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## An Improved Model of Direct Gauge Mediation <sup>1</sup>

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### Abstract

We present a new, improved model of gauge mediation of dynamical SUSY Breaking: the model does not have gauge messengers or  $\sim 10$  TeV scalars charged under the Standard Model (SM), thus avoiding the problem of negative (mass)<sup>2</sup> for supersymmetric SM (SSM) scalars faced by some earlier models. The gauge mediation is direct, *i.e.*, the messengers which communicate SUSY breaking to the SSM fields carry quantum numbers of the gauge group which breaks SUSY. These messenger fields couple to a modulus field. The model has a very simple particle content: the modulus and the messengers are the only chiral superfields (other than the SSM fields) in the model. The inverted hierarchy mechanism is used to generate a local SUSY breaking minimum for the modulus field in a perturbative regime thus making the model calculable.

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# 1 Introduction

There has been a substantial effort in the last few years in building models of gauge mediated dynamical SUSY breaking (for a review, see [1]). These models are predictive since they have only a few parameters and solve the supersymmetric flavor problem: the supersymmetric contributions to the flavor changing neutral currents (FCNC's) are negligible since the scalars with the same gauge quantum numbers are degenerate. Typically, these models contain vector-like fields with Standard Model (SM) gauge quantum numbers, called “messengers”, which couple to a flat direction (modulus) of the model. This modulus develops a vacuum expectation value (vev) in its scalar and  $F$  components resulting in a non-supersymmetric spectrum for the messengers. The messenger fields then communicate SUSY breaking to the supersymmetric SM (SSM) sparticles at one/two loops. Dine *et al.* first constructed models in which the messengers are *not* a part of the dynamical SUSY breaking (DSB) sector [2]. Recently, models with the messengers as an integral part of DSB sector (*i.e.* carrying the quantum numbers of the gauge group which breaks SUSY) have been built [3, 4, 5, 6, 7, 8, 9, 10]. We will call these *direct* gauge mediation (GM) models. We refer the reader to the review by Giudice and Rattazzi [1] for other models of gauge mediation.

We briefly mention the mechanisms used to stabilize the modulus and break SUSY in some of the existing models of direct GM and also the problems faced by some of these models.

In the models of references [3] and [4], a dynamical superpotential leads to a run-away behaviour along a classical flat direction. Non-renormalizable operators stabilize the potential at large expectation values. In the model of reference [3], the SUSY breaking scale is so high that the supergravity contributions to the scalar soft masses dominate over the GM contributions possibly leading to non-degenerate scalar masses and thus FCNC's. This problem was overcome in the model of reference [4]. The models of both references [3] and [4] have scalars charged under the SM originating in the

DSB sector with soft masses  $\sim 10$  TeV. These scalars drive the SSM scalar  $(\text{mass})^2$  negative through two loop Renormalization Group Evolution (RGE) [4, 11].

References [5, 6] use the “inverted hierarchy mechanism” (loop corrections) to generate a SUSY breaking minimum (which might be a local minimum) along a direction with a constant non-zero potential energy at tree level. These models have massive gauge multiplets with a non-supersymmetric spectrum. The SSM scalars get a negative contribution to their soft  $(\text{mass})^2$  by coupling at two loops to these heavy gauge multiplets (we will refer to them as “gauge” messengers) [12]. To avoid this problem, the authors of [7] used a singlet as the modulus and, to realize the inverted hierarchy, had to add another singlet and some extra matter fields.

A new mechanism of SUSY breaking, which could be used for direct gauge mediation, was discussed in [10]. In this model, different mechanisms lift the flat directions in different regions of the classical moduli space. This model requires a dynamical assumption (about non-calculable terms in the Kähler potential) to work.

In this letter, we present a new, improved model of direct gauge mediation: there are no gauge messengers or  $\sim 10$  TeV scalars charged under the SM, thus avoiding the problems of negative SSM scalar  $(\text{mass})^2$  faced by the models in [3, 4, 5, 6]. The model has a very simple structure and particle content: the only chiral superfields (in addition to the SSM fields) in the model are the modulus and the messengers, unlike the model in [7]. The model uses the inverted hierarchy mechanism to stabilize the vev of the modulus in a perturbative regime so that the model is calculable.

## 2 The Model

The gauge group and global symmetry group of the model are

$$SU(2)_1 \times SU(2)_2 \times [SU(6) \times U(1)_R], \quad (1)$$

where the group in brackets is the global symmetry group. Later, we will identify part of the  $SU(6)$  global symmetry with the SM gauge group. The particle content is

$$\begin{aligned}\Sigma &\sim (\mathbf{2}, \mathbf{2}) \times (\mathbf{1}, \mathbf{2}) \\ Q &\sim (\mathbf{2}, \mathbf{1}) \times (\mathbf{6}, \mathbf{0}) \\ \bar{Q} &\sim (\mathbf{1}, \mathbf{2}) \times (\bar{\mathbf{6}}, \mathbf{0}).\end{aligned}\tag{2}$$

The only renormalizable superpotential consistent with the gauge and global symmetries is

$$W = \lambda \Sigma Q \bar{Q},\tag{3}$$

where the gauge and global indices are appropriately summed over. The non-anomalous  $U(1)_R$  symmetry was imposed to forbid mass terms for the fields. With the superpotential in Eqn.(3), this is the only non-anomalous  $U(1)$  symmetry of the model. Consider the  $D$ -flat direction parametrized by  $\det \Sigma$ . Up to global and gauge symmetries, the  $\Sigma$  vev along this direction is

$$\langle \Sigma \rangle = \frac{1}{\sqrt{2}} \text{diag}[v, v].\tag{4}$$

This vev breaks  $SU(2)_1 \times SU(2)_2$  to the diagonal  $SU(2)_D$  at the scale  $v/\sqrt{2}$ . Three components of  $\Sigma$  are eaten by the super-Higgs mechanism. The remaining component (the superfield  $1/\sqrt{2} \text{tr} \Sigma$ ) which is the flat direction is massless and is a singlet of  $SU(2)_D$ . We denote this superfield (and the vev of its scalar component) by  $v$ . Assume  $v \gg \Lambda_1, \Lambda_2$  where the  $\Lambda_1, \Lambda_2$  are the dynamical scales of the  $SU(2)$  groups so that the gauge couplings are weak at the scale  $v$  and it suffices to use tree level matching and one loop running of gauge couplings.<sup>3</sup> We match the holomorphic gauge couplings of  $SU(2)_{1,2}$  and  $SU(2)_D$  at the scale  $v/\sqrt{2}$ :  $1/g_D^2(v/\sqrt{2}) = 1/g_1^2(v/\sqrt{2}) + 1/g_2^2(v/\sqrt{2})$  (this is the first matching condition).<sup>4</sup> The diagonal  $SU(2)_D$  has 12 fundamentals,  $Q, \bar{Q}$  and thus, at one loop, its gauge coupling does not run between

<sup>3</sup>See, however, the second footnote on page 6.

<sup>4</sup>The canonical gauge couplings should be matched at the mass of the heavy gauge boson  $\sim \sqrt{g_1^2 + g_2^2} v/\sqrt{2}$ . Using the Shifman-Vainshtein formula [13] for the relation between

the scale  $v/\sqrt{2}$  and the scale  $\lambda/\sqrt{2} v$  where all of the  $Q, \bar{Q}$  become heavy (assume, for simplicity, that  $\lambda < 1$ ). Below the scale  $\lambda v/\sqrt{2}$ , we then have a pure  $SU(2)$  gauge theory with the singlet superfield  $v$ . At the scale  $\lambda v/\sqrt{2}$  where  $Q, \bar{Q}$  are integrated out, we set the gauge coupling of the  $SU(2)_D$  theory with  $Q, \bar{Q}$  equal to the gauge coupling of the pure  $SU(2)$  gauge theory (this is the second matching condition). The two matching conditions give the holomorphic dynamical scale,  $\Lambda_L$ , of the pure  $SU(2)_D$ :

$$\left(\frac{\Lambda_L}{\lambda v/\sqrt{2}}\right)^6 = \left(\frac{\Lambda_1}{v/\sqrt{2}}\right)^2 \left(\frac{\Lambda_2}{v/\sqrt{2}}\right)^2. \quad (5)$$

This theory undergoes gaugino condensation giving the superpotential<sup>5</sup>

$$\begin{aligned} W_{eff} &= 2 \Lambda_L^3 \\ &= \sqrt{2} \lambda^3 \Lambda_1 \Lambda_2 v. \end{aligned} \quad (6)$$

The low energy theory (below  $\Lambda_L$ ) has only the field  $v$  and has  $F_v = \sqrt{2} \lambda^3 \Lambda^2$  where  $\Lambda^2 = \Lambda_1 \Lambda_2$ . At tree level, the superfield  $v$  has the canonical Kähler potential  $v^\dagger v$  in the low energy theory. Thus, the model breaks SUSY with a constant vacuum energy  $2 \lambda^6 \Lambda^4$ . The vev  $v$  is undetermined at this level. To determine  $v$ , we need to include the corrections to the Kähler potential of  $v$  due to the wavefunction renormalization  $Z$  for  $\Sigma$  [5, 6, 14]. This is the only modification to the potential since the superpotential in Eqn.(6) is exact [14].

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the canonical and holomorphic gauge couplings, we can show that the holomorphic gauge couplings have to be matched at the scale  $v/\sqrt{2}$  [14].

<sup>5</sup>A priori, there could be a superpotential term induced by the instantons in the broken  $SU(2)$  gauge group (even though the  $SU(2) \times SU(2)$  gauge group is not completely broken) [15]. However, along the flat direction  $\det \Sigma$  (for  $v \gg \Lambda_1, \Lambda_2$ ), Eqn.(6) is the only superpotential consistent with the non-anomalous  $R$ -symmetry and anomalous  $U(1)$  symmetries acting on  $Q$  and  $\bar{Q}$  and thus no other superpotential term can be generated. The instantons cannot generate this superpotential since any superpotential generated by instantons is proportional to  $\Lambda_1^{(2n)} \Lambda_2^{(2m)}$  where  $n$  and  $m$  are non-negative integers (the one loop beta-function of both  $SU(2)_1$  and  $SU(2)_2$  is 2).

The effective low energy Lagrangian is:

$$\mathcal{L} = \int d^4\theta Z(v) v^\dagger v + \int d^2\theta \sqrt{2} \lambda^3 \Lambda^2 v + \text{h.c.} \quad (7)$$

(There is no renormalization of  $Z$  below the scale  $v$  since all the fields coupling to  $\Sigma$  become heavy at  $\sim v$ .) The potential is then [5, 6, 14]

$$V(v) = \frac{2 \lambda^6 \Lambda^4}{Z(v)}. \quad (8)$$

Since  $v \gg \Lambda_1, \Lambda_2$ , we can compute  $Z$  in perturbation theory. The one loop RGE for  $Z$  is

$$\frac{dZ(t)}{dt} = \frac{2Z(t)}{16\pi^2} \left( \frac{3}{2}(g_1(t)^2 + g_2(t)^2) - 6\lambda(t)^2 \right). \quad (9)$$

where  $g_1$  and  $g_2$  are the gauge couplings of the two  $SU(2)$  gauge groups and  $t \sim \ln v$ . The potential develops a minimum along  $v$  via the “inverted hierarchy mechanism” [16] as follows. At large momentum scales, it is possible that the Yukawa contribution in the above equation dominates, since the gauge couplings are asymptotically free while the Yukawa coupling can grow with energy. This makes  $Z$  smaller as  $v$  increases. Similarly, for small values of  $v$ , the gauge contribution to Eqn.(9) dominates. Thus, for small  $v$ ,  $Z$  increases with energy. In other words, the potential decreases with energy for small  $v$  and increases with  $v$  for large  $v$ . So, we get a minimum of the potential for  $v$  such that  $\lambda \sim g$  so that  $dZ/dt$  is zero. This scale can be naturally much larger than  $\Lambda_{1,2}$  since the RG scaling is logarithmic in  $v$ . We need  $v \gg \Lambda_{1,2}$  so that we are in the perturbative regime. Upon gauging the  $SU(3) \times SU(2) \times U(1)$  subgroup of the global  $SU(6)$  symmetry, the  $Q, \bar{Q}$  act as two  $5 + \bar{5}$  messengers since they have a supersymmetric mass  $\lambda v/\sqrt{2}$  and a SUSY breaking  $(\text{mass})^2, \lambda F_v/\sqrt{2} Q\bar{Q}$  (here  $Q$  and  $\bar{Q}$  denote the scalar components). Thus the SM gauge couplings remain perturbative upto the GUT scale even for small values of  $v$ . We need  $F_v/v = \sqrt{2} \lambda^3 \Lambda^2/v \sim 10 - 100$  TeV to get SSM scalar and gaugino masses in the range 100 GeV - 1 TeV.

The scalar contained in the superfield  $v$  acquires a mass  $\sim F_v/v \times (1/16\pi^2) \sim \text{few } 100 \text{ GeV}$  (for  $\lambda, g \sim 1$ ) once the local minimum develops. This scalar is not charged under the SM. The spontaneous breaking of the  $U(1)_R$  symmetry produces a Nambu-Goldstone boson which consists mainly of the pseudoscalar in  $v$ . This ‘‘R-axion’’ can acquire a mass greater than about 100 MeV if we add a constant term to the superpotential to cancel the cosmological constant once SUSY is made local [17]. Then, the R-axion is safe from astrophysical constraints. The fermion in  $v$  is the Goldstino and is eaten by the gravitino when SUSY is made local. This model has no gauge messengers: the broken  $SU(2)$  gauge multiplet does have a non-supersymmetric spectrum, but it does not couple to the SSM scalars at one or two loops. The phenomenology of this model is similar to that of conventional gauge mediation with two families of messengers and messenger scale  $v$  [18].

There are other flat directions,  $Q^2$  and  $\bar{Q}^2$  (with the  $SU(6)$  symmetry global), which are not lifted by the superpotential of Eqn.(3). Consider the flat direction  $Q_1, Q_2 \neq 0$ . Along this direction,  $SU(2)_1$  is broken and  $\Sigma$  and two  $\bar{Q}$  s become massive. The low energy theory is  $SU(2)_2$  with four fundamentals ( $\bar{Q}$ ) which has a moduli space with a quantum modified constraint [19]. Thus, no superpotential is generated along the  $Q^2$  flat direction<sup>6</sup>. The Kähler potential is known for  $Q_1, Q_2 \gg \Lambda_1^2, \Lambda_2^2$  (it is canonical in  $Q$ ). Thus, there is a SUSY minimum along the  $Q^2$  flat direction for  $Q^2 \gg \Lambda_1^2, \Lambda_2^2$ . A similar analysis is true for the  $\bar{Q}^2$  direction.

For vev’s  $\sim O(\Lambda_1, \Lambda_2)$  along the flat directions, the  $SU(2)$ ’s are strongly coupled (at the scale of the vevs) and hence the above analysis is not valid. For example, along the flat direction  $\det \Sigma$ , for  $v$  not much larger than  $\Lambda_{1,2}$ , we can still integrate out  $Q, \bar{Q}$  but there will be higher loop and non perturbative

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<sup>6</sup>A priori, the instantons of the completely broken  $SU(2)_1$  group can generate a superpotential. Any non-perturbative superpotential has charge 2 under the non-anomalous  $U(1)_R$  symmetry. The only field with a non-zero charge under this symmetry,  $\Sigma$ , is heavy along this flat direction. Thus we expect no superpotential to be generated along this flat direction.

effects in the matching of the dynamical scales, Eqn.(5).<sup>7</sup> Also, for vev's  $\lesssim O(\Lambda_1, \Lambda_2)$ , the fields  $Q, \bar{Q}$  along the flat direction  $\det \Sigma$  and similarly  $\Sigma, \bar{Q}$  along the  $Q^2$  flat direction can not be integrated out, *i.e.*, these fields appear as additional light degrees of freedom in the superpotential. The Kähler potential along the flat directions is also not calculable for small vev's. We require a weakly coupled description for this purpose.

Each  $SU(2)$  considered separately has four flavors and thus has a dual description. In [22], dual theories to the  $SU(2) \times SU(2)$  models of this kind were constructed and it was shown that the dual theories have the same infrared physics as the original theories. We checked that in these dual theories, on adding the superpotential of Eqn.(3), the non-perturbative superpotential of Eqn.(6) is generated and that along the flat direction corresponding to  $Q^2$  no superpotential is generated (for vev's along the flat directions such that we are in the perturbative regime). However, in these dual theories, there is always one gauge group which is strongly coupled in the infra-red so that the dynamics and the Kähler potential for small vev's are still not calculable. It is possible that there is a SUSY minimum near the origin (in addition to the SUSY minimum along  $Q^2$  or  $\bar{Q}^2 \gg \Lambda_1^2, \Lambda_2^2$  with  $\Sigma = 0$ ).

Since the model is non-chiral, *i.e.*, we can add mass terms to all the fields (with the  $SU(6)$  global), we do expect a global SUSY preserving minimum. Thus, the minimum we obtained along the  $v$  direction is only a local minimum. We can estimate the tunneling rate from this "false" vacuum to the SUSY preserving minimum along the  $Q^2$  or  $\bar{Q}^2 \neq 0, \Sigma \sim 0$  flat directions. The potential energy,  $E$ , of the false vacuum is  $\sim \Lambda_1^2 \Lambda_2^2$ . The distance in field space,  $\Delta \Phi$ , from the false vacuum to the true vacuum is  $\sim v$ . Since,  $\Delta \Phi \gg E^{1/4}$ , the tunneling action from the false vacuum to the true vacuum

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<sup>7</sup>Using the techniques of [20, 21], *i.e.*, invariances under various  $U(1)$  symmetries, it is possible to show that, in this case, the matching of Eqn.(5) is exact even non-perturbatively.

can be estimated as [6, 8]:

$$S \sim 2\pi^2 \frac{(\Delta\Phi)^4}{E} \sim 2\pi^2 \frac{v^4}{\Lambda_1^2 \Lambda_2^2}. \quad (10)$$

Thus, the tunneling rate is negligibly small since  $v \gg \Lambda_1, \Lambda_2$  which was required for a perturbative calculation.

We did a numerical analysis to find out the range of possible values of  $v$ . We proceed as follows. We choose a value of  $v$  and choose  $\Lambda$  such that  $\Lambda^2/v \sim 10 - 100$  TeV to get SSM scalar and gaugino masses in the range 100 GeV - 1 TeV. We assume that the wavefunction  $Z$  is 1 at the Grand Unification (GUT) scale. Assuming  $\Lambda_1 = \Lambda_2$ , for simplicity, gives the  $SU(2)$  couplings at the GUT scale. When we gauge the SM subgroup of the global  $SU(6)$  symmetry, the Yukawa couplings for the different  $SU(6)$  components of  $Q, \bar{Q}$  are no longer the same at all energies due to RG scaling. For simplicity, we assume that the  $\lambda$ 's are all equal at the GUT scale. The value of  $v$  along with the weak scale values of the SM gauge couplings gives us the SM gauge couplings at the GUT scale. Then, with these boundary conditions (at the GUT scale), we numerically solve the RGE's for  $Z$  and  $\lambda$ 's to determine the value of  $\lambda$  at the GUT scale which gives a minimum at the chosen value of  $v$ . There is no solution for  $\lambda$  if  $v \lesssim 10^{10}$  GeV. The reason is as follows. For  $v \lesssim 10^{10}$  GeV and  $\Lambda^2/v \sim 10 - 100$  TeV, we get  $v/\Lambda_{1,2} \lesssim 10^3$  which implies that  $g_{1,2}(v)$  and hence the  $\lambda(v) \sim g_{1,2}(v)$  (required for a minimum at  $v$ ) are  $\sim 2 - 3$ . This results in the Yukawa couplings hitting their Landau poles below the GUT scale. We checked that it is possible to get a consistent minimum for values of  $v$  between  $10^{10}$  GeV and  $10^{15}$  GeV for  $O(1)$  values of  $\lambda$  and the  $SU(2)$  gauge coupling at the GUT scale. For larger values of  $v$  and hence  $F_v$ , the supergravity contribution to the scalar masses begins to dominate spoiling the degeneracy of the squarks and sleptons.

In summary, we have presented a simple model of direct gauge mediation which uses the inverted hierarchy mechanism to generate a local SUSY breaking minimum. The model does not have gauge messengers or  $\sim 10$  TeV

scalars charged under the SM and thus it avoids the problem of negative SSM scalar (mass)<sup>2</sup> faced by some of the earlier models of direct GM. This model works for messenger scales between  $10^{10}$  GeV and  $10^{15}$  GeV.

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