra with software packages such as `FeynHiggs` and `softsusy` [22], we allow -3GeV difference. The yellow region represents the section where the lightest Higgs mass is between 111.4GeV and 114.4GeV. On the other hands, the red region can be said to be definitely ruled out by the direct Higgs search bound of the LEP experiment.

In this figure, we can see that there is a tachyonic region at the large $\tan\beta$. This feature appears similarly in some parameter space in the gaugino mediation. This can be easily understood when we consider the mass matrix of the stau in (2.13). To use the hidden sector strong RG effects as a solution of B_μ problem, the squared scalar masses are not free from $16\pi^2$ suppression. Therefore, the diagonal parts of the stau mass matrix are rather small compared to the off diagonal parts. Here we look more carefully the off diagonal parts. The off diagonal parts are composed with tau mass, the trilinear coupling A and $\mu\tan\beta$. In `softsusy`, μ and B_μ are fitted by the proper EWSB; therefore, we do not have to worry about this. We investigate the possibility that the parameter space can

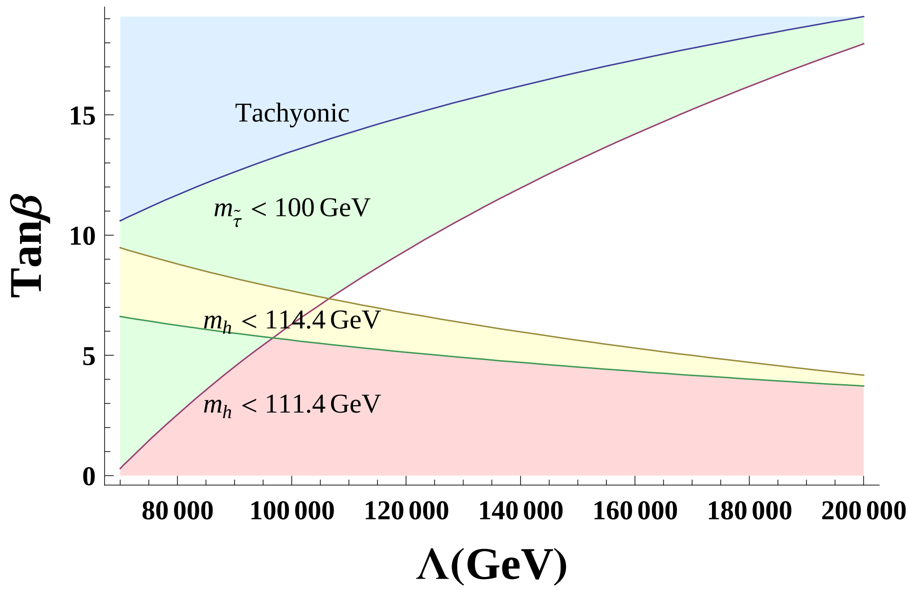

Figure 1: Plot of forbidden region in $16\pi^2$ suppression case, with $A = -\mu$, $\Lambda_{\rm CFT} = 10^8$GeV and the number of messengers $= 1$. The yellow and the red parts are mass bound for the lightest Higgs.

be enlarged. Let us consider the case that the 'effective' messenger scale to be different from 10^8GeV. Varying $\Lambda_{\rm CFT}$ from 10^6GeV to 10^{10}GeV, we find that the pattern does not change significantly. As denoted above we have the messenger scale as 10^{14}GeV; thus, the 'effective' messenger scale can not be larger than it.

4. Numerical analysis: more on the valid region

The mechanism which we investigate, also provides generating μ on a theoretical base. Here we will check whether the relation in (2.6) can be satisfied at the scale where $16\pi^2$ suppression does appear. We will stay in the region where the low energy spectra appear to be valid. First of all, we should keep in mind that in `softsusy` μ and B_μ are fitted by the requirement of the proper EWSB. The method to consider the hidden sector RG running effects has a unique property. Since the hidden sector RG effects which make the B_μ comparable to μ^2, affect the operators which are universally proportional to $SS^\dagger$; therefore, the squared scalar masses suffer such a suppression. These effects are revealed in the boundary conditions which we have chosen at the 'effective' messenger scale. Let us return to the start point of our analysis. We consider a direct interaction between messengers and Higgs like (1.1). With this superpotential, we derive these relations:

$$\begin{aligned}
\mu &= \frac{\xi_1\xi_2}{16\pi^2}\Lambda f(\lambda_1/\lambda_2)\left[1+\mathcal{O}\left(\frac{F^2}{M_{\rm mess}^4}\right)\right] \\
B_\mu &= \frac{\xi_1\xi_2}{16\pi^2}\Lambda^2 f(\lambda_1/\lambda_2)\left[1+\mathcal{O}\left(\frac{F^2}{M_{\rm mess}^4}\right)\right] \qquad (4.1)\\
B_\mu &= \Lambda\mu,
\end{aligned}$$

where $\lambda_{1,2}$ are coupling constants between the messenger and the goldstino supermultiplet. f is a function appearing after we integrate out the messenger fields. To satisfy phenomenological low energy requirements, we introduce hidden sector RG effects. As a result we get the suppression factor appear in (2.8). We set the boundary conditions to represent such factors and we get the valid region. Now let us refer to $\tilde{B_\mu}$[2] as the postulated value by `softsusy` to satisfy the low energy requirements and B'_μ[3] as one obtained by the trace back RG of $\tilde{B_\mu}$ respectively. The ratio $\frac{B'_\mu}{B_\mu}$, which affects the squared scalar masses, is set as boundary conditions at the 'effective' messenger scale. Then we expect that the region which passes our consistency test, would appear as a band in the valid region of section 3. The band should be under the control of the value $\xi_1\xi_2$. However, we should not miss a point that μ is also dependant on the value of $\xi_1\xi_2$. If we use the third relation of (4.1), we can eliminate this $\xi_1\xi_2$ dependance.

Our strategy is very simple. Once we get the $\tilde{\mu}$ and $\tilde{B_\mu}$ at the electroweak scale. We will follow the MSSM RG flow to the 'effective' messenger scale in the valid parameter space so that we can get μ' and B'_μ. Moreover, it is natural to identify μ with μ'. Then we will check whether the factor we get by the trace back RG is the same as we set as boundary conditions and the evaluation of the μ is consistent with (4.1). Now let us turn to the figure (2). In this case, we provide the ratio δ between $\mu \times \Lambda$ and $B'_\mu \times$ suppression factor obtained by the `softsusy`

$$\rho = \frac{B'_\mu \times (\text{suppression factor})}{\mu \times \Lambda}. \tag{4.2}$$

Here we see that the valid region in figure (1) has a negative ρ in figure (2). This means that B'_μ have an opposite sign to μ. Here we want to look into (4.1) carefully. The third relation says that if we fix μ to be positive real, then the sign of B_μ is dependent on the sign of Λ. That is, if Λ is positive, then B'_μ as well as B_μ should be positive, since the suppression factor does not change the sign of B'_μ. If the suppression factor change the sign, it will affect the sign of the squared scalar masses. Of course, there are studies on this case, i.e. the negative squared scalar masses by allowing a large mixing [28], but we will leave this topic to the further study.[4] This sign problem might be accidental at $\Lambda_{\rm CFT}$, so we are not sure whether this can be really problematic. The visible sector contribution is suppressed as much as B_μ above $\Lambda_{\rm CFT}$. Therefore, the dominant contribution comes from the hidden sector, and it is dependent on the sign of B'_μ. To see explicitly, let us check this. Terms, which run for RG of B_μ in the visible sector, are the linear terms shown in (2.8); they experience the same hidden sector RG effect as μ. We can divide the RG property of

[2]From now on, the tilded represent the value which obtained by `softsusy` to satisfy the requirements of the low energy

[3]From now on, the primed are obtained by the trace back RG. For example, μ' is the result of the trace back RG of $\tilde{\mu}$.

[4]We run the program for the negative suppression, and we find that in this case this mechanism pass the our consistency test. We, however, need more clarification.

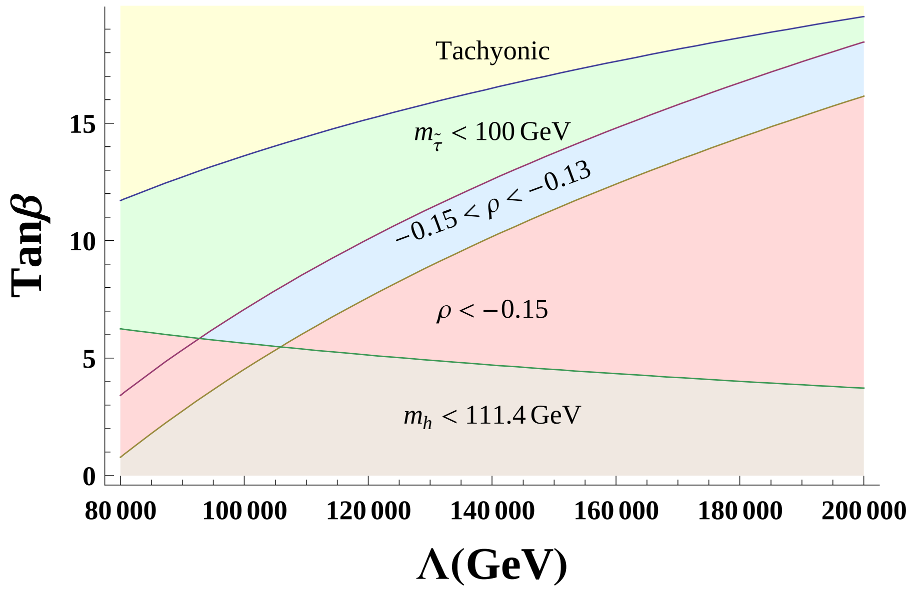

Figure 2: The set up is the same as the previous one except that the yellow part is excluded and we provide the ratio ρ between the theoretical prediction and the result of the trace-back RG. The blue region is $-0.15 < \rho < -0.13$, and the red is $\rho < -0.15$.

B'_μ into two part, the visible sector contribution and the hidden sector contribution:

$$\begin{aligned}\delta B'_\mu &= \delta(\text{visible part}) + \delta(\text{hidden contribution}) \\ &= \delta(B \times \mu \times \text{suppression factor}) \\ &= \delta(B \times \mu) \times (\text{suppression factor}) + B \times \mu \times \delta(\text{suppression factor}),\end{aligned}$$

where $B = \frac{B_\mu}{\mu}$, and the suppression factor is $\left(\frac{\mu_R}{M_{\text{mess}}}\right)^\alpha$ above Λ_{CFT}. As denoted above, the hidden sector contribution depends on the sign of B'_μ; thus, we should check whether the visible sector contribution can flip the sign. Here we assume that the visible sector enjoys the ordinary MSSM RG. It is sufficient to check the RG of B'_μ with the ordinary MSSM RG equation to the real messenger scale. Let us turn off the hidden sector contribution for a while, i.e. $\delta(\text{hidden sector}) = 0$. Then we run the visible sector RG of B'_μ to the messenger scale. As a result, we find that B'_μ does not get positive with the visible sector contribution only (See figure 3). We may doubt whether this sign problem is originated from the rather small effective messenger scale. In figure (4), we see that this problem is generic; especially, $\frac{B'_\mu}{\mu^2}$ becomes smaller as $\tan\beta$ increases. On the other hand, this result is highly dependent on the boundary condition, especially on the trilinear coupling A. If we accept the minimal Yukawa coupling of the messenger fields, the dominant contribution to A is derived by the Higgs-messenger Yukawa coupling, and this is the option we choose. If we choose A differently, for example to be μ, then the result change seriously though this is not the ordinary case.

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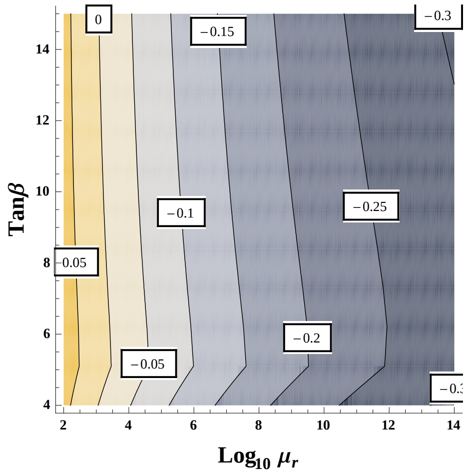

Figure 3: The RG of B'_μ with $\Lambda = 1.5 \times 10^5 GeV$ within the visible sector only. Here we see that the visible sector cannot make $\frac{B'_\mu}{\mu^2}$ positive.

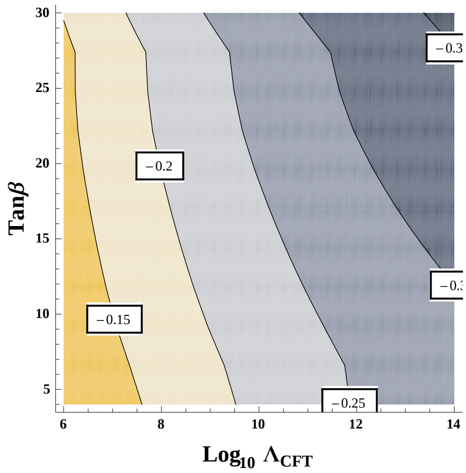

Figure 4: The effect of varying $\Lambda_{\rm CFT}$ scale with $\Lambda = 1.5 \times 10^5 GeV$. Here we can see that the sign problem is generic in this mechanism. The contours represent the value of $\frac{B'_\mu}{\mu^2}$.

5. Conclusion

In the present study, we investigate the low energy spectra of the B_μ/μ solution provided by the strong hidden sector. Via the strong hidden sector RG effects, the squared scalar masses suffer $16\pi^2$ suppression. As a result, diagonal parts of the mass matrices of the scalar can be relatively small compared with other cases. Especially stau might give a constraint in the parameters space. Using `softsusy`, we observe that there exists tachyonic sector for the large $\tan\beta$. In the region which appears to be valid in the low energy spectra test, we trace back to the 'effective' messenger scale along the MSSM RG flow. Then we compare the factor which we obtain by the trace back RG and the factor which we have chosen as the boundary conditions at the 'effective' messenger scale. During this study, we find that there is a sign problem, which seems generic in this mechanism.

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