Semi-direct Gauge-Yukawa Mediation Postprint

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Full Text

Preamble

Semi-direct Gauge-Yukawa Mediation

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Abstract

We propose semi-direct Gauge-Yukawa mediation of supersymmetry (SUSY) breaking. The messenger fields mediating SUSY breaking to the visible sector do not directly couple with the goldstino field, and instead they have gauge and Yukawa interactions with some primary messenger fields which couple directly with the goldstino fields. From explicit Feynman diagram calculations for the SUSY breaking soft masses, we find that the SUSY particle spectra can be realistic. In particular, this generalization of semi-direct gauge mediation solves the massless gaugino problem since the holomorphic soft mass terms of the messenger fields can be generated by Yukawa couplings.

We also provide arguments that this scenario can be realized naturally in some dynamical SUSY-breaking models such as ISS-like models.

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Introduction and Motivation

The mechanism for supersymmetry (SUSY) breaking remains an open question. Due to the tree-level sum rule for the SUSY-breaking spectrum, namely the supertrace theorem, SUSY must be broken spontaneously in the hidden sector [**?**]. The crucial issue then becomes how to mediate the SUSY-breaking effects to the visible sector. Several schemes have been proposed for this purpose, among which Gauge Mediated SUSY Breaking (GMSB) is a particularly promising approach because it automatically avoids the notorious flavor problem [**?**] (for reviews, see, e.g., [**?**]).

In particular, direct gauge mediation has recently attracted much attention, in which the messenger fields play a role in determining the SUSY-breaking vacua [**?**]. However, this attractive approach suffers from several problems, such as the suppressed gaugino mass problem [**?**], the Landau pole problem, and the fine-tuning problem. On the other hand, indirect gauge mediation typically does not have the gaugino mass problem but has the disadvantage of requiring messengers to be added by hand. A recent study of such messenger gauge mediation is given in Ref. [**?**].

In both direct and indirect gauge mediation, the messenger fields Φ must carry SUSY-breaking information in the form of mass splittings ($m^2_{\text{fermionic}} - m^2_{\text{bosonic}} \neq 0$) between the fermionic and bosonic components, which are degenerate in a SUSY-preserving theory. However, how the messengers obtain such mass splittings is unknown. In previous studies, a singlet $X = M + \theta^2 F$ is simply introduced as the SUSY-breaking source that couples directly to the messengers. This is the so-called Minimal Gauge Mediation (MGM) and has been widely considered in phenomenology.

This simplification in MGM may conceal important questions. For instance, why must the messengers couple directly to X to obtain mass splitting? In principle, the mass splitting can be achieved in a cascade manner: some primary messengers ϕ_0 (neutral under the Standard Model (SM) gauge group G_{SM}) mediate SUSY breaking to some secondary messengers ϕ_1 via hidden sector gauge interactions with a gauge group G_h or via Yukawa interactions. This dynamical process can be successive if necessary until finally mediating SUSY breaking to the visible sector.

In this paper, we assume the simplest cascade pattern, namely the two-step cascade mediation, as depicted in Fig. 1 [Figure 1: see original paper]. In fact, such a two-step cascade mediation via pure gauge interaction is precisely the semi-direct gauge mediation studied intensively in the literature [**?**]. However, the analogous mediation mechanism with additional Yukawa interactions has not been studied, and this will be examined in the present work.

Conventional semi-direct gauge mediation suffers from the massless gaugino problem because the secondary messengers do not acquire holomorphic soft terms [?]. However, with Yukawa-assisted mediation, the holomorphic soft terms arise at one-loop order, thus allowing gauginos to become massive. We shall show that a calculable cascade SUSY-breaking scenario with the minimal Kähler potential, based on a very general setup, is viable only when both gauge and Yukawa interactions are present. In other words, this Hybrid Cascade Gauge-Yukawa Mediation (HCGYM) can provide viable SUSY particle spectra in some cases. The crucial point is the presence of a non-vanishing supertrace in the messenger sector that links to the visible sector and modifies the sfermion masses drastically.

Our cascade scenario has one essential difference from gauge mediation with a messenger sector added by hand: our scenario can be realized dynamically in models such as Intriligator-Seiberg-Shih (ISS)-like models [?]. In other words, the messengers in our scenario are a natural part of the hidden sector that breaks SUSY. We simply take a different approach by gauging the so-called global flavor symmetry in the hidden sector. However, in our study we will start from a class of weakly coupled effective O' Raifeartaigh (OR) models (some characterized by cubic terms and some large global symmetry gauged properly), without elaborating on their dynamical realization.

[Figure 1: see original paper] A schematic diagram showing the structure of the two-step HCGYM.

The paper is organized as follows. In Section II we propose a general framework for HCGYM and calculate the SUSY breaking soft terms via the secondary messengers. Our calculations will be performed explicitly in a very general setup. In Section III we analyze under what conditions realistic soft masses in the visible sector can be achieved, and comment on possible general features of HCGYM. Finally, discussions and conclusions are given in Section IV.

II. Hybrid Cascade Gauge-Yukawa Mediation

A. The Primary SUSY-Breaking Mediation

We first focus on SUSY breaking mediation without specifying any concrete model. Generally, a cascade mediation via gauge and Yukawa (hybrid) interactions has a typical structure as follows:

(i) A hidden sector, which breaks SUSY, is composed of the Goldstino superfield X and the primary vector-like messenger fields $(\phi_i, \bar{\phi}_i)$ with i being the gauge/flavor index of the hidden sector group G_h such as $SU(N)$.

(ii) The secondary messenger fields $(F, \bar{F})$, which are charged under both G_h and the SM gauge group, and the vector-like fields $(f, \bar{f})$, which carry only SM quantum numbers, both have supersymmetric mass terms and obtain SUSY-breaking effects via interactions with the hidden gauge fields as well as from cubic Yukawa coupling terms (discussed below).

(iii) Some cubic terms must be present, which couple one primary and one secondary messenger field with the SM vector-like fields, or couple two secondary messengers with one primary messenger. However, cubic terms with two primary fields are not allowed by SM gauge invariance. Note that if we turn off these cubic terms, the hidden sector SUSY breaking is not affected because the secondary messengers are not relevant to the SUSY-breaking dynamics, although they couple to the hidden sector. In this case, our scenario reduces to the semi-direct gauge mediation discussed in Ref. [?].

We stress that our scenario is significantly different from semi-direct gauge mediation [?]. Semi-direct gauge mediation [?] interpolates between MGM and direct gauge mediation. It introduces a messenger sector that does not affect SUSY breaking and does not have to couple to X, but uses this sector to probe the hidden sector only via hidden gauge interactions. In our approach, direct renormalizable couplings in the superpotential play a central role. As a result, the spectrum of the secondary messengers and the visible fields is greatly modified. Such cubic terms can arise in dynamical models, and moreover their mass scale can be dynamically determined. In other words, our framework regards the messenger sector as a built-in part of dynamical model building, rather than something introduced by hand.

With this general setup, we propose to use an effective OR model that exhibits the dynamical structure. In our convention, we choose a basis where the light messengers are diagonal while the hidden sector takes a general form:

$$W = X\bar{\phi}_a\phi_a + (\lambda_{ab}X + m_{ab})\bar{\phi}_a\phi_b + \lambda_a\phi_a\bar{F}_i\bar{f} + \bar{\lambda}_a\bar{\phi}_aF_if + M_FF_i\bar{F}_i + M_ff\bar{f}$$

Here, the index a does not refer to any symmetry, and we require $M_{F,f} \ll M_h$, where M_h is defined as the typical mass scale in the hidden sector, while $\sqrt{F}$ determines the SUSY-breaking scale. Note that multiple pairs of secondary messengers give additive contributions and thus do not change our conclusion. By the way, we may embed the SM gauge group into the group G_h, which then induces the massless gaugino problem [?] where $\log\det M_{\text{primary}} = 0$ denotes the supersymmetric mass matrix for the hidden sector messenger.

Our general setup can be equipped with R-symmetry. With R-symmetry, our framework can be considered as a generalization of (extra)ordinary GMSB [?] with cubic terms and the extra gauge group G_h. The particle charges under R-symmetry are: $R(X) = 2$, $R(F)+R(\bar{F}) = 2$, $R(f)+R(\bar{f}) = 2$, $R(\bar{\phi}_a)+R(\phi_b) = 0$ (for $\lambda_{ab} \neq 0$), $R(\phi_a) + R(\bar{F}) + R(f) = 2$ (for $\lambda_a \neq 0$), $R(\bar{\phi}_a) + R(\phi_b) = 2$ (for $m_{ab} \neq 0$), $R(\bar{\phi}_a) + R(F) + R(\bar{f}) = 2$ (for $\bar{\lambda}_a \neq 0$). To make gauginos massive, R-symmetry must be broken. This can be achieved by introducing proper structures in the hidden sector, leading to spontaneous breaking of R-symmetry radiatively [?] or at tree-level [?]. Here, we will not scrutinize the R-symmetry breaking mechanism. Instead, we simply assume it can be realized.

Next, we discuss the SUSY breaking encoded in the secondary messengers. They feel SUSY-breaking effects through two ways. One is through conventional gauge mediation at the two-loop level, of the non-holomorphic form such as $m_F^2 \sim (g_h^4/(16\pi^2)^2)|F/M_h|^2$ (because F and f are similar, for simplicity we will focus only on the former). They also receive direct one-loop contributions from the Yukawa couplings, which are calculated explicitly in Appendix IV A and can be extracted using the wave function renormalization method [?]. In short, these three terms are given by:

$$m_{F,g}^2 \simeq \frac{g_h^4}{(16\pi^2)^2}\frac{|F|^2}{M_h^2}, \quad m_{F,Y1}^2 = \frac{\lambda^4}{(16\pi^2)^2}(N+2)\frac{|F|^2}{M_h^2}, \quad m_{F,Y2}^2 = \frac{\lambda^2 g_h^2}{(16\pi^2)^2}\frac{|F|^2}{M_h^2}$$

where g_h is the gauge coupling of G_h and its value will be constrained by phenomenology. The explicit expressions of the Σ_i functions are given in Eqs. (37), (38), and (39). Such terms enter the diagonal elements of the bosonic components of the secondary messengers and behave as D-term SUSY breaking. In addition to these non-holomorphic terms, the other soft terms are holomorphic terms, namely $m_{\rm hol}^2 F\bar{F}$ + h.c., which are generated at one-loop level due to Yukawa interactions and proportional to the secondary messenger masses:

$$m_{\rm hol}^2 \equiv m_{F\bar{F}}^2 = m_{F^*\bar{F}^*}^2 = i(\Sigma_1 + \Sigma_2 + \Sigma_3) \sim \frac{\lambda^2}{16\pi^2}\frac{F}{M_h}M_F$$

where we have set $M_F \simeq m_f$. This term contributes to the non-diagonal mass terms of the secondary messengers and acts as conventional F-term SUSY breaking. It is absent in conventional semi-direct gauge mediation, rendering gauginos massless at leading order in SUSY breaking.

A comment is in order. For semi-direct gauge mediation, there is a robust gaugino screening theorem [?]. It is derived through the renormalization of the real physical superfield:

$$R(\mu) = S(\mu) + S^\dagger(\mu) + 8\pi^2 \log\det Z_r(\mu)$$

where $S(\mu)$ is the holomorphic gauge coupling that runs only at one-loop and r runs over light fields at scale μ. The screening depends heavily on the replacement $\mu^2 \to \mu^2 Z_F(\mu_F)$ after integrating out all messenger scales, leading to $R(\mu)$ being independent of X at leading order. However, just as in the analysis of [?] which introduced chiral messengers so that the two chiralities of the messengers have different wave functions and thus invalidate the replacement, in our case the chiral Yukawa coupling between secondary and primary messengers can also invalidate the replacement. Consequently, the screening is avoided.

By the way, in our framework we assume that G_h is not Higgsed. If it is Higgsed, the gauge mediation receives modifications controlled by a new parameter $y =$

M_V^2/M_h^2, where M_V is the vector boson mass scale [?]. When $y \to 0$, we have the complete gauge group, while conversely we have no gauge symmetry at the messenger scale. A sufficiently high M_V is not a desired case, and the reason will be discussed below.

B. Gaugino and Sfermion Masses

Having studied the SUSY breaking mediated to the secondary messengers by hybrid mediation, we now calculate the SUSY breaking mediated to the visible sector by conventional GMSB. In contrast to simple GMSB models where the hidden sector messenger spectrum has a vanishing supertrace, in our cascade framework the supertrace is nonzero due to radiative corrections that essentially arise from the holomorphic soft mass term $m^2_{F\bar{F}}$. In the basis $(F, \bar{F})$, the mass matrix of the scalar component is:

$$M^2_{\text{scalar}} = \begin{pmatrix} M_F^2 + m_F^2 & m^2_{\text{hol}} \\ m^{2*}_{\text{hol}} & M_F^2 + m^2_{\bar{F}} \end{pmatrix}$$

where $m_F^2 = m^2_{\bar{F}} + D$ and $D = m^2_{\bar{F}} - m^2_F$ measures the supertrace. The holomorphic term m^2_{hol} can be effectively treated as a spurion superfield X which has a vanishing vacuum expectation value (VEV) for its lowest component but a non-zero θ^2 component. From this perspective, our framework provides a natural realization of multi-spurion fields, which is helpful for implementing general gauge mediation [?]. The term $D = m^2_{\bar{F}} - m^2_F$ measures the supertrace and has deep implications for soft masses in the visible sector.

The explicit calculations have been carried out in Ref. [?], and we use their results to calculate the soft masses of gauginos and sfermions. In the limit of small SUSY breaking $D, F \ll M_F^2$, the generic sfermion mass squared is approximately given by:

$$m^2_{\tilde{f}} \simeq \sum_a \left(\frac{g_a^2}{16\pi^2}\right)^2 C_a \left[2\frac{|F|^2}{M_F^2} + 4D\log\frac{\Lambda}{M_F}\right]$$

where C_a is the Dynkin index related to the three SM gauge groups, and Λ is the ultraviolet cutoff scale related to a high scale at which D and F are generated—in fact, this is the primary messenger scale $M_{F,f}$. This logarithmic UV-dependence is a characteristic contribution of cascade gauge mediation. As for gaugino masses, the D-type contribution does not affect them significantly, and thus they take the conventional form:

$$M_{\lambda_a} \simeq \frac{g_a^2}{16\pi^2}\frac{F}{M_F}$$

Now some comments are due regarding the properties of soft masses in the presence of D-type contributions. From Eq. (10) we obtain:

$$m_{\tilde{f}}^2 \sim \frac{g^4}{(16\pi^2)^2}\frac{|F|^2}{M_F^2}\left[1+\frac{2D}{|F|^2/M_F^2}\log\frac{\Lambda}{M_F}\right]$$

Thus, depending on the sign of D, the sfermion mass squared can be either enhanced ($D<0$) or reduced ($D>0$). As calculated in the previous section, there are several sources of non-holomorphic terms, and in the following we will show that a pure Yukawa cascade model is not viable in practice.

First, in cascade gauge mediation, a contribution from pure gauge interaction is necessary. In the limit $g_h \to 0$, generically the one-loop effect is dominant and scales as (dropping the gauge index a for simplicity):

$$D \sim m_F^{(1)2} - m_{\bar{F}}^{(1)2} \sim \frac{\lambda^4}{(16\pi^2)}\frac{|F|^2}{M_h^2}\mathrm{sign}(D)$$

We will find that the sign is not uniquely predicted but depends on parameters in the hidden sector. On the other hand, the holomorphic term scales as:

$$m_{\mathrm{hol}}^2 \sim \frac{\lambda^2}{16\pi^2}\frac{F}{M_h}M_F$$

which is suppressed by the light supersymmetric mass term of the secondary messengers. Then we find that the D-type contribution dominates over the sfermion mass term. Concretely, in Eq. (10) the second term is about $(16\pi^2/\lambda^2)$ times larger than the first term, which is at the same order as the gaugino mass. As a result, it leads to a split spectrum, which will incur large fine-tuning in the Minimal Supersymmetric Standard Model (MSSM).

The above conclusion is obtained only when the holomorphic and non-holomorphic SUSY-breaking terms are generated at the same loop level. We note that in some cases, two-loop order SUSY-breaking contributions can dominate over one-loop contributions that are suppressed by higher powers of M_F/M_h due to subtle cancellations [?]. This happens when the hidden sector essentially contains only one mass scale, as in minimal GMSB with X being the spurion field. In our model, this is approximately satisfied in the limit $M_F \to 0$, but unfortunately a realistic spectrum cannot be obtained. In fact, the D-type contribution scales as:

$$D \sim \frac{\lambda^4}{(16\pi^2)^2}\frac{|F|^2}{M_h^2}$$

Then the sfermion mass is given by:

$$m_{\tilde{f}}^2 \sim \frac{g^4}{(16\pi^2)^2}\frac{|F|^2}{M_h^2}\left[1-\frac{2(N+2)\log(\Lambda/M_h)}{16\pi^2}\lambda^4\right]$$

which is always negative and unacceptable since it breaks the SM gauge group G_{SM} at the messenger scale.

Now we consider the second limit $\lambda, \bar{\lambda} \to 0$ with nonzero g_h. In this case we recover the conventional semi-direct gauge mediation scenario. Studies in Refs. [**?**] pointed out that, independent of hidden sector details, the doubly charged messengers do not acquire F-type SUSY-breaking information from it, and thus the gaugino mass vanishes at leading order (gaugino screening). In Ref. [**?**] some complicated chiral messenger fields were proposed to overcome this difficulty. In our HCGYM approach, we attempt to solve this by turning on Yukawa couplings.

In our HCGYM scenario, we find that for proper values of g_h/λ, the sfermion mass squared can be reduced as a result of the non-vanishing supertrace. The sfermion masses receive D-type contributions as:

$$m_{\tilde{f}}^2 \sim \frac{g^4}{(16\pi^2)^2}\frac{|F|^2}{M_h^2}\left[1+\frac{2D}{|F|^2/M_h^2}\log\frac{\Lambda}{M_h}\right] \sim \frac{g^4}{(16\pi^2)^2}\frac{|F|^2}{M_h^2}\left[1+\mathrm{sign}(D)\frac{16\pi^2}{\lambda^2}\right]$$

Therefore, for $\mathrm{sign}(D)=1$, there exists cancellation between gauge and Yukawa contributions. This D-type contribution approximately vanishes and the sparticle spectrum reduces to the MGM case when:

$$\frac{g_h^2}{\lambda^2} \sim \frac{1}{N+2}$$

In fact, we can tune the couplings to obtain our required soft masses. Note that the reduced sfermion spectrum generically requires tuning of the ratio λ/g_h, which arises from cancellation between the one-loop effect ($\propto \lambda^2$) and the two-loop effect ($\propto g_h^2$). By the way, increasing N helps to lower the required value of g_h, which may however induce a Landau pole for the $SU(3)_C$ gauge coupling.

If two-loop Yukawa corrections are dominant, different results are obtained. Then the D-type contribution is:

$$D \sim \frac{\lambda^4}{(16\pi^2)^2}\frac{|F|^2}{M_h^2}, \quad m_{\mathrm{hol}}^2 \sim \frac{\lambda^2}{16\pi^2}\frac{F}{M_h}M_F$$

Due to the fact that both D and F-type contributions come from the two-loop level, they have the same sign, and the additional gauge contribution makes the sfermion mass quite large (negative).

Our above analysis is just a rough estimation. In the next subsection we will give a concrete example and perform numerical calculations to show that our framework indeed works.

C. A Concrete Example and A Preliminary Dynamical Realization

To demonstrate that our HCGYM scenario can satisfy the required properties, we now study numerically a concrete example. For this simple example the superpotential is assumed to be:

$$W = X\bar{\phi}_1\phi_1 + \mu(\bar{\phi}_1\phi_2 + \bar{\phi}_2\phi_1) + \lambda\phi_2 F\bar{f} + \bar{\lambda}\bar{\phi}_2\bar{F}f + m_F F\bar{F} + m_f f\bar{f}$$

where X is regarded as the spurion superfield, both $(\phi, \bar{\phi})$ and $(F, \bar{F})$ belong to the fundamental representation of the hidden gauge group $SU(N)$, while $(f, \bar{f})$ is neutral. We set $X = M + F_X\theta^2$ with $M = 1$, $\mu = 1$, and $F_X = 0.1$. Then we obtain the D and F terms as:

$$D = m_F^2 - m_{\bar{F}}^2 = \frac{0.07}{(16\pi^2)^2}\lambda^2\frac{|F_X|^2}{M^2}, \quad m_{\text{hol}}^2 = m_{F\bar{F}}^2 = \frac{0.008}{16\pi^2}\lambda^2\frac{F_X}{M}M_F$$

Another messenger pair $(f, \bar{f})$ also obtains soft D- and F-type corrections which are enhanced by a group factor N, but receives no corresponding gauge contribution since they are charged under G_{SM} only. With all these contributions, from Eq. (10) we find that for sfermion masses the D-type contribution has a ratio:

$$R \equiv \frac{m_{\tilde{f}}^2(D\text{-type})}{m_{\tilde{f}}^2(F\text{-type})} \sim \frac{2(1+N)\log(\Lambda/M_F)}{(16\pi^2/\lambda^2) + g_h^4(0.07\lambda^2)^2(N^2+1)/2N - 0.008}$$

Note that this result is independent of M_F. In Fig. 2 [Figure 2: see original paper] we plot this ratio as a function of g_h (setting $\lambda = 1$). As mentioned previously, the ratio can be negative and thus can reduce sfermion masses with some tuning of λ/g_h.

[Figure 2: see original paper] The ratio R of D-type and F-type contributions as a function of g_h. Here we take $N = 4$ and $\lambda = 1$, $\Lambda/M_F = 10^4$. The right-handed half of the curve covers the region with $R < 0$.

Let us comment on the dynamical realization of our framework. In principle, any conventional semi-direct gauge mediation model can be upgraded to our HCGYM scenario by coupling the secondary messengers to the hidden sector. A more natural model can be found in Ref. [?], which studied metastable SUSY-breaking vacua in $SU(N_c)$ supersymmetric QCD (SQCD) with N_f vector-like quarks in the superconformal window $3/2N_f < N_c < 2N_f$. Based on the SQCD in Ref. [?], with some deformations, Ref. [?] studied its Seiberg dual description:

$$W = Kp\bar{q} + L\bar{p}q + Nq\bar{q} + Mpp$$

The flavor symmetry is $SU(N_f^{(1)}) \times SU(N_f^{(2)})$ and the magnetic gauge group is $SU(\tilde{N}_c)$ with $\tilde{N}_c = 2N_f - N_c$. Various fields are assigned in representations as:

	$SU(N_f^{(1)})$	$SU(N_f^{(2)})$	$SU(\tilde{N}_c)$	$U(1)_R$
K	$\bar{N}_f^{(1)}$	N_f	1	2
L	$N_f^{(1)}$	$\bar{N}_f$	1	2
N	1	$\mathrm{Adj}+1$	1	2
M	$\mathrm{Adj}+1$	1	1	2
p	$N_f^{(1)}$	1	$\tilde{N}_c$	0
$\bar{p}$	$\bar{N}_f^{(1)}$	1	$\bar{\tilde{N}}_c$	0
q	1	$N_f^{(2)}$	$\tilde{N}_c$	0
$\bar{q}$	1	$\bar{N}_f^{(2)}$	$\bar{\tilde{N}}_c$	0

Aside from the last line, this gives the content of the ISS model [?]. It has a SUSY-breaking vacuum, with superfields parameterized as:

$$\tilde{q} = \begin{pmatrix} \mu + \sigma_1 & \phi_2 \\ \phi_3 & \mu + \sigma_2 \end{pmatrix}, \quad N = \begin{pmatrix} \sigma_3 & \phi_4 \\ \phi_5 & \sigma_6 \end{pmatrix}, \quad X = p\tilde{p}, \quad Y = p\tilde{q}, \quad \tilde{Y} = \tilde{p}q$$

Expanding around this vacuum, we obtain:

$$W \supset X\phi_1\phi_2 - \mu^2 X + \mu(\phi_1\phi_4 + \phi_2\phi_3) + \mu(\phi_5\phi_8 + \phi_6\phi_7) + Y\phi_2\phi_5 + \tilde{Y}\phi_1\phi_6 + M_p\phi_5\phi_6$$

As shown in Ref. [?], when the flavor symmetry satisfies the relation:

$$\frac{3}{2}N_f^{(2)} < \tilde{N}_c < N_f^{(1)} + N_f^{(2)} < 3N_f^{(2)}$$

the theory is weakly coupled in the IR and has tree-level SUSY-breaking vacua. Choosing $N_f^{(1)} = 2$, $N_f^{(2)} = 5$, $N = 3$, and embedding $G_{\mathrm{SM}} \subset SU(N_f^{(2)})$, we find that ϕ_{1-4} play the role of primary messengers while ϕ_{5-8} are secondary messengers. From Eq. (27) it is clear that hidden sector SUSY-breaking effects mediated by gauge and Yukawa interactions to the secondary messengers indeed appear in the model. Of course, this is not a realistic model since R-symmetry is not broken. To make it realistic, further modifications in the hidden sector are required, which is beyond the scope of this paper.

III. Discussion and Conclusion

We proposed a pattern of cascade gauge-Yukawa mediation of SUSY breaking based on explicit Feynman diagram calculations for the soft mass terms. We found that realistic soft masses for gauginos and sfermions can be obtained when conventional messenger fields acquire SUSY-breaking effects through gauge and Yukawa interactions with the hidden sector. This hybrid scenario can easily avoid the massless gaugino problem and reduce sfermion masses by fine-tuning between λ and g_h.

Finally, we note that very recently such a cascade pattern was utilized for SUSY breaking (not mediation) in Refs. [**?**]. And R-symmetry spontaneous breaking at the two-loop level [**?**] is also a result of cascade effects.

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IV. Appendix

A. Soft Mass Terms from the One-Loop Yukawa Mediation

Here we present explicit one-loop calculations for the soft terms of the secondary messengers, using dimensional regularization. First, the mass-squared matrix of the primary messenger bosonic components $(\phi, \bar{\phi}^*)$ and the Dirac fermion mass matrix are respectively given by:

$$B = \begin{pmatrix} (\lambda X + m)^\dagger(\lambda X + m) & (\lambda X + m)^\dagger F_X \\ F_X^\dagger(\lambda X + m) & (\lambda X + m)(\lambda X + m)^\dagger \end{pmatrix}, \quad M_F = \lambda X + m$$

where M_F is an $N_p \times N_p$ matrix while B is $2N_p \times 2N_p$. They can be diagonalized by unitary matrices U, N_L, and N_R respectively:

$$U^\dagger B U = M_1, \quad N_L^\dagger M_F N_R = M_2$$

For convenience in our following calculations, we parameterize the matrix M_1 in block form:

$$M_1 = \begin{pmatrix} C & D \\ D^\dagger & C^T \end{pmatrix}, \quad M_2 = \mathrm{diag}(m_{1i}, m_{2i})$$

Then the soft terms are given by:

$$V_{\text{soft}} = (m^2_{F\bar{F}} F\bar{F} + \text{h.c.}) + m^2_{FF^*} F^* F + m^2_{\bar{F}\bar{F}^*} \bar{F}^* \bar{F}$$

The messenger fields $(F, \bar{F})$ get SUSY-breaking soft mass terms through one-loop diagrams shown in Fig. 3 [Figure 3: see original paper] and Fig. 4 [Figure 4: see original paper], which are given by:

$$m^2_{F\bar{F}} = m^2_{F^*\bar{F}^*} = i(\Sigma_1 + \Sigma_2 + \Sigma_3), \quad m^2_{FF^*} = m^2_{\bar{F}\bar{F}^*} = i(\Sigma_4 + \Sigma_5 + \Sigma_6)$$

[Figure 3: see original paper] Feynman diagrams contributing to the F-type soft terms.

[Figure 4: see original paper] Feynman diagrams contributing to the D-type soft terms.

The Σ functions are given by:

$$\Sigma_1 = \frac{1}{4\pi^2} \left[\bar{\lambda}_a (\lambda_{ab} X + m_{ab})^* (A_{bi} A^\dagger_{ic})^* \lambda_c - \frac{1}{2} \right] f(m^2_{1i}, M^2_f) + (A \to C, m_1 \to m_2)$$

$$\Sigma_2 = \frac{1}{4\pi^2} \left[\bar{\lambda}_a (C_{ai} C^\dagger_{ib})^* (\lambda_{bc} X + m_{bc})^* \lambda_c - \frac{1}{2} \right] f(m^2_{2i}, M^2_f) + (C \to A, m_2 \to m_1)$$

$$\Sigma_3 = \frac{1}{4\pi^2} \left[\bar{\lambda}_a (A_{ai} C^\dagger_{ib})^* (\lambda_{bc} X + m_{bc})^* \lambda_c - \frac{1}{2} \right] f(m^2_{1i}, M^2_f) + (A \leftrightarrow C, m_1 \leftrightarrow m_2)$$

Repeated indices should be summed over. The first two terms come from scalar loops (the left plot in Fig. 3), and $\Sigma_1 \propto M_f$ is orders of magnitude smaller than the second term. That is, the holomorphic corrections are dominated by terms proportional to the low-energy mass of the messengers themselves. The fermion loop contributes terms $[\lambda_a U_{Rai} U^\dagger_{Lib} \bar{\lambda}_b / (4\pi^2)] M_{2i} M_f f(M^2_{2i}, M^2_f)$.

The non-holomorphic terms from Fig. 4 are given by:

$$\Sigma_4 = \frac{1}{4\pi^2} \left[\bar{\lambda}_a (\lambda_{ab} X + m_{ab}) (A_{bi} A^\dagger_{ic})^* (\lambda_{cd} X + m_{cd})^* \bar{\lambda}^*_c \right] \left[\log \frac{m^2_{1i}}{M^2_f} - 1 \right] + (A \to C, m_1 \to m_2)$$

$$\Sigma_5 = \frac{1}{4\pi^2} \left[\bar{\lambda}_a C_{ai} C^\dagger_{ib} \bar{\lambda}^*_b \right] \left[-\frac{1}{2} \log \frac{m^2_{2i}}{M^2_f} \right]$$

$$\Sigma_6 = \frac{1}{4\pi^2}\left[\bar{\lambda}_a A_{ai} C^\dagger_{ib} \bar{\lambda}^*_b\right]\left[2\log\frac{M^2_{Li}}{m^2_{1i}} - \frac{1}{2}\log\frac{m^2_{1i}}{M^2_f}\right] + (A \leftrightarrow C, m_1 \leftrightarrow m_2)$$

The function $f(x, y)$ is defined as:

$$f(x, y) = x\log\frac{x}{y} - x + y$$

The messenger fields $(f, \bar{f})$ have similar expressions except for a flavor index.

B. Soft Mass Terms from the Two-Loop Yukawa Mediation

The two-loop corrections dominate when $M_F = 0$ in the hidden sector. Now we calculate such two-loop corrections explicitly via the wave function renormalization method. We use the notation of Ref. [**?**]. The superpotential is given by:

$$W = \lambda_a X \bar{\phi}_{ai}\phi_{ai} + \lambda_a \phi_{ai}\bar{F}_i f + \bar{\lambda}_a \bar{\phi}_{ai} F_i \bar{f} + M_F F_i \bar{F}_i + M_f f \bar{f}$$

In the following calculations only one pair of secondary messenger fields is introduced, thus $\lambda_a = \lambda = \bar{\lambda}_a$. According to this method, we must calculate the threshold effect after integrating out the heavy messengers. The soft terms can be extracted from the wave functions of light fields that get renormalization from messengers, which are given by:

$$m^2_{\text{soft}}|_{Q=M_h} = \left.\frac{\partial(\Delta\gamma)}{\partial\log Q}\right|_{Q=M_h} F^2$$

The beta-function, anomalous dimension, and their change when crossing the messenger scale are given by:

$$\beta_{\lambda^2} = \frac{\lambda^4}{16\pi^2}(N+2), \quad \gamma_F = -\frac{\lambda^2}{16\pi^2}, \quad \Delta\gamma_F = -\frac{\lambda^2}{16\pi^2}\theta(Q - M_h)$$

With this quantity we obtain:

$$m^2_{F\bar{F}} = \frac{\lambda^2}{16\pi^2}\frac{F}{M_h}M_F, \quad m^2_F = 0$$

Similarly, we have:

$$m^2_{F,Y1} = \frac{\lambda^4}{(16\pi^2)^2}(N+2)\frac{|F|^2}{M^2_h}, \quad m^2_{F,Y2} = \frac{\lambda^2 g^2_h}{(16\pi^2)^2}\frac{|F|^2}{M^2_h}$$

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Note: Figure translations are in progress. See original paper for figures.