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Light Gauginos and Conformal Sequestering

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Abstract

In a wide class of direct and semi-direct gauge mediation models, it has been observed that the gaugino masses vanish at leading order. It implies that there is a hierarchy between the gaugino and sfermion masses, invoking a fine-tuning problem in the Higgs sector via radiative corrections. In this paper, we explore the possibility of solving this anomalously light gaugino problem exploiting strong conformal dynamics in the hidden sector. With a mild assumption on the anomalous dimensions of the hidden sector operators, we show that the next to leading order contributions to the gaugino masses can naturally be in the same order as the sfermion masses. μ/B_μ problem is also discussed.

1 Introduction

Gauge mediated supersymmetry breaking [1, 2, 3, 4] is a highly predictive and attractive way of transmitting supersymmetry breaking of the hidden sector to the SSM (See [5, 6, 7] for reviews). Also, given the fact that gauge interactions do not distinguish flavors, phenomenologically dangerous flavor changing neutral currents are naturally suppressed.

One drawback of gauge mediation is that gauginos tend to be light compared to sfermions. First of all, gauge interactions themselves do not break R-symmetry, so R-symmetry must be broken in the hidden sector for gauginos to obtain non-vanishing masses. But as was studied in [8], the absence of supersymmetry preserving vacuum for a theory with generic superpotential requires R-symmetry be spontaneously broken, thus introducing a phenomenologically unfavorable R-axion. This problem can be avoided using recently proposed metastable supersymmetry breaking vacua [9]. However, there is another problem: broken R-symmetry itself does not guarantee sufficiently large gaugino masses. Actually, in a wide class of direct gauge mediation models (See e.g. [10, 11, 12, 13, 14]) and also in any semi-direct gauge mediation models [15, 16, 17], the leading order contribution to gaugino masses vanishes irrespective of whether R-symmetry is broken or not¹.

Anomalously light gauginos are problematic because relatively heavy sfermions induce a large correction to the Higgs mass, reintroducing the hierarchy problem. One possible way out may be to take the messenger scale to be very close to the supersymmetry breaking scale so that the subleading corrections are to be in the same size as the leading contribution. However, as was studied in [19], such a model is severely constrained by the recent Tevatron bound on the sparticle masses and the mass bound on a light gravitino.

Recently, Komargodski and Shih shed light on the origin of the light gauginos. In [20], they related the vanishing gaugino masses at leading order and global structure of the vacua in renormalizable theories, and showed, based on the study of generalized O’Raifeartaigh models, that the pseudomoduli space must have a tachyonic direction somewhere to generate sizable gaugino masses. This analysis opens up a new possibility to avoid the anomalously light gaugino problem. Namely, the leading order gaugino mass generally does not vanish if supersymmetry is broken in uplifted metastable vacua. This idea was initially employed in [21] and further discussed in [22, 23]².

In this paper, we propose an alternative solution to this anomalously light gaugino mass problem. As was studied in [25], due to the gaugino screening effect, it is hard to generate

¹For F -term breaking models with calculable messenger sector, one can show that there is an upper bound for the ratio of gaugino masses to sfermion masses [18].

²For recent studies for generating leading order gaugino masses in semi-direct gauge mediation, see [24].

the leading order gaugino mass at one-loop level even if we impose a non-canonical Kähler potential of a messenger. Thus, we instead reduce the sfermion masses relative to the gaugino masses by exploiting conformal sequestering [26, 27]³: With an appropriate choice of a hidden sector which is assumed to approach a strongly coupled fixed point below the messenger scale, sfermion masses can be suppressed relative to the gaugino masses by strong hidden sector renormalization group effects. Thus, one can make the sfermion mass parameters be lower or in the same order as the gaugino masses at the end of conformal sequestering. Then, the sfermion masses are driven to be in the same order as gaugino masses by standard model renormalization group effects.

In this work, we only discuss how conformal sequestering works for a messenger model with vanishing leading order gaugino masses and do not discuss the details of the hidden sector. But embedding our scenario into direct and semi-direct gauge mediations where the gaugino masses are zero at leading order would be possible and an interesting future direction.

The rest of the paper is organized as follows. In section 2, we review light gaugino mass problem and gaugino screening. Then, we comment on UV sensitivity of a non-renormalizable model and discuss a contribution of a heavy messenger which is either coupled to a light messenger or not. In section 3, we discuss the effects of conformal sequestering on sfermion and gaugino masses as well as μ , B_μ terms. Although μ is required to be fine-tuned in a technically natural way, we find that our scenario can be phenomenologically viable with all the dimensionless parameters other than Higgs-hidden sector coupling be of order one. In the Appendix, we present a model of Higgs-messenger coupling which generates μ and B_μ at one-loop with a messenger sector exhibiting vanishing leading order gaugino masses.

2 Anomalously Light Gaugino Problem

In this section, we review some known facts on gaugino masses at leading order in supersymmetry breaking scale. If the SUSY breaking scale is much smaller than messenger masses, one can reliably use the technique of analytic continuation into superspace [34, 25].

2.1 Leading order gaugino mass

We first review anomalously small gaugino mass problem which have been observed quite frequently in direct gauge mediation models. Suppose a messenger sector having the following

³For applications to gauge mediation, see, e.g., [28, 29, 30, 31, 32, 33].

general superpotential interaction with supersymmetry breaking field $\langle X \rangle = M + \theta^2 F$,

$$W = \sum_{ab} \mathcal{M}(X)_{ab} \phi^a \tilde{\phi}^b. \quad (2.1)$$

where the messenger mass matrix $\mathcal{M}_{ab}(X)$ is a holomorphic function of X . In this case, the gaugino masses are generated by integrating out the messengers $\phi^a, \tilde{\phi}^b$. The formula is given [34, 21] by

$$m_\lambda = -\frac{g_{\text{SM}}^2}{16\pi^2} F \frac{\partial}{\partial X} \log \det \mathcal{M}(X). \quad (2.2)$$

As was studied in [20], if the fermion mass matrix $\mathcal{M}(X)$ includes a zero eigenvalue at some point in the pseudomoduli space spanned by X , then the corresponding bosonic mode becomes tachyonic in a region including the point where the fermionic zero mode appears. In this case, the determinant of $\mathcal{M}(X)$ is a function of X . On the other hand, if the eigenvalues of $\mathcal{M}(X)$ are non-vanishing, the bosonic modes of ϕ^a are stable everywhere⁴ in the pseudomoduli space and $\det \mathcal{M}(X)$ is constant, which yields vanishing leading order gaugino masses. Since the sfermion masses are generally generated at this order, the gauginos are light relative to the sfermions if the pseudomoduli space is stable everywhere.

2.2 Gaugino screening

Now we consider a more general setup in which the Kähler potential of messengers is non-canonical: it has a dependence on the supersymmetry breaking field X . Let us derive the gaugino mass formula utilizing analytic continuation into superspace [25]. Suppose the theory has N pair of messengers $\phi^a, \tilde{\phi}^a$ ($a = 1, \dots, N$) which are fundamentals and anti-fundamentals of the standard model gauge interactions, respectively. A generic Kähler potential we consider is

$$K = \sum_a Z_a(X, X^\dagger) (\phi^{a\dagger} e^{V_{\text{SM}}^{(\phi)}} \phi^a + \tilde{\phi}^{a\dagger} e^{V_{\text{SM}}^{(\tilde{\phi})}} \tilde{\phi}^a), \quad (2.3)$$

where $Z_a(X, X^\dagger)$ are some real functions of X, X^\dagger . Finally, the superpotential is given by (2.1).

One can extract the gaugino masses generated by integrating out the messenger fields from the wave function renormalization gauge chiral superfield. One should, however, use the physical gauge coupling R rather than the holomorphic one, since the holomorphic coupling is not invariant under field rescaling [25]. As pointed out in [25], contributions from messenger interactions to the gaugino masses are suppressed by additional loop factors. Thus, a non-canonical Kähler potential cannot contribute to the leading order gaugino mass. To see

⁴The presence of a supersymmetric vacuum at large X should be a small effect.

this, one may write down the physical coupling below the messenger scale. For the sake of simplicity, we assume the fermion mass matrix of the messengers is constant: $\mathcal{M}(X) = m$, so $W = m\phi\tilde{\phi}$. The physical mass is defined using wavefunction renormalization Z_M of the messenger at the scale,

$$\mu_m^2 = \frac{|m|^2}{Z_M(\mu_m)^2}.$$

Below this scale, the physical coupling is given by

$$R(\mu) = R'(\mu_0) + \frac{b}{16\pi^2} \log \frac{\mu^2}{\mu_0^2} + \frac{1}{16\pi^2} \log \frac{|m|^2}{\mu_0^2 Z'_M(\mu_0)^2} + \frac{T_G}{8\pi^2} \log \frac{\text{Re}S(\mu)}{\text{Re}S'(\mu_0)} - \sum_r \frac{T_r}{8\pi^2} \log \frac{Z_r(\mu)}{Z'_r(\mu_0)},$$

where r runs all matter fields in the SSM, μ_0 is the cut-off scale of the theory, and b is a coefficient of beta function below the messenger scale. $S(\mu)$ is a holomorphic gauge coupling and primed quantities are the ones above the messenger scale. Here, we see that $Z_M(\mu_m)$ dependence drops out at low energy. Thus, a non-canonical Kähler potential does not contribute to the leading order of gaugino mass. Moreover, we could have assumed a spurion dependence of Kähler potential at the cut-off scale. Plugging the definition of real coupling $R'(\mu_0)$ at the cut-off scale, we see that $Z'_M(\mu_0)$ dependence also cancels out. Therefore, the leading order gaugino masses are not affected by spurion dependence of Kähler potential of messengers at all. However, if we impose a spurion dependence in $S'(\mu_0)$, it definitely contributes. Although it is nothing but adding gaugino masses by hand at the cut-off scale, it is contained in a frame work of the gauge mediation [4], since it vanishes in turning off the gauge coupling of the SSM. Usually in calculable models, these contributions, if exist, are generated by a heavy messenger around the cut-off scale and small compared to the leading term. Nevertheless, as we will see below, in some case it can be dominant and comparable to the sfermion masses.

2.3 Next to leading order gaugino mass

Since the sfermion masses generally arise at leading order, the vanishing gaugino masses at leading order implies that there is a hierarchy between gaugino and sfermion masses. One may consider the next to leading order gaugino masses to solve the hierarchy. There are several sources for non-vanishing gaugino masses at next to leading order⁵:

1. While the gaugino masses leading order in F at one-loop are prohibited, there is no problem for having non-vanishing gaugino masses at higher order in F . Explicit calculations

⁵Here, we focus on the gaugino masses generated at or above the messenger scale. At low-energy, the Higgs-Higgsino loop can also generate non-zero gaugino masses [35].

show that the next leading order contribution arises at $\mathcal{O}(F^3/M_{\text{mess}}^5)$ [10]. One might hope that the gaugino masses can be comparable to sfermion masses if $F/M_{\text{mess}}^2 \sim 1$. However, these higher order corrections are suppressed by small numerical coefficients in known examples and not sufficient to solve the hierarchy. Also, there is a phenomenological constraint on such a low scale mediation model [19].

2. Another possibility is that the gaugino masses are generated in $\mathcal{O}(F)$, but at higher loop level. In this case, using wave-function renormalization technique [25], one can explicitly show that the leading order gaugino mass in F at two-loop also vanishes if one-loop contribution does. Thus, the leading contribution is generated at best from three-loop diagrams. This contribution includes additional loop factors, so should be suppressed compared to the leading order.
3. As discussed in the previous subsection, there could be a contribution of a heavy messenger at the cut off scale or above, which would be of order $\mathcal{O}(F/M_{\text{heavy}})$. This type of contribution is suppressed to the leading order soft masses by a factor of $\mathcal{O}(M_{\text{mess}}/M_{\text{heavy}})$ compared to the leading order soft masses.

In any case, the gaugino masses are suppressed compared to the leading order and so to sfermion masses. This, combined with the current experimental lower bound for gaugino masses, indicates that the scales of the sfermion masses should be much higher than that of electroweak symmetry breaking, giving rise to fine tuning for the Higgs mass via top-stop loops. This is in contrast to the fact that relatively heavy gauginos at the messenger scale does not cause any problem because the sfermion masses are driven to be in the same order as gaugino masses at a lower scale by standard model renormalization group effects. In the next section, we explore the possibilities of solving this fine tuning problem using strong conformal dynamics in the hidden sector.

3 Conformal Sequestering

In the previous section, we learned that the vanishing of the gaugino masses at leading order gives rise to fine tuning associated to heavy sfermions and just modifying the Kähler potential does not generate a large enough gaugino masses. In this section, we show that the gaugino masses can naturally be larger than or comparable to the sfermion masses if the hidden sector flows into a strongly coupled fixed point below the messenger scale. We first see how the strong conformal dynamics in the hidden sector reduces the sfermion masses. Then we show how the next to leading order gaugino masses are sequestered, and provide a set of constraints the

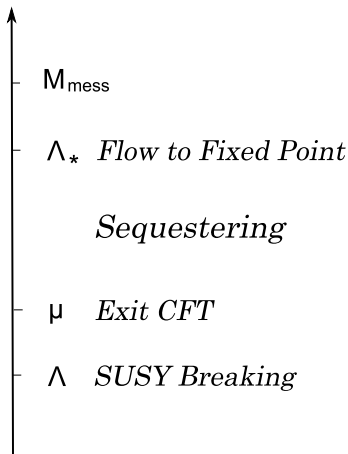


Figure 1: *Energy scales in our scenario*

parameters should satisfy to avoid fine tuning. As the hidden sector is assumed to be strongly coupled, one cannot calculate the anomalous dimensions of the non-holomorphic operators. But we expect that given a large number of SCFTs, one can find appropriate SCFTs which satisfy the constraints. To illustrate, we discuss the explicit numbers which are consistent with the constraints. We end this section with a discussion on μ/B_μ problem. We assume that supersymmetry is broken by a gauge singlet throughout this section.

3.1 Sequestering of scalar masses

Let us assume that a singlet chiral superfield X eventually gets a non-zero F -component vev and breaks SUSY. Before the hidden sector flows into the IR fixed point, messengers are integrated out at the scale M_{mess} . It generates interactions between X , X^\dagger and the SSM matters in Kähler potential and subsequently the sfermion masses. The operators which generate the sfermion masses are summarized as⁶

$$\int d^4\theta T(X, X^\dagger) Q_{\text{SM}}^\dagger Q_{\text{SM}}.$$

and the D -component of $T(X, X^\dagger)$ induces the sfermion masses. In general, the function $T(X, X^\dagger)$ can be expanded as follows:

$$T(X, X^\dagger) = c_0 |X|^2 + \dots$$

where the ellipsis contains a holomorphic and anti-holomorphic function of X which do not contribute to the scalar masses as well as higher order interaction terms.

⁶We impose messenger parity [1, 36] to prevent dangerous hypercharge D -terms from being generated in the hidden sector.

Our goal is to suppress these terms so that the sfermion masses can be lower than or in the same order as next to leading order gaugino masses. To do this, let us suppose that at some scale Λ_* below M_{mess} , the theory flows into a strongly coupled CFT (see Figure 1). Since the theory is strongly coupled, the anomalous dimensions of various operators such as X and $|X|^2$ are expected to be order one⁷. Then, the terms in the Kähler potential, or equivalently those in $T(X)$, are suppressed until the theory exits from the conformal regime at a scale μ . If the anomalous dimensions of X and $|X|^2$ are given by γ_X and $2\gamma_X + \alpha_*$, respectively, the operators in $T(X, X^\dagger)$ are sequestered as

$$T(X, X^\dagger) \rightarrow \epsilon^{2\gamma_X + \alpha_*} c_0 |X|^2 + \dots, \quad (3.1)$$

where the epsilon is defined by

$$\epsilon = \mu/\Lambda_*.$$

Note that the anomalous dimension of $|X|^2$ is not just the twice of the anomalous dimension γ_X because it is not a chiral operator and this fact is crucial to suppress the sfermion masses relative to the gaugino masses. Then, after the theory exits from a strongly coupled CFT at μ , supersymmetry is broken at the scale $\Lambda^2 := F < \mu^2$.

In this scenario, the leading order of the scalar masses below the scale μ are given by

$$m_s^2 = \epsilon^{2\gamma_X + \alpha_*} \left(\frac{g_{\text{SM}}^2}{16\pi^2} \right)^2 c_0 |F|^2 \sim \epsilon^{2\gamma_X + \alpha_*} \left(\frac{g_{\text{SM}}^2}{16\pi^2} \right)^2 \frac{|F|^2}{M_{\text{mess}}^2}, \quad (3.2)$$

where we took $c_0 \simeq 1/|M_{\text{mess}}|^2$ because this operator was generated by integrating out the messengers with mass M_{mess} ⁸.

Let us then discuss the sequestering of the gaugino masses. As discussed in the previous section, the following three contributions to the gaugino masses can be dominant when the leading order contribution (2.2) vanishes:

1. The first possibility is the case where the gaugino masses generated at order F^3/M_{mess}^5 at one-loop level dominate over others. These gaugino masses typically originate from non-holomorphic operators including D or \bar{D} . For example, a term

$$\int d^4\theta \frac{X^\dagger X X^\dagger D^2 X}{M_{\text{mess}}^6} W^\alpha W_\alpha,$$

⁷Since the conserved current is not renormalized, one have to be careful about global symmetries which may prevent these operators from acquiring large anomalous dimensions [26, 27].

⁸For simplicity, we also assumed that higher order terms have larger anomalous dimensions and so suppressed by larger powers of ϵ . Since the hidden sector is strongly coupled, however, there is no way to justify this assumption.

generates gaugino masses of order F^3/M_{mess}^5 if one assumes that the lowest component of X is in the same order as M_{mess} . This type of operators is not protected by the superconformal algebra, so one cannot determine the precise scaling at strong coupling. Nevertheless, one expects that for an appropriate choice of SCFT, the gaugino masses can be at the same order as the sfermion masses after sequestering. Let us assume that the anomalous dimension of the operator generating the $\mathcal{O}(F^3/M_{\text{mess}}^5)$ correction to be $4\gamma_X + \gamma_*$. Then, the gaugino masses after conformal sequestering are

$$m_\lambda^{1\text{-loop}} \sim \epsilon^{4\gamma_X + \gamma_*} \frac{g_{\text{SM}}^2}{16\pi^2} \frac{F^3}{M_{\text{mess}}^5} > m_\lambda^{3\text{-loop}}, m_\lambda^{\text{heavy}}.$$

If the condition $\epsilon^{\alpha_*/2 - 3\gamma_X - \gamma_*} \simeq F^2/M_{\text{mess}}^4$ holds, the gaugino mass is comparable to the sfermion masses.

2. The second possibility comes from the leading contribution in F at three loop level. As noted in the previous section, at two loop order the leading contribution in F also vanishes. This contribution dominates over others when

$$m_\lambda^{3\text{-loop}} \sim \epsilon^{\gamma_X} \left(\frac{g_{\text{SM}}^2}{16\pi^2} \right)^3 \frac{F}{M_{\text{mess}}} > m_\lambda^{1\text{-loop}}, m_\lambda^{\text{heavy}}.$$

If $\epsilon^{\alpha_*/2} \sim (g_{\text{SM}}^2/16\pi^2)^2$, the gaugino and scalar masses are in the same order. Since the gaugino masses are proportional to g_{SM}^6 , this scenario would induce a heavy gluino due to the comparatively large g_3 . This is problematic because the squark masses are roughly in the same order as gluino masses after the standard model renormalization group effects are taken into account, invoking a fine tuning problem. This problem can be avoided if the gauge couplings are close to each other at the messenger scale (*i.e.*, if the messenger scale is large) and will be discussed in the next subsection.

3. The last possibility is coming from a decoupled heavy messenger around the cut-off scale. There could be a contribution to the gaugino mass at leading order in F at one loop level. Since it typically comes from a operator $XW_\alpha W^\alpha$ generated at the heavy messenger scale M_{heavy} , one expect the following scaling under the sequestering.

$$m_\lambda^{\text{heavy}} \sim \epsilon^{\gamma_X} \frac{g_{\text{SM}}^2}{16\pi^2} \frac{F}{M_{\text{heavy}}} > m_\lambda^{1\text{-loop}}, m_\lambda^{3\text{-loop}}.$$

When this contribution dominates over others, by taking $\epsilon^{\alpha_*/2} \simeq M_{\text{mess}}/M_{\text{heavy}}$ one can make gaugino mass and scalar mass are in the same order.

Which effect can be dominant is the subject of the next subsection.

3.2 Constraints on the parameters

For simplicity, let us assume that the hidden sector enters into the conformal regime somewhere close to the messenger scale $\Lambda_* \sim M_{\text{mess}}$ and it exits near the scale of supersymmetry breaking, $\mu \sim \sqrt{F}$. For the first scenario, where one-loop F^3 contribution is dominant, to be realized, three-loop contribution should be sufficiently suppressed. This condition leads to

$$\epsilon^{3\gamma_X + \gamma_* + 4} > \left(\frac{g_{\text{SM}}^2}{16\pi^2} \right)^2. \quad (3.3)$$

This inequality implies that the lower bound of ϵ becomes larger for larger $3\gamma_X + \gamma_*$. Especially, if $3\gamma_X + \gamma_*$ is positive, ϵ must be larger than $\mathcal{O}(10^{-1})$. If we assume that the coupling in the hidden sector is small at the messenger scale, this may not be enough to guarantee the room for the small coupling constant at the messenger scale to become large at the conformal fixed point, so the conformal sequestering may not happen. For strongly coupled messengers, however, this is not a problem. Another constraint is that the gaugino masses should be comparable to or larger than the sfermion masses. This implies

$$\epsilon^{\alpha_*/2 - 3\gamma_X - \gamma_* - 4} \lesssim 1. \quad (3.4)$$

This condition may require large $|\alpha_*|$ and $|\gamma_*|$ while they can still be order one. Since $\epsilon < 1$, one can rewrite the inequality (3.4) as

$$\frac{\alpha_*}{2} - \gamma_* \gtrsim 4 + 3\gamma_X. \quad (3.5)$$

In any CFT, the conformal dimension of a gauge singlet, Lorentz scalar operator must be no less than one, thus γ_X must be positive. For $\gamma_X \gtrsim 2$, the left hand side of (3.4) should be larger than 10, so it requires somewhat large $|\alpha_*|$ and $|\gamma_*|$, while for small γ_X , $|\alpha_*|$ and $|\gamma_*|$ can be order one. Finally, we take the mass of the lightest chargino to be of order 10^2 GeV to avoid fine tuning:

$$\frac{g_2^2}{16\pi^2} M_{\text{mess}} \epsilon^{4\gamma_X + \gamma_* + 6} \sim 10^2 \text{ GeV}, \quad (3.6)$$

which determines the size of M_{mess} .

The argument above may be too abstract, so we can use numerical values to see that this scenario can be realized with order one anomalous dimensions. Let us take

$$\gamma_X \sim 0.5, \quad \gamma_* \sim -4, \quad \epsilon \sim 10^{-1.5}.$$

Then, (3.3) is satisfied. Eq.(3.4) determines the value of α_* to be $\gtrsim 3$, so let us just take $\alpha_* \sim 3$. The messenger scale and the supersymmetry breaking scale are fixed by (3.6) and

$$M_{\text{mess}} \sim 10^{11} \text{ GeV}, \quad \sqrt{F} \sim 10^{9.5} \text{ GeV}.$$

For this choice of parameters, we can also calculate the gravitino mass

$$m_{3/2} = \frac{F}{M_{\text{Pl}}} \sim \mathcal{O}(1 \text{ GeV}).$$

So, the anomaly and gravity mediation effects are also suppressed compared to the gauge mediation contributions. Note that the gravity mediation effect is also suppressed by conformal sequestering.

For the second scenario to be realized, the leading order contribution at three-loop should be larger than the F^3 contribution at one-loop. This is the case when

$$\epsilon^{3\gamma_X + \gamma_* + 4} < \left(\frac{g_{\text{SM}}^2}{16\pi^2} \right)^2. \quad (3.7)$$

The gaugino masses are larger than the sfermion masses after the hidden sector renormalization group effect is considered when

$$\epsilon^{\alpha_*/2} \lesssim \left(\frac{g_{\text{SM}}^2}{16\pi^2} \right)^2. \quad (3.8)$$

The condition that the lightest chargino is somewhere close to 10^2 GeV implies

$$\left(\frac{g_2^2}{16\pi^2} \right)^2 \epsilon^{\gamma_X + 2} M_{\text{mess}} \sim 10^2 \text{ GeV}. \quad (3.9)$$

To make the discussion concrete, let us again plug some numerical value to these conditions and see if they can be realized. Let us take

$$\gamma_X \sim 0, \quad \gamma_* \gtrsim 0, \quad \epsilon \sim 10^{-2}.$$

Then, (3.7) is satisfied. In order for (3.8) to be satisfied, α_* should be $\gtrsim 4$. Finally, our scenario provides an appropriate lightest chargino mass for

$$M_{\text{mess}} \sim 10^{12} \text{ GeV}, \quad \sqrt{F} \sim 10^{10} \text{ GeV}.$$

For these parameters, the gravitino mass is estimated as

$$m_{3/2} \sim \mathcal{O}(10 \text{ GeV}).$$

In this scenario, one also needs to check that the hierarchy between the gluino mass and the lightest chargino mass is sufficiently small at the messenger scale to avoid the hierarchy problem. At 10^{12} GeV , $\alpha_2/\alpha_3 \sim 3/4$, implying that the ratio between the gluino mass and the lightest chargino is of order one. We conclude that the hierarchy between the gluino masses and the lightest chargino mass is small enough to avoid a fine-tuning problem.

For the third scenario to be realized, the effect of heavy messengers should dominate over the other contributions. One-loop, F^3 contribution is naturally small if γ_* is large. The condition that the contribution from heavy messengers is larger than that from three-loop contribution gives

$$\frac{M_{\text{mess}}}{M_{\text{heavy}}} > \left(\frac{g_{\text{SM}}^2}{16\pi^2} \right)^2. \quad (3.10)$$

For the light messengers to play some role, the gaugino masses from the heavy messenger should be comparable to the sfermion masses from the light messengers. This gives a constraint

$$\frac{M_{\text{mess}}}{M_{\text{heavy}}} \sim \epsilon^{\alpha_*/2}. \quad (3.11)$$

Note that the contribution to sfermion masses from heavy messengers are suppressed by $m_{\text{mess}}/M_{\text{heavy}}$ relative to that from light messengers. Finally, we take the lightest chargino masses to be around 10^2 GeV, so

$$\epsilon^{\gamma_X+2} \frac{g_2^2}{16\pi^2} M_{\text{mess}} \sim 10^2 \text{ GeV}. \quad (3.12)$$

All the conditions are satisfied, for example, if we take

$$M_{\text{mess}} \sim 10^{11} \text{ GeV}, \quad M_{\text{heavy}} \sim 10^{13} \text{ GeV}, \quad \sqrt{F} \sim 10^9 \text{ GeV},$$

$$\gamma_X \sim 1, \quad \alpha_* \sim 2.$$

With these parameters, the gravitino mass is given by

$$m_{3/2} \sim \mathcal{O}(10^{-1} \text{ GeV}).$$

Again, the gravity and anomaly mediation effects are suppressed.

3.3 μ/B_μ Problem

μ and B_μ parameters can be defined as coefficients of supersymmetric and non-supersymmetric holomorphic Higgs bilinear terms in the Lagrangian, respectively:

$$\int d^2\theta \mu H_u H_d, \quad B_\mu H_u H_d. \quad (3.13)$$

For natural electroweak symmetry breaking, both μ and B_μ should be typically in the same order as electroweak scale. B_μ is a supersymmetry breaking parameter, so should be in the same order as electroweak symmetry breaking scale if one assumes that hierarchy problem is solved by supersymmetry. However, μ is a supersymmetry preserving parameter and there is

a priori no reason why it is also in the same order. This is so-called μ problem, and requires direct coupling between Higgs and hidden sector.

In gauge mediated supersymmetry breaking, there is another problem called B_μ problem. Generally in models of gauge mediation, μ and B_μ are generated at the same loop order. For a model given in Appendix A, for example, both of them are generated at one-loop. Typically,

$$\mu = \frac{\lambda_u \lambda_d}{16\pi^2} \frac{F^\dagger}{M_{\text{mess}}}, \quad B_\mu = \frac{\lambda_u \lambda_d}{16\pi^2} \frac{|F|^2}{M_{\text{mess}}^2}, \quad (3.14)$$

where $\lambda_{u,d}$ is a coupling between $H_{u,d}$ and messengers and should be small for perturbation theory to be valid. This implies

$$\mu^2 \sim \frac{\lambda_u \lambda_d}{16\pi^2} B_\mu \ll B_\mu, \quad (3.15)$$

which prevents natural electroweak symmetry breaking. For a general study and earlier attempts on μ/B_μ problem, see [37] and references therein.

Unlike the scenarios in [29, 30], where gaugino masses and μ are naturally in the same order, the gaugino masses are typically smaller than μ in our scenario. This implies that we must manually set them to be in the same order by unnaturally fine-tuned sequestering or small coupling constants $\lambda_{u,d}$. In this sense, our scenario does not solve the μ problem naturally. Nevertheless, B_μ can be naturally in the same order as μ and other soft masses after conformal sequestering and standard model renormalization group effects are taken into account, following the line of arguments employed in [29, 30]. Here, we should emphasize that μ must be fine-tuned to the same order as gaugino masses at the messenger scale, while B_μ just has to be smaller than μ and does not require any fine-tuning—if B_μ is smaller than μ at the messenger scale, it is automatically driven to the same order as μ due to the standard model renormalization group effects.

For concreteness, we work on a model where the next to leading order gaugino masses come from F^3 -terms at one-loop, while the above argument also works in other models discussed in Section 3.1. Also, we consider all the soft parameters related to the Higgs sector for completeness. Among the soft terms, the μ , and $A_{u,d}$ terms arise from operators linear in X . They are typically generated at one-loop by

$$c_\mu \int d^4\theta X^\dagger H_u H_d, \quad c_{A_{u,d}} \int d^4\theta X^\dagger H_{u,d}^\dagger H_{u,d}, \quad (3.16)$$

Here, c_μ and $c_{A_{u,d}}$ are numerical coefficients and proportional to $\lambda_u \lambda_d$ and $\lambda_{u,d}^2$, respectively. The scaling of these operators under hidden sector renormalization group flow is characterized

by γ_X , the anomalous dimension of X . To have μ comparable to the gaugino masses, we should set λ_u and λ_d to satisfy

$$\lambda_u \lambda_d \sim g_{\text{SM}}^2 \epsilon^{3\gamma_X + \gamma_* + 4}. \quad (3.17)$$

This condition may require small λ_u and λ_d , while that can be achieved if, e.g., the Higgs or messengers are composite.

The terms generated by operators quadratic in X require a bit of care. The operators relevant to the Higgs sector are

$$c_{m_{u,d}} \int d^4\theta X X^\dagger H_{u,d} H_{u,d}^\dagger, \quad c_{B_\mu} \int d^4\theta X X^\dagger H_u H_d, \quad (3.18)$$

The numerical coefficient are, up to order one coefficient, typically $(g_{\text{SM}}^2/16\pi^2)^2$, $\lambda_u \lambda_d/16\pi^2$, respectively. At first sight, one may expect that $c_{m_{u,d}}$ and c_{B_μ} are suppressed by hidden sector renormalization group effects. As was pointed out in [29, 30], however, the operators quadratic in X mix with those linear in X and what ends up being suppressed is a specific combination of the coefficients $c_{m_{u,d}}$, c_{B_μ} , c_μ and $c_{A_{u,d}}$. A careful operator product analysis [38] shows that the combinations that are sequestered by the anomalous dimension of XX^\dagger are

$$c_{m_{u,d}} - \frac{1}{2}C(2 - \alpha_*)(|c_\mu|^2 + |c_{A_{u,d}}|^2), \quad c_{B_\mu} - \frac{1}{2}C(2 - \alpha_*)\text{Re}(c_\mu(c_{A_u}^* + c_{A_d}^*)),$$

where C is the OPE coefficient of X and X^\dagger and assumed to be order one. Since $m_{u,d}^2 = -(c_{m_{u,d}} - |c_{A_{u,d}}|^2)|F|^2$ and $B_\mu = -(c_{B_\mu} - \text{Re}(c_\mu(c_{A_u}^* + c_{A_d}^*)))|F|^2$, they are naturally in the same order as μ and $A_{u,d}$ after sequestering if α_* is sufficiently large and λ_u and λ_d are taken to be in the same order. The value of α_* is constrained by

$$\frac{\lambda_u \lambda_d}{16\pi^2} \gtrsim \epsilon^{\alpha_*},$$

which can be rewritten, using (3.17), as

$$\frac{g_{\text{SM}}^2}{16\pi^2} \gtrsim \epsilon^{\alpha_* - 3\gamma_X - \gamma_* - 4}.$$

This is satisfied for the choice of parameters in Section 3.2.

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Appendix A Generating μ and B_μ

In this Appendix, we present a simple model that generates μ and B_μ at leading order. Let us consider a simple messenger superpotential,

$$W_{\text{mess}} = m_1(X)(\phi_1\bar{\phi}_1 + \phi_2\bar{\phi}_2) + m_2(X)(\phi'_1\bar{\phi}'_1 + \phi'_2\bar{\phi}'_2),$$

with coupling to Higgs being given by

$$W_{\text{Higgs}} = \lambda_u H_u \phi_1 \phi_2 + \lambda_d H_d \bar{\phi}_1 \bar{\phi}_2 + \lambda'_u H_u \phi'_1 \phi'_2 + \lambda'_d H_d \bar{\phi}'_1 \bar{\phi}'_2,$$

The condition for anomalously small gaugino is given by

$$\frac{\partial}{\partial X}(m_1(X)m_2(X)) = 0. \tag{A.1}$$

For this superpotential, the μ -term at the leading order was calculated in [39]:

$$\begin{aligned} \mu &= -\frac{F}{16\pi^2} \left[\lambda_u \lambda_d \left. \frac{m'_1(X)}{m_1(X)} \right|_0 + \lambda'_u \lambda'_d \left. \frac{m'_2(X)}{m_2(X)} \right|_0 \right] \\ &= -\frac{F}{16\pi^2} \left[\lambda_u \lambda_d \left(\left. \frac{\partial}{\partial X} \ln m_1(X)m_2(X) \right|_0 \right) + (\lambda'_u \lambda'_d - \lambda_u \lambda_d) \left. \frac{m'_2(X)}{m_2(X)} \right|_0 \right]. \end{aligned}$$

The first term in the second line vanishes because of (A.1), but the second term is non-zero in general (except for the case with $\lambda_u \lambda_d = \lambda'_u \lambda'_d$). We conclude that the μ -term is generically generated at the leading order. Similarly, these superpotential interactions also generate B_μ -term. The result is

$$B_\mu = -\frac{F}{16\pi^2} \left[\lambda_u \lambda_d \left(\left. \frac{m'_1(X)}{m_1(X)} \right|_0 \right)^2 + \lambda'_u \lambda'_d \left(\left. \frac{m'_2(X)}{m_2(X)} \right|_0 \right)^2 \right],$$

which is also non-vanishing. In general, μ and B_μ do not vanish if each messenger couples to Higgs superfields differently.

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