

## A simple model of low-scale direct gauge mediation

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**ABSTRACT:** We construct a calculable model of low-energy direct gauge mediation making use of the metastable supersymmetry breaking vacua recently discovered by Intriligator, Seiberg and Shih. The standard model gauge group is a subgroup of the global symmetries of the SUSY breaking sector and messengers play an essential role in dynamical SUSY breaking: they are composites of a confining gauge theory, and the holomorphic scalar messenger mass appears as a consequence of the confining dynamics. The SUSY breaking scale is around 100 TeV nevertheless the model is calculable. The minimal non-renormalizable coupling of the Higgs to the DSB sector leads in a simple way to a  $\mu$ -term, while the  $B$ -term arises at two-loop order resulting in a moderately large  $\tan \beta$ . A novel feature of this class of models is that some particles from the dynamical SUSY breaking sector may be accessible at the LHC.

**KEYWORDS:** Supersymmetry Breaking, Supersymmetry Phenomenology.

While supersymmetry elegantly solves the fine tuning problem of the Higgs mass, and may even explain the origin of the weak scale by relating it to the supersymmetry breaking scale, a generic supersymmetric extension of the standard model (SM) itself raises a number of problems. These problems include the  $\mu$ -problem (why the single dimensionful supersymmetric parameter is related to the supersymmetry breaking scale) and the little hierarchy problem (which is a percent level fine-tuning problem emerging from the non-observation of the Higgs and superpartners at LEP2).

One of the main issues that was appreciated early on in supersymmetric model building is the problem of flavor changing neutral currents (FCNCs): for generic soft supersymmetry breaking scalar masses there are additional one-loop diagrams without GIM suppression contributing to FCNC's. This problem is quite generic in models with high scale supersymmetry breaking, where non-trivial flavor physics is likely to affect the soft breaking scalar masses. This issue led to interest in gauge mediated SUSY breaking [1, 2] (GMSB), where the scale of supersymmetry breaking can be below the flavor breaking scale, and the soft masses themselves are generated via SM gauge interactions. As a result the soft breaking mass terms will only depend on the SM quantum numbers and be flavor independent. While many realistic models were constructed (see [3] for a review), they were quite complicated and typically had several layers of interactions (messengers) to communicate SUSY breaking to the Standard Model fields. Simplifying these models so that messengers would directly participate in the dynamics of the dynamical supersymmetry breaking (DSB) sector proved difficult. Even though viable direct gauge mediation models exist [4], they typically require rather large messenger scales. While these scales could be sufficiently low to provide significant theoretical control in studying the dynamics of the DSB sector, one of the main promises of gauge mediation — the possibility that in models with a low SUSY breaking scale the DSB sector itself could in principle be directly observable in future experiments — was never realized. DSB models without a hierarchy of scales are typically strongly coupled and as a result one can at best establish the existence of a SUSY breaking minimum but not the details of the spectrum.

In this paper we make use of the recent discovery by Intriligator, Seiberg, and Shih [5] (ISS) of metastable SUSY breaking vacua. From the model building point of view the main new feature of the models of [5] is that the supersymmetry breaking vacua are located near the origin of the moduli space yet are calculable. This raises the hope that a calculable low scale direct mediation model can be obtained. In this paper we show the first example of such a model. As in the ISS case the DSB sector of our model has a fairly simple dual description in the UV: it is just SUSY QCD with some masses and higher dimensional operators added. The higher dimensional operators can be suppressed by scales as high as  $10^{11}$  GeV. Supersymmetry breaking is triggered by dynamical symmetry breaking and while the SUSY breaking scale is as low as 100 TeV, the effective low energy theory is calculable. Since the fine-tuning depends logarithmically on SUSY breaking scale, it is significantly reduced in our model. Furthermore, the  $\mu$ -term could be generated by the same dynamics that leads to SUSY breaking.

**A simple ISS-type model.** We will start out with a simple toy model and gradually

add a few features in order to make a realistic SUSY breaking model. We intend to make use of the ISS models by embedding the SM gauge group into the flavor symmetry of the DSB sector. Thus the flavor symmetry should at least be SU(5). However since at least one field charged under the flavor symmetry gets a VEV, the minimal size of the flavor symmetry in the DSB sector is SU(6). To focus on the simplest possibility we assume that there is no gauge group in the magnetic description. Thus we are led to consider the following fields charged under the global symmetries

|                      | SU(6)         | U(1) | U(1) <sub>R</sub> |
|----------------------|---------------|------|-------------------|
| $\tilde{\phi}$       | □             | 1    | 0                 |
| $\tilde{\bar{\phi}}$ | □̄            | -1   | 0                 |
| $\widetilde{M}$      | <b>Ad + 1</b> | 0    | 2                 |

(1)

with the superpotential

$$W_1 = \tilde{\bar{\phi}} \widetilde{M} \tilde{\phi} - hf^2 \text{Tr} \widetilde{M} . \tag{2}$$

The global symmetries of this model are just the symmetries of an s-confining SU(5) gauge theory with 6 massive flavors [6, 7]. Indeed, we can identify  $\widetilde{M}$ ,  $\tilde{\phi}$ , and  $\tilde{\bar{\phi}}$  with mesons, baryons, and antibaryons of the electric description respectively. The linear term in the superpotential above then arises from a mass term in the microscopic theory with the identification  $m\Lambda \approx hf^2$ , while the cubic term is required to ensure the correct mapping of the two descriptions. Finally instantons generate an operator

$$W_{\text{inst.}} = \frac{\det \widetilde{M}}{\Lambda^3} , \tag{3}$$

where  $\Lambda$  is the intrinsic holomorphic (dynamical) scale of the microscopic theory. As explained in ref. [5], the term in eq. (3) is responsible for ensuring that there is a SUSY preserving global minimum with

$$\langle \widetilde{M} \rangle \sim f \left( \frac{\Lambda}{f} \right)^{\frac{3}{5}} . \tag{4}$$

There is also a metastable SUSY breaking vacuum at  $\langle \widetilde{M} \rangle \sim 0$ , which can have a lifetime much longer than the age of the Universe for  $f \ll \Lambda$ . Near this SUSY breaking vacuum, the instanton term is an irrelevant operator that we can ignore (unless we want to calculate the tunnelling rate to the true vacuum state). SUSY is broken since the matrix  $\tilde{\bar{\phi}} \tilde{\phi}$  has rank one, so

$$\frac{\partial W_1}{\partial \widetilde{M}_I^J} = \tilde{\bar{\phi}}^I \tilde{\phi}_J - hf^2 \delta_J^I \neq 0 . \tag{5}$$

In order to minimize the scalar potential energy, one flavor (in an appropriate basis) will get a VEV,  $\tilde{\bar{\phi}}^K \tilde{\phi}_K = hf^2$ , and the global SU(6) symmetry will be spontaneously broken to SU(5).

**Gauging the flavor symmetry.** We will gauge the standard model (SM) subgroup of SU(5) in which case vacuum alignment [8] will prefer the VEV to align so as to preserve the gauge symmetry. Thus it is convenient to write the fields in a form where the unbroken symmetry is manifest, so we split  $\tilde{\phi}_J$  into  $\phi_j$  and  $\psi$  (where  $\langle\psi\rangle \neq 0$ ) and similarly

$$\widetilde{M} = \begin{pmatrix} M_i^j & N^j \\ \bar{N}_i & X \end{pmatrix}. \tag{6}$$

So our field content can be rewritten as<sup>1</sup>

|                   |           |                 |        |              |               |     |           |                 |
|-------------------|-----------|-----------------|--------|--------------|---------------|-----|-----------|-----------------|
|                   | $\phi$    | $\bar{\phi}$    | $\psi$ | $\bar{\psi}$ | $M$           | $X$ | $N$       | $\bar{N}$       |
| SU(5)             | $\square$ | $\bar{\square}$ | 1      | 1            | <b>Ad + 1</b> | 1   | $\square$ | $\bar{\square}$ |
| U(1) <sub>R</sub> | 0         | 0               | 0      | 0            | 2             | 2   | 2         | 2               |

(7)

with a superpotential

$$W_2 = \bar{\phi}M\phi + \bar{\psi}X\psi + \bar{\phi}N\psi + \bar{\psi}\bar{N}\phi - hf^2(\text{Tr}M + X). \tag{8}$$

Let us discuss the dynamics entailed by this superpotential. The equation of motion for  $X$  leads to a non-zero  $\bar{\psi}\psi$  VEV. This in turn marries  $\phi$  ( $\bar{\phi}$ ) with  $\bar{N}$  ( $N$ ) making sure that they are massive and do not obtain VEVs. Finally, SUSY is broken by the  $\mathcal{F}$ -component of  $\text{Tr}M$ . Rescaling  $\text{Tr}M$  so that its kinetic term is canonically normalized, we obtain

$$\mathcal{F}_{\text{Tr}M} = \sqrt{5}hf^2. \tag{9}$$

As pointed out in [5] there are a number of massless states at tree-level. Some of these are goldstone bosons of the spontaneously broken SU(6) symmetry. Since SU(6) can be explicitly broken by the superpotential and is certainly broken by gauging the SM subgroup, these fields can obtain masses at one-loop level once SUSY is broken. There are also massless scalars corresponding to pseudo-flat directions of the O’Raifeartaigh model. As shown in [5] these will also obtain soft masses at one-loop which can be easily analyzed using the Coleman-Weinberg potential. In particular, the field with the non-vanishing  $\mathcal{F}$ -term,  $\text{Tr}M$ , is stabilized at the origin. Finally, the fermionic components of  $M$  remain massless.

**Communicating SUSY breaking to the SM fields.** Let us now describe how SUSY breaking is communicated to the SM superpartners. The fields  $\bar{\phi}$  and  $\phi$  couple directly to the SUSY breaking and obtain holomorphic soft mass terms. In addition,  $\bar{N}$  and  $N$  will obtain holomorphic soft mass terms due to the supersymmetric mixing (through  $\langle\psi\rangle$ ) with  $\bar{\phi}$  and  $\phi$  in the superpotential. Once the SM subgroup of SU(5) is gauged these fields will act as messengers. It is important to notice that in the absence of a VEV for  $\text{Tr}M$  the mass matrix of scalar messengers would have one vanishing eigenvalue at tree level. This is problematic for two reasons. First, we cannot allow light scalars with SM charges. Second, messenger multiplets with a vanishing supertrace of the mass matrix and massless scalars

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<sup>1</sup>We only show quantum numbers under symmetries relevant for the following analysis.

do not usually contribute to soft scalar masses at two-loops (although non-trivial mixing of the messengers in our model modifies the calculation of gauge mediated contributions). There is one more problem in the model presented so far — at the minimum of the potential there is an accidental discrete  $R$ -symmetry which forbids gaugino masses [5]. In this case this is a  $Z_{10}$   $R$ -symmetry. This is the discrete subgroup of the  $U(1)_R$  in (7) that is left unbroken by (3). In fact, the same  $R$ -symmetry forbids the soft masses for fermionic components of  $M$ . Thus in order to solve these problems we need to generate a scalar VEV for the field  $M$ . Below we will show two possibilities for how to achieve that.

Once we have a VEV for  $\text{Tr}M$ , the fermion messenger mass matrix will have the see-saw form

$$m_f = \begin{pmatrix} \langle M \rangle & \langle \psi \rangle \\ \langle \bar{\psi} \rangle & 0 \end{pmatrix}. \tag{10}$$

A Majorana gaugino mass will now be generated at one-loop as in more standard gauge mediation models. Notice that to leading order in the SUSY breaking parameter  $\mathcal{F}$  the gaugino mass is proportional to  $\text{Tr}(m_f^{-1}\mathcal{F})$  and it vanishes in our model.<sup>2</sup> However, terms which are higher order in  $\mathcal{F}$  are non-vanishing (the leading  $\mathcal{F}^3/m_f^5$  term was calculated in [9] and the result to all orders in  $\mathcal{F}$  can be found in [10]). Since  $\mathcal{F}/m_f^2 \sim 1$  in our model these contributions are not suppressed. The masses of both the scalar and fermionic superpartners (while qualitatively similar) will differ from predictions of usual gauge mediation scenarios, due to a more complicated messenger spectrum.

Scalar components of  $M$  will obtain contributions to their masses both from the Coleman-Weinberg potential and the usual gauge mediated loops (except for two scalars that are neutral under the Standard Model). Finally, we note that the fermionic components of  $M$  will obtain masses both from gauge mediation and from Yukawa coupling to messengers,  $\bar{\phi}M\phi$ . Thus these particles may have TeV scale masses, and be accessible at the LHC. Since  $M$  contains  $SU(5)$  adjoints some of the scalars and fermions will be colored and thus can be easily produced. The details of the spectrum and experimental signatures will be studied elsewhere.

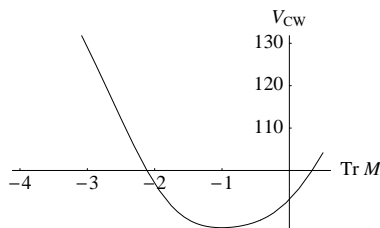
**Generating the  $M$ -VEV via singlet interactions.** We have seen above that in order to generate a non-vanishing gaugino mass we need to generate a scalar VEV for  $M$ . Here and in the next paragraph we show two possibilities for that. While the models seem quite similar the implications of the two mechanisms are actually quite different. In both cases we modify the theory slightly by adding the singlets  $S, \bar{S}, Z, \bar{Z}$ . In the first case these fields will be fundamental singlets and the discrete  $R$ -symmetry is broken via explicit superpotential interactions:

$$W = \bar{\phi}M\phi + \bar{\psi}X\psi + \bar{\phi}N\psi + \bar{\psi}\bar{N}\phi - hf^2(\text{Tr}M + X) + (d\text{Tr}M + m)S\bar{S} + m'(S\bar{Z} + Z\bar{S}). \tag{11}$$

The new superpotential terms necessarily break the  $R$ -symmetry, due to the simultaneous presence of both the  $\text{Tr}M S\bar{S}$  and the  $S\bar{S}$  terms. The last terms in this superpotential

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<sup>2</sup>We thank Hitoshi Murayama for pointing this out to us.



**Figure 1:** Plot of the Coleman-Weinberg potential (in arbitrary units). The horizontal axis is  $\text{Tr}M$  in units of  $\sqrt{\mathcal{F}}$ .

(proportional to  $m'$ ) ensure that supersymmetry is not restored via VEV's of  $S$  and  $\bar{S}$ . There are now two interactions contributing to the one-loop potential for  $M$ . Loops of the singlets  $S$  and  $\bar{S}$  will tend to generate an  $M$  VEV to cancel the mass term for these fields while loops of  $\bar{\phi}$  and  $\phi$  will generate a positive contribution to the mass squared of  $\text{Tr}M$ . Generically the vacuum will be shifted from the origin. In order to keep the  $S, \bar{S}$  fields from obtaining VEV's (so that SUSY breaking originates fully from the F-term of  $M$ , and not partly from the singlet sector) we need to assume  $m'^2 > dhf^2$ . The one-loop Coleman-Weinberg potential for spontaneously broken supersymmetric theories

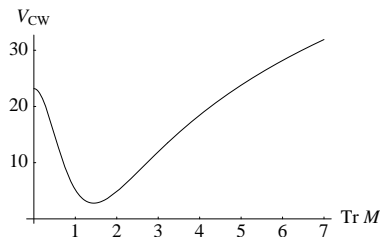
$$\frac{1}{64\pi^2} \text{STr} \mathcal{M}^4 \log \frac{\mathcal{M}^2}{\Lambda^2} \tag{12}$$

can be evaluated for the potential following from (11). In order for the minimum of the potential to be significantly shifted from the origin, the interaction strength of the  $S, \bar{S}$  fields to  $\text{Tr}M$  should not be very small, otherwise their effect around the origin will be negligible. As an example we show the Coleman-Weinberg potential along the  $\text{Tr}M$  direction in units of  $\sqrt{hf^2}$ , for parameters  $m = 4, m' = 1.5, d = 0.4$  (again in the same units). We can see that for these parameters the minimum is at  $\text{Tr}M \sim \sqrt{5hf^2}$ . This implies that messenger multiplets obtain supersymmetric contributions to their masses,  $m_m$ , which are comparable to splittings withing the multiplet,  $\mathcal{F}/m_m^2 \sim 1$ .

**Generating the  $M$ -VEV via gauge interactions.** A perhaps more elegant solution for generating the VEV for  $M$  is by using the mechanism of [11]. Instead of adding the explicit  $mS\bar{S}$  mass term one can maintain the discrete R-symmetry in the superpotential, and only break it spontaneously via the VEV of  $M$ . This has the added benefit that imposing this discrete R-symmetry can forbid some (but not all) other unwanted terms in the superpotential (for example  $\text{Tr}M^2$  which would restore supersymmetry).

To achieve the spontaneous breaking of the R-symmetry (by forcing the  $M$  VEV from the origin) we gauge a U(1) symmetry under which  $S, Z$  have charges +1 and  $\bar{S}, \bar{Z}$  have charges -1. The superpotential will now be

$$W = \bar{\phi}M\phi + \bar{\psi}X\psi + \bar{\phi}N\psi + \bar{\psi}\bar{N}\phi - hf^2(\text{Tr}M + X) + d\text{Tr}M\bar{S}\bar{S} + m'(S\bar{Z} + Z\bar{S}) . \tag{13}$$



**Figure 2:** Plot of the Coleman-Weinberg potential for the case with a U(1) gauge symmetry (in arbitrary units). The horizontal axis is  $\text{Tr} M$  in units of  $\sqrt{\mathcal{F}}$ .

In order for these U(1) gauge fields to contribute to the CW potential we need to pick the parameters of the theory such that  $S$  obtains a VEV. The reason for the additional contributions to the CW potential is that in this case the  $Z, \bar{Z}$  directions will be related to the  $M$  VEV. As a consequence SUSY breaking will not fully originate from the dynamical sector, but the  $S, Z$  sector will also contribute. Minimizing the CW potential one can find minima similar to those in [11] for a wide range of perturbative U(1) couplings  $g$  and small couplings  $d$ . An example for such a minimum is shown in figure 2 for  $d = .01$ ,  $g = 0.1$ , and  $m' = 0.05$  (in units of  $\sqrt{\mathcal{F}}$ ).

**Generating the SUSY breaking linear term via dynamics.** We now turn to the origin of the linear term which leads to SUSY breaking. It can be generated from a condensate of an additional supercolor sector (this possibility has recently been also emphasized in [12]). One of the simplest possibilities is an  $\text{SU}(2)_{\text{sc}}$  gauge group with 2 flavors. Thus our complete SUSY breaking model has the following fields

|                              | $\phi$    | $\bar{\phi}$    | $\psi$ | $\bar{\psi}$ | $M$                        | $X$ | $N$       | $\bar{N}$       | $S$ | $\bar{S}$ | $Z$ | $\bar{Z}$ | $p$       | $\bar{p}$       | $T$ |
|------------------------------|-----------|-----------------|--------|--------------|----------------------------|-----|-----------|-----------------|-----|-----------|-----|-----------|-----------|-----------------|-----|
| $\text{SU}(2)_{\text{sc}}$   | 1         | 1               | 1      | 1            | 1                          | 1   | 1         | 1               | 1   | 1         | 1   | 1         | $\square$ | $\bar{\square}$ | 1   |
| $\text{U}(1)_{\text{gauge}}$ | 0         | 0               | 0      | 0            | 0                          | 0   | 0         | 0               | 1   | -1        | 1   | -1        | 0         | 0               | 0   |
| $\text{SU}(5)$               | $\square$ | $\bar{\square}$ | 1      | 1            | $\mathbf{Ad} + \mathbf{1}$ | 1   | $\square$ | $\bar{\square}$ | 1   | 1         | 1   | 1         | 1         | 1               | 1   |
| $\text{SU}(2)$               | 1         | 1               | 1      | 1            | 1                          | 1   | 1         | 1               | 1   | 1         | 1   | 1         | $\square$ | $\bar{\square}$ | 1   |

(14)

The full superpotential of the model is

$$\begin{aligned}
 W = & \bar{\phi} M \phi + \bar{\psi} X \psi + \bar{\phi} N \psi + \bar{\psi} \bar{N} \phi - \frac{h}{\Lambda_{UV}^2} p^2 \bar{p}^2 (\text{Tr} M + X) \\
 & + c T p \bar{p} + d \text{Tr} M S \bar{S} + m' (S \bar{Z} + Z \bar{S}) .
 \end{aligned}
 \tag{15}$$

This superpotential can be enforced for example by the discrete  $Z_{10}$  R-symmetry under which  $M, X, N, \bar{N}, Z, \bar{Z}, T$  have charge 2, and  $p, \bar{p}, \phi, \bar{\phi}, \psi, \bar{\psi}, S, \bar{S}$  have charge 0. This is anomaly free under the  $\text{SU}(2)_{\text{sc}}$  and the  $\text{U}(1)_{\text{gauge}}$ . However, it is still not the most generic superpotential allowed by the symmetries. The term  $X S \bar{S}$  would also be allowed by the symmetries, but has to be assumed to vanish. However, other dangerous terms like  $\text{Tr} M^2$  are excluded by this discrete symmetry. The couplings to the singlet  $T$  give mass terms for the mesons  $p \bar{p}$ . Since the supercolor sector has a deformed quantum moduli space this

leaves the “baryons”  $B = p^2$  and  $B = \bar{p}^2$ , with non-zero VEVs. Thus the strong dynamics enforces

$$p^2 \bar{p}^2 = B \bar{B} = \Lambda_{\text{sc}}^4 = f^2 \Lambda_{UV}^2 \quad (16)$$

where  $\Lambda_{\text{sc}}$  is the intrinsic holomorphic scale for the supercolor gauge group, and the last part of the equation can be taken as the definition of the scale  $f$ .

**Generating the  $\mu$ -term.** To complete the construction of a phenomenologically relevant model we need to generate the terms in the Higgs potential necessary for electroweak symmetry breaking. To obtain a  $\mu$  term we add the superpotential term

$$W_\mu = \beta \frac{p^2 \bar{p}^2}{\Lambda_{UV}^3} H_u H_d . \quad (17)$$

Then around the scale where the supercolor group gets strong the model generates a  $\mu$  term of order

$$\mu \sim \beta f \left( \frac{\Lambda_{\text{sc}}}{\Lambda_{UV}} \right)^2 . \quad (18)$$

The soft SUSY-breaking  $B$ -term, however, vanishes at tree level. As noticed in [13] a  $B$ -term is however generated at two loops (since  $M_2$  itself is generated at one-loop order)

$$B \sim \frac{3\alpha_2}{2\pi} M_2 \mu \ln \left( \frac{\mathcal{F}}{M_2 \mu} \right) . \quad (19)$$

This leads to a  $B$ -term that is small compared to the square of electroweak scale (by a factor of  $\alpha_2$ ) and consequently results in a large  $\tan \beta$ , of order 10–50.

**The microscopic description of the theory.** We will assume that all the singlets  $S$ ,  $\bar{S}$ ,  $Z$ ,  $\bar{Z}$ , and  $T$  are elementary fields both in the effective and microscopic description. Above the scale  $\Lambda$  the microscopic dual description is a generalized SUSY QCD:

|           | SU(5) | SU(2) <sub>SC</sub> | U(1) <sub>gauge</sub> | SU(5) | SU(2) | SU(2) | Z <sub>10</sub> |         |
|-----------|-------|---------------------|-----------------------|-------|-------|-------|-----------------|---------|
| $Q$       | □     | 1                   | 0                     | □     | 1     | 1     | 1               | ), (20) |
| $\bar{Q}$ | □̄    | 1                   | 0                     | □̄    | 1     | 1     | 1               |         |
| $q$       | □     | 1                   | 0                     | 1     | 1     | 1     | 1               |         |
| $\bar{q}$ | □̄    | 1                   | 0                     | 1     | 1     | 1     | 1               |         |
| $p$       | 1     | □                   | 0                     | 1     | □     | 1     | 0               |         |
| $\bar{p}$ | 1     | □                   | 0                     | 1     | 1     | □     | 0               |         |
| $S$       | 1     | 1                   | 1                     | 1     | 1     | 1     | 0               |         |
| $\bar{S}$ | 1     | 1                   | −1                    | 1     | 1     | 1     | 0               |         |
| $Z$       | 1     | 1                   | 1                     | 1     | 1     | 1     | 2               |         |
| $\bar{Z}$ | 1     | 1                   | −1                    | 1     | 1     | 1     | 2               |         |
| $T$       | 1     | 1                   | 0                     | 1     | □     | □     | 2               |         |



with a superpotential

$$\begin{aligned}
 W = & m S\bar{S} + m'(S\bar{Z} + Z\bar{S}) + cT p\bar{p} + \beta \frac{p^2\bar{p}^2}{\Lambda_{UV}^3} H_u H_d \\
 & + \frac{\tilde{d}}{\Lambda_{UV}} \text{Tr} Q\bar{Q} S\bar{S} - \frac{\tilde{h}}{\Lambda_{UV}^3} p^2\bar{p}^2 (\text{Tr} Q\bar{Q} + q\bar{q}) .
 \end{aligned}
 \tag{21}$$

The mapping between the two descriptions is  $M \leftrightarrow Q\bar{Q}/\Lambda$ ,  $X \leftrightarrow q\bar{q}/\Lambda$ ,  $N \leftrightarrow Q\bar{q}/\Lambda$ , etc. Thus (with  $\tilde{h}$  and  $\tilde{d}$  of order 1) we expect that natural values of  $h$  and  $d$  are of order

$$h \sim d \sim \frac{\Lambda}{\Lambda_{UV}} .
 \tag{22}$$

**Estimate of scales.** We are now ready to present some typical values for the scales in this theory. To have a low-scale model of SUSY breaking we are assuming that  $\sqrt{\mathcal{F}} \sim 100$  TeV. We are also assuming that the messenger mass scale determined by  $\text{Tr} M$  is of the same order  $\sim 100$  TeV. The highest possible value for the UV scale  $\Lambda_{UV}$  can be obtained by calculating the Landau pole for the QCD coupling. This depends quite sensitively on the mass of the components of the SU(5) adjoint field  $M$ . If we assume that the leading contribution to their masses (along with the superpartners of the SM fields) are the gauge mediated contributions at 1 TeV, then we find the scale for the Landau pole to be around  $\text{few} \cdot 10^{11}$  GeV. This can be increased slightly (by a factor of at most 10) by decreasing the ratio  $\Lambda/\Lambda_{UV}$ . Thus a safe assumption would be a choice satisfying  $\Lambda_{UV} \leq 10^{11}$  GeV. Finally, we need to generate a  $\mu$ -term of order 100 GeV. The scales thus should satisfy the relations:

$$\frac{\Lambda\Lambda_{sc}^4}{\Lambda_{UV}^3} \sim \mathcal{F} \sim (100 \text{ TeV})^2, \quad \frac{\Lambda_{sc}^4}{\Lambda_{UV}^3} \sim \mu \sim 100 \text{ GeV}, \quad \Lambda_{UV} \leq 10^{11} \text{ GeV}.
 \tag{23}$$

Assuming we choose the parameters corresponding to the potential in figure 2 we can satisfy these constraints by:

$$\Lambda \sim \Lambda_{sc} \sim 10^8 \text{ GeV}, \quad \Lambda_{UV} \sim 10^{10} \text{ GeV}, \quad m' \sim 5 \cdot 10^3 \text{ GeV}, \quad \tilde{d} \sim 1.
 \tag{24}$$

In this case the bounce action interpolating between the SUSY breaking vacuum and the SUSY preserving vacuum can be estimated to be

$$S_b \sim (\Delta M)^4 / \mathcal{F}^2 \sim 5 \cdot 10^8,
 \tag{25}$$

thus the tunneling rate is suppressed by a factor of  $e^{-10^8}$ , which suggest that the metastable vacuum should have a lifetime  $\tau$  much longer than the age of the Universe ( $\sim 4 \times 10^{17} s$ ). A back-the-envelope-estimate gives

$$\tau \sim \frac{1}{100 \text{ GeV}} \frac{1s}{10^{24} \text{ GeV}^{-1}} \sqrt{\frac{2\pi}{S_b}} e^{S_b} \sim 10^{2 \cdot 10^8} s .
 \tag{26}$$

**The little hierarchy.** The minimal model described above provides a simple implementation of minimal gauge mediation with a single messenger field (up to the deviations discussed in the previous section due to the mixing of the messenger field with non-messengers). One of the main drawbacks of minimal gauge mediation is the little hierarchy problem. There are usually three separate sources for the little hierarchy problem in gauge mediated models. The first and most important is the the large mass ratio of the squark and slepton/Higgs masses dictated by the quantum numbers of the messenger fields, and is specific to gauge mediated models. The other two sources of fine tuning are generic to supersymmetric extensions of the SM. One of these is the logarithmic running of the soft mass parameters, which in this case is cut off at the messenger scale and could logarithmically enhance the soft mass parameters appearing in the Higgs potential, which thus usually requires a small stop mass to avoid fine-tuning. Finally, one also needs to make sure that the Higgs mass is above the 115 GeV LEP2 bound, which usually requires a heavy stop mass.

One of these issues is naturally resolved here, since we can take the messenger scale to be around 100 TeV, therefore the logarithmic enhancement of the soft masses is very small. The stop/slepton mass ratio can be lowered for example by changing the number of doublet messengers vs. triplet messengers [14]. Another even simpler possibility would be to change the hypercharge assignments of the messengers. However, these are usually incompatible with unification. Unification is however problematic in our model anyway due to the Landau pole. At the scale  $\Lambda_{UV}$  one would need to UV complete the model, which would likely involve taking a dual of the color gauge group, resulting in a cascading gauge theory, without a conventional perturbative unification (but rather unifying onto string theory in a warped throat [15]).

As for the Higgs bound, it strongly depends on how the  $\mu$ -term is generated. With the operator given in eq. (17) there is no additional Higgs quartic term generated, however, one can easily imagine extensions of this model where the operator leading to the  $\mu$ -term will contain NMSSM-type additional sources for a quartic term thus relaxing the fine-tuning from the Higgs mass constraint.

**Summary.** We have presented a calculable low-scale model of direct gauge mediation. The supersymmetry breaking scale can be as low as 100 TeV, and there is no hierarchy between the messenger masses and the SUSY breaking scale. The messengers play an integral part in SUSY breaking: they are composites of the dynamics that breaks supersymmetry and in the magnetic picture it is the structure of the interactions of the messengers that actually results in SUSY breaking. These interactions are such that the SUSY breaking holomorphic mass term for the messengers emerges naturally. In order to generate a real mass for the messengers one needs to shift the VEV of the meson  $M$  from the origin, which can be achieved by adding additional interaction terms in the superpotential or by including an additional U(1) gauge interaction. A  $\mu$ -term can be generated from dynamics, and some of the sources of the SUSY fine tuning problem can be eliminated. We expect that the phenomenology of the model will be quite distinctive due to the presence of additional TeV scale particles and modifications of the traditional GMSB spectrum. The main drawback

for now is the usual Landau-pole problem that is simply appearing to the the presence of the SU(5) adjoint superfield  $M$ .

One of the most interesting features of the model is a rich phenomenology at the TeV scale. Particles from the DSB sector may be directly observable at the LHC. In particular, both the scalar and fermionic components of  $M$  can have TeV scale masses. Scalars in  $M$  obtain mass due to the Coleman-Weinberg potential as well as gauge mediated contributions at two-loops. Since the Coleman-Weinberg potential is generated at one loop, it will give a leading contribution to the scalar masses as long as the  $M\phi\bar{\phi}$  coupling is large. If the coupling turns out to be small, the leading contribution to masses of scalar components in  $M$  will arise from gauge mediation. Fermionic components of  $M$  obtain mass both from GMSB loops as well as one-loop contributions arising from  $M\phi\bar{\phi}$  coupling. The latter coupling is of order one, so the new fermions will roughly have masses comparable to the gluino mass. The new scalar and fermionic particles at the TeV scale include SU(5) adjoints which transform as  $(8, 1)_0$ ,  $(1, 3)_0$ ,  $(3, 2)_{1/6}$ , and  $(\bar{3}, 2)_{-1/6}$  under  $SU(3)_c \times SU(2)_w \times U(1)$ .

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**Note added.** While this manuscript was in preparation refs. [16, 17] appeared, which also make use of metastable vacua for gauge mediation. In refs. [17] the messengers do not play an essential role in the supersymmetry breaking dynamics. The main difference from ref. [16] is that there the R-symmetry is broken via a mass term  $N\bar{N}$  (in the notation of eq. (8)) while here this is achieved by generating a VEV for  $M$ . A preliminary version of this work has been presented in [18].

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