# Yukawa deflected gauge mediation

Z. Chacko*
*Department of Physics, University of California, Berkeley, California 94720*
*and Theoretical Physics Group, Lawrence Berkeley National Laboratory, Berkeley, California 94720*

Eduardo Pontón†
*Department of Physics, Yale University, New Haven, Connecticut 06511*
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We consider models that are natural extensions of those where supersymmetry is broken at low energy scales and transmitted to visible matter by gauge interactions. We investigate the situation where the quark and lepton superfields of the minimal supersymmetric standard model (MSSM) are localized to a brane in a higher dimensional space while the messenger fields and the sector that breaks supersymmetry dynamically are localized to another brane in the same space. The MSSM gauge and Higgs fields are assumed to propagate in the bulk. If some of the messenger fields and the Higgs fields have the same quantum numbers, this allows the possibility of mixing between these fields so that the physical Higgs and messenger fields are admixtures of the brane and bulk fields. This manifests itself in direct couplings of the quark and lepton fields to the physical messengers which are proportional to the MSSM Yukawa couplings and hence preserve the flavor structure of the Cabibbo-Kobayashi-Masakawa (CKM) matrix. The result is new contributions to the soft supersymmetry breaking parameters that are related to the Yukawa couplings and which therefore naturally satisfy the constraints from flavor changing neutral currents. For messenger scales greater than 1000 TeV these new contributions are parametrically of the same order of magnitude as gauge mediation. This scenario naturally avoids the cosmological problems associated with stable messengers and admits a simple and natural solution to the $\mu$ problem based on the next-to-minimal supersymmetric standard model.

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## I. INTRODUCTION

Gauge mediated supersymmetry breaking is arguably the most attractive candidate for a realistic mechanism of supersymmetry breaking [1]. In this scenario one assumes that there is a hidden sector in which supersymmetry is broken, and which couples to a set of messenger fields charged under the standard model gauge interactions. Supersymmetry breaking effects are then communicated to the visible sector fields through loop effects involving the gauge interactions. This leads to a viable and highly predictive spectrum of sparticles. Since supersymmetry breaking is communicated by gauge interactions the squark and slepton spectrum is nearly flavor diagonal and therefore in good agreement with the experimental constraints on flavor changing neutral currents (FCNC's).

In gauge mediation it is usually assumed that direct interactions between the minimal supersymmetric standard model (MSSM) fields and the messenger fields, if any, are very small since this would lead to new sources of flavor violation beyond the Cabibbo-Kobayashi-Masakawa (CKM) matrix [3,4]. The current constraints on flavor changing neutral currents place tight constraints on any such interactions [5]:

$$\frac{m_{ds}^2}{m_{ss}^2} \leqslant (6\times 10^{-3})\,\frac{m_{ss}}{1\ \text{TeV}}. \tag{1.1}$$

In this paper we consider a natural extension of gauge mediation with messenger-Higgs field mixing in which there are no sources of flavor violation apart from the CKM matrix itself. We consider the situation where the quark and lepton superfields of the MSSM are localized to a brane in a higher dimensional space, while the gauge and Higgs fields propagate in the bulk. The messenger fields and the supersymmetry breaking sector are assumed to be localized to another brane in the same space. If some of the messenger fields have the same quantum numbers as the Higgs fields, this allows the possibility of mixing between them so that the physical Higgs and messenger fields are admixtures of the brane and bulk fields. This manifests itself in the Lagrangian as direct couplings of the quark and lepton fields to the physical messengers that are proportional to the MSSM Yukawa couplings and therefore preserve the flavor structure of the CKM matrix. The result is new contributions to the soft scalar masses that are related to the Yukawa couplings of the standard model fermions, and which therefore naturally satisfy the constraints from FCNC's. The extra dimensions are assumed to be sufficiently small that four dimensional gauge coupling unification is unaffected. This also allows other potentially large sources of supersymmetry breaking such as anomaly mediation [6], gaugino mediation [7], and radion mediation [8] to be neglected. This scenario, which we call "Yukawa deflected gauge mediation," naturally avoids the cosmological problems associated with stable messengers. We further investigate the $\mu$ problem of Yukawa deflected gauge mediation in the context of the next-to-minimal supersymmetric standard model (NMSSM). In the context of a

*Electronic address: zchacko@thsrv.lbl.gov
†Electronic address: eduardo.ponton@yale.edu

 

specific model, we demonstrate that it is indeed possible to generate the correct pattern of symmetry breaking with a realistic spectrum of masses.

Our idea is in the spirit of an earlier suggestion by Dvali and Shifman [9] that the Higgs doublets of the MSSM are in fact also the messengers of gauge mediated supersymmetry breaking. In that case there are also contributions to the scalar masses related to the Yukawa couplings and constrained by the CKM matrix, but obtaining a light Higgs doublet is not simple.

## II. MESSENGER-HIGGS FIELD MIXING

Consider a gauge mediated supersymmetry (SUSY) breaking model with two pairs of messengers $(Q_m, \overline{Q}_m)$ $(m=1,2)$ transforming as $\mathbf{5}$ and $\overline{\mathbf{5}}$ of SU(5). Under $SU(3)\times SU(2)\times U(1)_Y$, these decompose as $(\mathbf{3},\mathbf{1},-\frac{2}{3})\oplus(\mathbf{1},\mathbf{2},-\mathbf{1})$ and $(\overline{\mathbf{3}},\mathbf{1},\frac{2}{3})\oplus(\mathbf{1},\mathbf{2},\mathbf{1})$. The SU(2) doublets in the messengers which we denote by $Q_{iu}$ and $\overline{Q}_{id}$ have the same quantum numbers as the MSSM Higgs fields $\tilde{H}_u$ and $\tilde{H}_d$ and can therefore mix with them. They can also have additional Yukawa couplings to the MSSM quarks $q_i$ and $u_i^c$ and leptons $l_i$ and $e_i^c$. In general, these new Yukawa couplings will lead to additional flavor violation outside the CKM matrix, and must therefore be forbidden by a symmetry, such as the messenger number symmetry which exists if the only messenger coupling in the superpotential is $XQ\overline{Q}$. Here $\langle X\rangle = M+F\theta^2$ is a chiral superfield that parametrizes supersymmetry breaking.

We now consider the situation where the MSSM quarks and leptons are localized to a brane in a five dimensional space, while the MSSM gauge and Higgs fields live in the bulk of the space. The messenger fields and the sector that breaks supersymmetry dynamically [2] are assumed to live on another brane. The extra dimension is assumed to be sufficiently small that $1/r \geqslant M_{\rm GUT}$ and gauge coupling unification goes through exactly as in four dimensions. However, $r$ is assumed to be sufficiently larger than the inverse cutoff of the higher dimensional theory so that the exchange of massive bulk states with mass of order of the cutoff does not alter our conclusions about the form of the effective theory below the scale $1/r$. Here we assume that there are no other light bulk fields beyond those of supergravity and the MSSM gauge and Higgs fields.

A 5D gauge multiplet consists of the gauge field $A_M$ $(M=0,...,4)$, a real adjoint scalar $\sigma$, and a fermion $\lambda$. We assume that the fifth dimension is compactified on an $S^1/Z_2$ orbifold of radius $r$. The fixed points of the orbifold are "branes" on which the hidden and visible sectors can be localized. The $Z_2$ parity assignments of the gauge field are such that $A_5, \sigma$, and half of the $\lambda$ components are odd. These states will then get masses of order $1/r$, and the surviving degrees of freedom make up an $\mathcal{N}=1$ gauge multiplet (see, e.g., Refs. [10,11] for details). A 5D hypermultiplet consists of two $N=1$ chiral multiplets, one of which is necessarily even and the other odd under the orbifold. Once again the odd states are projected out and are not present in the effective theory below the scale $1/r$. Therefore the Higgs doublets of the MSSM $\tilde{H}_u$ and $\tilde{H}_d$ are assumed to emerge from two different hypermultiplets.

As a consequence of the higher dimensional nature of the theory any Yukawa couplings between the messenger fields and the MSSM quarks and leptons are forbidden by locality. However, mixing between the messengers and the Higgs fields is still allowed. After integrating out the extra dimension the superpotential of the higher dimensional theory has the form

$$\begin{aligned} W = {} & X\left[\sum_{m=1}^{2}\lambda_m Q_m\overline{Q}_m + \tilde{\lambda}_d Q_{1u}\tilde{H}_d + \tilde{\lambda}_u \tilde{H}_u \overline{Q}_{2d}\right] \\ & + [\tilde{y}_{U,ij}\tilde{H}_u q_i u_j^c + \tilde{y}_{D,ij}\tilde{H}_d q_i d_j^c + \tilde{y}_{L,ij}\tilde{H}_d l_i e_j^c] \\ & + [\mu \tilde{H}_u \tilde{H}_d + \text{gauge kinetic terms}], \end{aligned} \tag{2.1}$$

where in order to avoid the dangerous term $X\tilde{H}_u\tilde{H}_d$ we have imposed the discrete symmetry $X\to -X$, $Q_{1u}\to -Q_{1u}$, $\overline{Q}_{2d}\to -\overline{Q}_{2d}$, with all other fields neutral.

From this expression it is clear that the physical doublet messengers are

$$\overline{M}_{1d} = \frac{\tilde{\lambda}_d\tilde{H}_d + \lambda_1\overline{Q}_{1d}}{\sqrt{\tilde{\lambda}_d^2+\lambda_1^2}}, \tag{2.2}$$

$$M_{2u} = \frac{\tilde{\lambda}_u\tilde{H}_u + \lambda_2 Q_{2u}}{\sqrt{\tilde{\lambda}_u^2+\lambda_2^2}}, \tag{2.3}$$

while the physical Higgs fields $H_u$ and $H_d$ are the orthogonal linear combinations. The superpotential rewritten in terms of these fields takes the form

$$\begin{aligned} W = {} & X\left[\sum_{m=1}^{2}\lambda_m Q_{mT}\overline{Q}_{mT} + \lambda_1' Q_{1u}\overline{M}_{1d} + \lambda_2' M_{2u}\overline{Q}_{2d}\right] \\ & + [y_{U,ij}H_u q_i u_j^c + y_{D,ij}H_d q_i d_j^c + y_{L,ij}H_d l_i e_j^c] \\ & + [y'_{U,ij}M_{2u} q_i u_j^c + y'_{D,ij}\overline{M}_{1d} q_i d_j^c + y'_{L,ij}\overline{M}_{1d} l_i e_j^c] + \cdots \end{aligned} \tag{2.4}$$

where $Q_{mT}$ and $\overline{Q}_{mT}$ denote the messenger $SU(3)_C$ triplets. The new couplings $\lambda'$, $y$, and $y'$ are related in a straightforward way to the old couplings $\lambda$ and $\tilde{y}$. In particular, note that the ratios

$$\frac{y_{U,ij}}{y'_{U,ij}} = k_U, \tag{2.5}$$

$$\frac{y_{D,ij}}{y'_{D,ij}} = k_D \tag{2.6}$$

are independent of the indices $i$ and $j$. This implies that the Yukawa couplings of the messengers to matter are proportional to the MSSM Yukawa couplings. Therefore the new supersymmetry breaking effects that emerge from the direct messenger matter couplings will be constrained by the CKM matrix and the sizes of the Yukawa couplings and will not give rise to large flavor violation.

Since the messenger doublets now have direct renormalizable couplings to the visible sector fields they are no longer stable and can directly decay into them. However, one may worry that this is not true of the the messenger triplets and that these will be stable, leading to cosmological difficulties. However, if the theory emerges from a supersymmetric grand unified theory the messenger triplets mix with the Higgs triplets, which have direct couplings to matter. The Higgs triplets are integrated out at the grand unified theory (GUT) scale. Then in the effective theory below the Higgs triplet mass there are direct couplings of the messenger triplets to visible fields suppressed by powers of the Higgs triplet mass. While these couplings are renormalizable and dimensionless they are small, of order $M/M_{\rm GUT}$. Nevertheless, they are easily large enough to allow the triplets to decay sufficiently rapidly so as to avoid the cosmological problems associated with stable messengers.

We now attempt to determine the size of the supersymmetry breaking contributions from these new direct messenger-matter interactions. The one loop contributions to the $(\text{scalar mass})^2$ from the Yukawa type couplings to the messengers vanish to leading order in $(F/M)^2$ [3]. The subleading one loop contributions of order $[y'^2/(16\pi^2)](F^4/M^6)$ are smaller than the leading two loop contributions which are of order $[y'^2/(16\pi^2)][g_3^2/(16\pi^2)]\times(F/M)^2$ provided $(F/M^2)\leqslant g_3/(4\pi)$. Notice that the two loop contributions are always parametrically of the same order as the usual gauge mediated contributions. Hence we will concentrate on the case where the messenger scale is large, $M\gg 10^6$ GeV, when the one loop contributions to the scalar masses can be safely neglected. Other contributions to the soft terms are trilinear $A$ terms which arise at one loop.

In the next section we give a derivation of the most general two loop contributions to the soft masses and the one loop contributions to the $A$ terms, at the messenger scale. Below we give the expressions for the model Eq. (2.4) keeping only the Yukawa couplings for third generation particles.

With the notation $y_t\equiv y_{U,33}$, $y_b\equiv y_{D,33}$, $y_\tau\equiv y_{L,33}$ and similarly for the new, primed Yukawa couplings, we find the following expressions for the new contributions to the soft masses at the messenger scale:

$$\Delta m^2_{\tilde q_3}=\frac{1}{128\pi^4}\Bigg[y_b'^2\bigg(-\frac{8}{3}g_3^2-\frac{3}{2}g_2^2-\frac{14}{9}g_Y^2+y_t'^2+3y_b'^2+\frac{1}{2}y_\tau'^2\bigg)+y_t'^2\bigg(-\frac{8}{3}g_3^2-\frac{3}{2}g_2^2-\frac{26}{9}g_Y^2+3y_t'^2\bigg)\Bigg]\left|\frac{F}{M}\right|^2,$$

$$\Delta m^2_{\tilde t^c}=\frac{1}{128\pi^4}\Bigg[y_t'^2\bigg(-\frac{16}{3}g_3^2-3g_2^2-\frac{52}{9}g_Y^2+6y_t'^2+y_b^2+y_b'^2\bigg)-y_b'^2y_t^2\Bigg]\left|\frac{F}{M}\right|^2,$$

$$\Delta m^2_{\tilde b^c}=\frac{1}{128\pi^4}\Bigg[y_b'^2\bigg(-\frac{16}{3}g_3^2-3g_2^2-\frac{28}{9}g_Y^2+y_t^2+y_t'^2+6y_b'^2+y_\tau'^2\bigg)-y_t'^2y_b^2\Bigg]\left|\frac{F}{M}\right|^2, \tag{2.7}$$

$$\Delta m^2_{\tilde l_3}=\frac{y_\tau'^2}{128\pi^4}\bigg(-\frac{3}{2}g_2^2-6g_Y^2+\frac{3}{2}y_b'^2+2y_\tau'^2\bigg)\left|\frac{F}{M}\right|^2,$$

$$\Delta m^2_{\tilde\tau^c}=2\Delta m^2_{L_i},$$

$$\Delta m^2_{H_u}=\frac{3y_t^2}{256\pi^4}[6y_t'^2+y_b'^2]\left|\frac{F}{M}\right|^2,$$

$$\Delta m^2_{H_d}=-\frac{1}{256\pi^4}[3y_b^2(y_t'^2+3y_b'^2)+3y_\tau^2y_\tau'^2+(3y_by_b'+y_\tau y_\tau')^2]\left|\frac{F}{M}\right|^2.$$

Here $g_Y$ is the hypercharge gauge coupling where the hypercharge is defined by $Q=T_3+Y/2$, and $T_3$ is the third SU(2) generator. These have to be added to the well known gauge mediated expressions

$$\begin{aligned}
m^2_{\tilde q_3}&=\frac{N}{128\pi^4}\bigg(\frac{20}{27}g_Y^4+\frac{3}{4}g_2^4+\frac{4}{3}g_3^4\bigg),\\
m^2_{\tilde t^c}&=\frac{N}{128\pi^4}\bigg(\frac{320}{27}g_Y^4+\frac{4}{3}g_3^4\bigg),\\
m^2_{\tilde b^c}&=\frac{N}{128\pi^4}\bigg(\frac{80}{27}g_Y^4+\frac{4}{3}g_3^4\bigg),\\
m^2_{\tilde l_3}&=\frac{N}{128\pi^4}\bigg(\frac{20}{3}g_Y^4+\frac{3}{4}g_2^4\bigg),\\
m^2_{\tilde\tau^c}&=\frac{N}{128\pi^4}\bigg(\frac{80}{3}g_Y^4\bigg),\\
m^2_{Hu,Hd}&=\frac{N}{128\pi^4}\bigg(\frac{20}{3}g_Y^4+\frac{3}{4}g_2^4\bigg),
\end{aligned} \tag{2.8}$$

where $N$ is the number of $\mathbf{5}\oplus\overline{\mathbf{5}}$ messenger pairs. We also find the following one loop contributions to the $A$ terms:

$$\begin{aligned}
A_t&=\frac{y_t}{16\pi^2}(3y_t'^2+y_b'^2)\frac{F}{M},\\
A_b&=\frac{y_b}{16\pi^2}(y_t'^2+3y_b'^2)\frac{F}{M},\\
A_\tau&=\frac{3y_\tau y_\tau'^2}{16\pi^2}\frac{F}{M}.
\end{aligned} \tag{2.9}$$

The expressions above show that the new contributions to the scalar masses are comparable to those from gauge mediation for the up sector of the third generation. This is also true for the down sector if $\tan\beta$ is large. Even for $\tan\beta \approx 10$ and $y_b \approx y_b'$ Yukawa deflection gives a 10% correction to the $(\text{mass})^2$ of the right handed bottom squark at the messenger scale.

## III. DERIVATION OF THE SOFT TERMS

In this section we derive the general expressions for the soft supersymmetry breaking terms induced at the messenger scale. These results can then be applied to theories with matter-messenger couplings like the models we are considering. The general formulas are most easily derived by the method of analytical continuation into superspace developed in [12]. We start by reminding the reader of the basic idea. If supersymmetry breaking in the messenger sector is parametrized by the vacuum expectation value (VEV) of a chiral superfield $\langle X\rangle = M + F\theta^2$, then the leading supersymmetry breaking contribution to the observable sector, in an expansion in powers of $F/M^2$, can be described within a supersymmetric framework. More precisely, if the parameters of the theory at a scale $\Lambda_{UV}$ above the messenger scale $M$ are fixed, then the low energy values of the wave function renormalization constants will depend, through their renormalization group (RG) evolution, on the scale $M$ at which the messenger fields are integrated out. The soft supersymmetry breaking parameters can then be incorporated by the replacement $M \to |X|$ in the Kähler potential, that is, by analytical continuation into superspace (in holomorphic terms the correct analytical continuation is given by $M \to X$). If the observable sector superfields are denoted by $Q'^a$, the low energy Lagrangian, in the presence of soft supersymmetry breaking, can then be written as

$$\mathcal{L} = \int d^4\theta\, Q'^{\dagger}_a \mathcal{Z}(|X|)^a_b Q'^b + \left( \int d^2\theta\, \lambda'_{abc} Q'^a Q'^b Q'^c + \text{H.c.} \right). \tag{3.1}$$

For simplicity, here we chose to show only the Yukawa couplings in the superpotential. The generalization to other operators will be evident in what follows. Note also that we allow for off-diagonal mixing in the kinetic terms, so that $\mathcal{Z}(|X|)$ is a general Hermitian matrix. The soft supersymmetry breaking terms can be read from the Lagrangian (3.1) after replacing $X$ by its VEV and expanding in powers of $\theta$: $\mathcal{Z}(|X|) = Z + 1/2(\partial Z/\partial M)(F\theta^2 + F^\dagger \bar{\theta}^2) + 1/4(\partial^2 Z/\partial M^2) F F^\dagger \theta^2 \bar{\theta}^2$, where $Z = \mathcal{Z}(M)$ is the usual wave function renormalization constant. To display these terms more clearly, it is convenient to perform the following (chiral) field redefinition $Q = Z^{1/2}[1 + Z^{-1}(\partial Z/\partial M) F\theta^2] Q'$, after which the Lagrangian becomes

$$\mathcal{L} = \int d^4\theta\, Q^\dagger_a Q^a + \left( \int d^2\theta\, \lambda_{abc} Q^a Q^b Q^c + \text{H.c.} \right) - \tilde{Q}^\dagger_a (m^2_{\tilde{Q}})^a_b \tilde{Q}^b - (A_{abc} \tilde{Q}^a \tilde{Q}^b \tilde{Q}^c + \text{H.c.}). \tag{3.2}$$

Here $\tilde{Q}$ is the scalar component of $Q$, $\lambda_{abc} = \lambda'_{a'b'c'} (Z^{-1/2})^{a'}_a (Z^{-1/2})^{b'}_b (Z^{-1/2})^{c'}_c$ are the renormalized Yukawa couplings, and the soft masses are given by

$$m^2_{\tilde{Q}} = -\frac{1}{4} Z^{-1/2} \left( \frac{\partial^2 Z}{\partial \ln M^2} - \frac{\partial Z}{\partial \ln M} Z^{-1} \frac{\partial Z}{\partial \ln M} \right) \times Z^{-1/2} \frac{F F^\dagger}{M M^\dagger} \tag{3.3}$$

while the $A$ terms are given by

$$A_{abc} = \frac{1}{2} \left( \lambda_{a'bc} \left[ Z^{-1/2} \frac{\partial Z}{\partial \ln M} Z^{-1/2} \right]^{a'}_a + \lambda_{ab'c} \left[ Z^{-1/2} \frac{\partial Z}{\partial \ln M} Z^{-1/2} \right]^{b'}_b + \lambda_{abc'} \left[ Z^{-1/2} \frac{\partial Z}{\partial \ln M} Z^{-1/2} \right]^{c'}_c \right) \frac{F}{M}. \tag{3.4}$$

In order to find explicit expressions for the soft parameters (3.3) and (3.4) at a scale $\mu$ one needs to solve for $Z(\mu;M)$ from its RG evolution equation

$$\frac{dZ}{dt} = \gamma Z, \tag{3.5}$$

where $\gamma$ is the matrix of anomalous dimensions and $t = \ln\mu$. In general, it is not possible to find closed expressions for the soft supersymmetry breaking parameters even at lowest loop order, except in a few simple cases [12]. It is, however, possible to write closed expressions for the soft parameters at the scale $\mu = M$, which can then be used as initial data for a numerical solution to the RG equations below the messenger scale. This is the strategy that we will follow, and our next task is to find the general formulas for the soft parameters just below the messenger scale. From Eqs. (3.3) and (3.4), we see that we need to evaluate the first and second derivatives of $Z(\mu;M)$ with respect to $\ln M$. In order to do this, we first note that by rescaling the fields, we can conveniently set $Z(\Lambda_{UV}) = 1$. Writing then $Z = 1 + \delta Z$ at an arbitrary scale and integrating Eq. (3.5), we formally obtain, for scales $\mu < M$,

$$\delta Z(t;M) = \int_{\ln\Lambda_{UV}}^{\ln M} dt'\, \gamma_>(t')[1 + \delta Z(t')] + \int_{\ln M}^{t} dt'\, \gamma_<(t';M)[1 + \delta Z(t';M)].$$

In writing this expression we took into account the fact that the anomalous dimensions can be discontinuous at $\mu = M$, and denoted by $\gamma_>$ ($\gamma_<$) the anomalous dimensions above (below) $M$. Our notation also reflects the fact that the anomalous dimensions as well as $\delta Z$ can depend on $M$ only *below* the messenger scale. Differentiating once with respect to $\ln M$, we find

$$\frac{d\,\delta Z(t;M)}{d\ln M}=\Delta\gamma(M)[1+\delta Z(M)]+\int_{\ln M}^{t}dt'\left\{\frac{d\gamma_<(t';M)}{d\ln M}[1+\delta Z(t';M)]+\gamma_<(t';M)\frac{d\,\delta Z(t';M)}{d\ln M}\right\},\tag{3.6}$$

where $\Delta_\gamma(M)\equiv\gamma_>(M)-\gamma_<(M)$ and we defined $\gamma_<(M)\equiv\gamma_<(t=\ln M;M)$. To obtain Eq. (3.6) we also used the fact that $\delta Z$ is continuous across $M$ (the anomalous dimensions are finite) so that $\delta Z(M)=\int_{\ln\Lambda_{UV}}^{\ln M}dt'\,\gamma_>(t')[1+\delta Z(t')]$ is well defined. Taking now a second derivative and evaluating at $\mu=M$ we obtain

$$\left.\frac{d^2\delta Z(t;M)}{d\ln M^2}\right|_{t=\ln M}=\frac{d\Delta\gamma(M)}{d\ln M}[1+\delta Z(M)]+\Delta\gamma(M)\frac{d\,\delta Z(M)}{d\ln M}-\frac{d\gamma_<(t;M)}{d\ln M}\times[1+\delta Z(M)]-\gamma_<(M)\left.\frac{d\,\delta Z(t;M)}{d\ln M}\right|_{t=\ln M}.$$

In order to simplify this expression, we note that (in a mass-independent scheme) the anomalous dimensions depend on $\ln M$ only through the gauge and Yukawa couplings, *both* of which will be generically denoted by $\lambda$. This implies that $d\gamma/d\ln M$ is of two loop order and the terms proportional to $\delta Z(M)$ are three loop effects, which we will neglect. Using also $d\,\delta Z(M)/d\ln M=\gamma_>(M)$ and, from Eq. (3.6), $[d\,\delta Z(t;M)/d\ln M]|_{t=\ln M}\simeq\Delta_\gamma(M)$, we can write

$$\left.\frac{d^2\delta Z(t;M)}{d\ln M^2}\right|_{t=\ln M}=\sum_\lambda\frac{d\Delta\gamma(M)}{d\lambda(M)}\frac{d\lambda(M)}{d\ln M}+\Delta\gamma(M)\gamma_>(M)-\gamma_<(M)\Delta\gamma(M)-\sum_\lambda\frac{d\gamma_<(t;M)}{d\lambda(t;M)}\left.\frac{d\lambda(t;M)}{d\ln M}\right|_{t=\ln M}+(\text{three loop order}).$$

It only remains to evaluate $d\lambda(t;M)/d\ln M$, which we can do starting from the corresponding RG equation. If $\beta[\lambda]$ is the $\beta$ function for $\lambda$, we can formally write for $\mu<M$

$$\lambda(t;M)=\lambda(\Lambda_{UV})+\int_{\ln\Lambda_{UV}}^{\ln M}dt'\,\beta_>[\lambda(t')]+\int_{\ln M}^{t}dt'\,\beta_<[\lambda(t';M)].$$

Differentiating with respect to $\ln M$ and evaluating at $\mu=M$, we get

$$\left.\frac{d\lambda(t;M)}{d\ln M}\right|_{t=\ln M}=\Delta\beta[\lambda(M)],\tag{3.7}$$

where $\Delta\beta[\lambda(M)]\equiv\beta_>[\lambda(M)]-\beta_<[\lambda(M)]$ and $\lambda(M)=\lambda(\Lambda_{UV})+\int_{\ln\Lambda_{UV}}^{\ln M}dt'\,\beta_>[\lambda(t';M)]$. From the expression for $\lambda(M)$ we also see that $d\lambda(M)/d\ln M=\beta_>[\lambda(M)]$. The second derivative can then be put in the following form:

$$\left.\frac{d^2\delta Z(t;M)}{d\ln M^2}\right|_{t=\ln M}=\sum_\lambda\left(\frac{d\Delta\gamma(M)}{d\lambda(M)}\beta_>[\lambda(M)]-\frac{d\gamma_<(M)}{d\lambda(M)}\Delta\beta[\lambda(M)]\right)+(\Delta\gamma(M))^2+[\gamma_>(M),\gamma_<(M)],\tag{3.8}$$

where $[A,B]=AB-BA$ is a commutator. We now have all the ingredients required to evaluate the soft parameters at the messenger scale. To lowest loop order we can replace all factors of $Z$ by 1 in Eqs. (3.3) and (3.4). Then using Eqs. (3.6) and (3.8) (evaluated at $\mu=M$) we obtain the final two loop expressions for the soft masses. In matrix notation these are

$$m_{\tilde Q}^2|_{\mu=M}=-\frac{1}{4}\left\{\sum_\lambda\left(\frac{d\Delta\gamma}{d\lambda}\beta_>[\lambda]-\frac{d\gamma_<}{d\lambda}\Delta\beta[\lambda]\right)+[\gamma_>,\gamma_<]\right\}\bigg|_{\mu=M}\frac{FF^\dagger}{MM^\dagger}.\tag{3.9}$$

For the $A$ terms, we obtain from Eqs. (3.4) and (3.6) the one loop result

$$A_{abc}|_{\mu=M}=\frac{1}{2}\left(\lambda_{a'bc}\Delta\gamma_a^{a'}+\lambda_{ab'c}\Delta\gamma_b^{b'}+\lambda_{abc'}\Delta\gamma_c^{c'}\right)\bigg|_{\mu=M}\frac{F}{M}.\tag{3.10}$$

Equations (3.9) and (3.10) are the main results of this section. These equations are understood to hold just below the messenger scale. In particular, the sums in Eq. (3.10) run only over the couplings in the effective low energy theory. Given a specific model it is now straightforward to calculate the induced soft terms at the messenger scale. Note that in the absence of direct matter-messenger couplings the anomalous dimensions of the observable fields are continuous at $\mu=M$. In this case only the second (and third) terms in Eq. (3.9) survive and one recovers the standard gauge mediated results when $\lambda$ is a gauge coupling.

## IV. THE $\mu$ PROBLEM

The models of Yukawa deflected gauge mediation naturally satisfy all constraints coming from neutral flavor changing processes, which could arguably be considered the most difficult challenge in theories of supersymmetry breaking. A second issue that should be addressed in any model of supersymmetry breaking is the origin of the Higgs bilinear term in the superpotential [13–19]

$$W=\mu H_u H_d . \tag{4.1}$$

In its most basic form the difficulty arises because, for phenomenological reasons, $\mu$ should be of the order of the weak scale. This scale is in turn related to the scale of supersymmetry breaking (if the hierarchy problem is to be solved by supersymmetry) and there is *a priori* no reason that the supersymmetric term (4.1) should be of weak scale order. It is then natural to assume that the $\mu$ term vanishes at the tree level and is generated only after supersymmetry breaking, for example from Kähler terms like [13]

$$K=\lambda H_u H_d\left(\frac{X^\dagger}{M}+\frac{XX^\dagger}{M^2}+\cdots\right), \tag{4.2}$$

where $\langle X\rangle = M+F\theta^2$. After supersymmetry breaking the first term in Eq. (4.2) generates the $\mu$ term (4.1) while the second generates the supersymmetry breaking term

$$V=B\mu H_u H_d . \tag{4.3}$$

As has been stressed in [14], in theories of gauge mediation the real challenge is to explain why $B$ and $\mu$ are of the same order. Since all other soft masses are generated at one loop, one needs $\lambda \sim 1/16\pi^2$ in order that $\mu \sim (1/16\pi^2)F/M$ has the correct size. The problem is then that $B\mu \sim (1/16\pi^2)(F/M)^2$ which implies the relation $B \sim (16\pi^2)\mu$. Indeed, generically both the $\mu$ and $B\mu$ terms are generated at the same loop order, which results in the previous relation. Such a large value of $B$ would require an unacceptable degree of fine-tuning to obtain a correct electroweak symmetry breaking pattern.

A very appealing solution to this problem is to introduce a new light standard model singlet field $S$ with superpotential couplings [20–22]

$$W=\lambda S H_d H_u - \frac{1}{3}\kappa S^3. \tag{4.4}$$

If supersymmetry breaking gives a negative mass squared to $S$, then in the process of electroweak symmetry breaking it will acquire a VEV and an effective $\mu = \lambda\langle S\rangle$ of the correct size will be generated. Similarly, the $B\mu$ term can arise from the $A$ term

$$V=A_\lambda S H_u H_d . \tag{4.5}$$

Unfortunately, in models of gauge mediation, both $m_s^2$ and $A_\lambda$ are very small and it has been shown that it is not possible to obtain a realistic symmetry breaking pattern [21]. In addition, there is always a light state associated with the spontaneous breaking of an approximate $R$ symmetry under which all superfields have $R$ charge 2/3. This symmetry is only broken by the term Eq. (4.5), which is very small in gauge mediation.

On the other hand, in models of Yukawa deflected mediation $A$ terms are generated at one loop as we have shown in Eq. (3.10) and thus they have the required order of magnitude to destroy the $R$ symmetry. Also, in these models $m_S^2$ can get a substantial negative contribution, which can lead to a sizable $S$ VEV. In this section we analyze the next-to-minimal supersymmetric standard model, defined by the replacement of the $\mu$ term in the MSSM by the superpotential Eq. (4.4), and show that it is possible to obtain realistic electroweak symmetry breaking.

We pause to note that in this model there is no supersymmetric $CP$ problem [23]. By redefining the phases of $S$ and $H_u$, we can assume without loss of generality that $\lambda$ and $\kappa$ are real. By rescaling $\overline{M}_{1d}$ and $M_{2u}$ in Eq. (2.4) we can assume that the proportionality constants $k_U$ and $k_D$ in Eq. (2.5) are real as well. We can also assume that the couplings $\lambda_i, \lambda_i'$ in the hidden brane are real by rotating the remaining messenger fields. Now all $CP$ phases will reside in the CKM matrix. In order to see this, one can rotate the matter superfields to the quark mass eigenbasis. By redefining the quark superfield phases, one can absorb, as usual, all but one of the CKM phases. In the quark mass eigenbasis the gauge symmetry is not explicit and, in particular, the Yukawa interactions between the matter and charged Higgs superfields are not flavor diagonal whereas those involving the neutral Higgs fields, by definition, are. However, the important point is that all Yukawa interactions can be written in terms of the physical CKM matrix and the real quark mass eigenvalues. Furthermore, since all field redefinitions are performed at the superfield level, there are no additional phases in any of the soft parameters. Therefore there is only one physical $CP$ violating phase.

In what follows we neglect for simplicity the $CP$ phase and assume that all parameters are real. We have the option of either restricting $S$ to a brane or allowing it to propagate in the bulk. Allowing $S$ to propagate in the bulk allows for a greater range of couplings, since it can now couple directly to the messenger triplets, as well as to the doublets.

In order to see the main features more easily, we will consider the case in which only $\tilde{H}_d$ and $S$ propagate in the bulk. Further, we will neglect the terms involving the smaller Yukawa couplings $y_b$ and $y_\tau$, as well as $y_b'$ and $y_\tau'$, which are proportional to them. If $\tan\beta$ is large, however, one should also include these couplings. The superpotential we consider at the messenger scale has the form

$$W=y_t H_u q_3 t^c - S H_u(\lambda H_d+\lambda_d \overline{M}_{1d}) - \frac{\kappa}{3}S^3. \tag{4.6}$$

While other couplings of $S$ are in principle allowed by symmetry, we are neglecting them here for purposes of simplicity. If such additional couplings are not large we do not expect them to significantly alter our conclusions. It is now straightforward to obtain the soft breaking terms that are induced after integrating out the messenger fields, from the general equations (3.9) and (3.10). We find, in addition to the standard gauge mediated contribution Eq. (2.8), the following nonvanishing new contributions to the soft masses of the observable fields, at the messenger scale:

$$\Delta m^2_{\tilde{q}_3} = -\frac{\lambda_d^2 y_t^2}{256\pi^4}\left|\frac{F}{M}\right|^2,$$

$$\Delta m^2_{\tilde{t}^c} = -\frac{\lambda_d^2 y_t^2}{128\pi^4}\left|\frac{F}{M}\right|^2,$$

$$\Delta m^2_{H_u} = \frac{\lambda_d^2}{128\pi^4}\left(2\lambda_d^2+\kappa^2-2g_Y^2-\frac{3}{2}g_2^2\right)\left|\frac{F}{M}\right|^2, \tag{4.7}$$

$$\Delta m^2_{H_d} = -\frac{\lambda_d^2\lambda^2}{64\pi^4}\left|\frac{F}{M}\right|^2,$$

$$\Delta m^2_{\tilde{S}} = -\frac{\lambda_d^2}{128\pi^4}(4g_Y^2+3g_2^2+4\kappa^2-4\lambda_d^2-3y_t^2)\left|\frac{F}{M}\right|^2.$$

We observe that $m_S^2$ can indeed be negative if $\kappa \sim 1$. In this region of parameter space $m^2_{H_u}$ also receives a positive contribution. Similarly, the nonvanishing one loop trilinear terms are

$$A_t = \frac{y_t\lambda_d^2}{16\pi^2}\frac{F}{M},$$

$$A_\lambda = \frac{3\lambda\lambda_d^2}{16\pi^2}\frac{F}{M}, \tag{4.8}$$

$$A_\kappa = \frac{\kappa\lambda_d^2}{8\pi^2}\frac{F}{M}.$$

These $A$ terms are defined in Eq. (3.2). In particular, we have not factored out the corresponding Yukawa coupling.

Equations (4.7) and (4.8) are all proportional to $\lambda_d$. Given the values of the various Yukawa couplings at the messenger scale as well as the supersymmetry breaking scale $F/M$, one can use the NMSSM renormalization group equations, to obtain the values of the various soft masses at the weak scale. As usual, the Higgs mass parameter is driven negative by the top Yukawa coupling and we find electroweak symmetry breaking minima for a large range of parameters. In order to reproduce the $Z$ boson mass $M_Z$, we require that the Higgs boson VEV's satisfy $v^2 \equiv v_u^2+v_d^2 = (174\ \mathrm{GeV})^2$. This fixes the overall scale. The minimization also determines $\tan\beta = v_u/v_d$, and one should try to adjust $y_t$ to reproduce $m_{\mathrm{top}} \sim 165$ GeV (the difference from the experimental value of about 175 GeV is attributed to QCD corrections). However, the fact that $y_t$ is attracted to its low energy quasifixed point leaves some freedom in the choice of $y_t$ at the messenger scale. This choice is important, however, in determining the evolution of various quantities such as $m^2_{H_u}$. In practice we take as arbitrary input parameters the values of the Yukawa couplings $y_t$, $\lambda$, $\kappa$, and $\lambda_d$ at the messenger scale as well as the messenger scale $M_{\mathrm{mess}}$.

We give two sample points in Table I. We used $g_Y = 0.1816$, $g_2 = 0.6486$, and $g_2 = 1.1005$ for the gauge coupling constants at the 1 TeV scale, and checked that the theory remains perturbative up to the GUT scale. The rest of the input parameters are given in the table, as well as the weak scale values for the various physical masses, which include the soft as well as the $D$-term contributions.

TABLE I. Sample points in parameter space for $N=1,2$ where $N$ is the number of $\mathbf{5}\oplus\overline{\mathbf{5}}$ messenger pairs. All masses are in GeV. $\alpha_3$ is the strong coupling constant and the sensitivity parameter is defined in the main text.

| | | $N=1$ | $N=2$ |
|---|---|---|---|
| Inputs | $M_{\mathrm{mess}}$ | $10^{11}$ | $10^{14}$ |
| | $y_t$ | 0.9 | 0.9 |
| | $\lambda$ | 0.15 | 0.2 |
| | $\kappa$ | 0.8 | 0.98 |
| | $\lambda_d$ | 0.748 | 0.994 |
| Neutralinos | $m_{\chi_1^0}$ | 108 | 132 |
| | $m_{\chi_2^0}$ | 165 | 179 |
| | $m_{\chi_3^0}$ | 173 | 208 |
| | $m_{\chi_4^0}$ | 315 | 380 |
| | $m_{\chi_5^0}$ | 1550 | 1410 |
| Charginos | $m_{\chi_1^\pm}$ | 135 | 155 |
| | $m_{\chi_2^\pm}$ | 315 | 382 |
| Higgs bosons | $\tan\beta$ | 7.1 | 5.5 |
| | $m_{h^0}$ | 115 | 115 |
| | $m_{H^0}$ | 467 | 500 |
| | $m_A$ | 466 | 500 |
| | $m_{H'^0}$ | 1220 | 1170 |
| | $m_{A'}$ | 1660 | 1360 |
| | $m_{H^\pm}$ | 473 | 505 |
| Sleptons | $m_{\tilde{e}_R}$ | 100 | 113 |
| | $m_{\tilde{e}_L}$ | 470 | 505 |
| | $m_{\tilde{\nu}_L}$ | 465 | 500 |
| Tops squarks | $m_{\tilde{t}_1}$ | 390 | 467 |
| | $m_{\tilde{t}_2}$ | 850 | 907 |
| Other squarks | $m_{\tilde{u}_L}$ | 1050 | 1105 |
| | $m_{\tilde{u}_R}$ | 1000 | 1065 |
| | $m_{\tilde{d}_L}$ | 1050 | 1105 |
| | $m_{\tilde{d}_R}$ | 960 | 1010 |
| Gluino | $M_3$ | 815 | 1030 |
| Sensitivity | $\alpha_3$ | 100 | 130 |
| | $\lambda_d$ | 90 | 130 |

We note that the next lightest supersymmetric particle (NLSP) is the right-handed stau (as in gauge mediation, the gravitino is the LSP). This is due to the effect of the $\mathrm{U}(1)_Y$ Fayet-Iliopoulos $D$ term [24] in the RG running of the soft masses:

$$\Delta\frac{d}{dt}m_i^2 = \frac{1}{16\pi^2}Y_i g_Y^2 \sum_j Y_j m_j^2, \tag{4.9}$$

where the sum runs over all fields and $Y_i$ is the hypercharge of the $i$th field. In pure gauge mediation this contribution vanishes, but it is in general not zero in the presence of Yukawa couplings. In our case we find, at the messenger scale,

$$\sum_j m_j^2 = \frac{\lambda_d^2}{32\pi^2}(3y_t^2 + 4\lambda^2 + 4\lambda_d^2 + 2\kappa^2 - 4g_Y^2 - 3g_2^2) \tag{4.10}$$

which is always positive since $y_t \sim 1$ cancels the smaller negative gauge contributions. Therefore, the fields having a positive (negative) hypercharge will receive a negative (positive) contribution from this term. The most important effect is on the right-handed sleptons and thus we expect the NLSP to correspond to the stau in this class of models.

A second distinctive feature is the relation $\Delta m_{t^c}^2 = 2\Delta m_{q_3}^2$, which holds, up to small corrections proportional to $y_b, y_\tau$, even when both $H_u$ and $H_d$ are allowed to propagate in the bulk.

When analyzing the spectrum at the weak scale it is important to include the radiative corrections to the lightest neutral Higgs boson mass [25]. The largest effect can be viewed as a top quark–top squark loop contribution to an effective quartic term in the effective potential below the top squark mass [26]. We include an estimate of this effect by adding the term

$$\Delta V_H = \left(\frac{3y_t^4}{8\pi^2}\ln\frac{m_{\tilde{t}}}{m_t}\right)(H_u^\dagger H_u)^2 \tag{4.11}$$

to the Higgs potential.

An important feature of these results is the amount of fine-tuning required to achieve electroweak symmetry breaking. We define the fractional sensitivity to a parameter $c$ (a coupling renormalized at $M_{\text{mess}}$) to be [27,28]

$$\text{sensitivity} = \frac{c}{v}\frac{\partial v}{\partial c}, \tag{4.12}$$

where $v$ is the Higgs boson VEV and the derivative is taken with all other couplings at the messenger scale held fixed. We find that the largest sensitivities are associated with $\alpha_3 = g_3^2/(4\pi)$ and $\lambda_d$. We note, however, that the sensitivities shown in the table are of the same order as the ones one would obtain for pure gauge mediation with tree level $\mu$ and $B_\mu$ terms fixed by the requirement of correct electroweak symmetry breaking (for the same values of $\tan\beta$ as shown in Table I). This amount of fine-tuning seems to be inherent in models in which the dominant soft breaking contributions arise from gauge mediation.

## V. CONCLUSIONS

Yukawa deflection alters the spectrum of gauge mediated supersymmetry breaking in a highly predictive manner while maintaining the requisite suppression of flavor changing neutral currents. It is an important effect for the third generation sparticles that have sizable Yukawa couplings. We have demonstrated that it can resolve in a simple and natural way the $\mu$ problem of gauge mediation, as well as the cosmological problems associated with stable messengers.

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