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FERMILAB-Pub-96-410

FERMILAB–Pub–96/410

UMD-PP-97-54

Submitted to Phys. Rev. D

Supersymmetric Models with Anomalous $U(1)$ Mediated Supersymmetry Breaking

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Abstract

We construct realistic supergravity models where supersymmetry breaking arises from the D-terms of an anomalous $U(1)$ gauge symmetry broken at the Planck scale. Effective action for these theories at sub-Planck energies (including higher dimensional terms in the superpotential) are severely restricted by the $U(1)$ symmetry and by the assumption they arise from an underlying renormalizable theory at a higher scale. Phenomenological consequences of these models are studied. It is found that they have the attractive feature that the gaugino masses, the A and B terms and the mass splittings between the like-charged squarks of the first two generations compared to their average masses can all be naturally suppressed. As a result, the electric dipole moment of the neutron as well as the flavor changing neutral current

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effects are predicted to be naturally small. These models also predict the value of the μ-term to be naturally small and have the potential to qualitatively explain the observed mass hierarchy among quarks and leptons. We then discuss examples of high scale renormalizable theories that can justify the choice of the the effective action from naturalness point of view.

I. INTRODUCTION

Supersymmetry provides ways to solve many of the puzzles of the standard model such as the stability of the weak scale under radiative corrections as well as the origin of the weak scale itself. Local supersymmetry provides a promising way to include gravity within the framework of unified theories of particle physics eventually leading the way perhaps to a theory of everything in string models. For this very good reason, supersymmetric extensions of the standard model (MSSM) have been the focus of intense theoretical activity [1] in recent years. Since experimental observations require supersymmetry to be broken, it is essential to have a knowledge of the nature and the scale of supersymmetry breaking in order to have a complete understanding of the physical implications of these theories. At the moment, we lack such an understanding and therefore it is important to explore the various ways in which supersymmetry breaking can arise and study their consequences.

There are several hints from the study of general class of MSSM which could perhaps be useful in trying to explore the nature of supersymmetry breaking. Two particular ones that rely on the supersymmetric sector of model are: (i) natural suppression of flavor changing neutral currents (FCNC) which require a high degree of degeneracy among squarks of different flavor and (ii) stringent upper limits on the electric dipole moment of the neutron (NEDM) which imply constraints on the gaugino masses as well as on the A and B terms of MSSM [2]. One could take the point of view that the above conclusions may be telling us something about the nature of supersymmetry breaking. If this is true, then it is important to isolate those SUSY breaking scenarios which realize the above properties in a simple manner and study their implications.

The standard way in which supersymmetry breaking is implemented in model building is to postulate the existence of a hidden sector where local supersymmetry is spontaneously broken and then find an appropriate way to transmit them to the visible sector. The various classes of models can be isolated depending on the way the SUSY breaking is transmitted. The two popular ones widely discussed in the literature are: (a) Polonyi type models where the SUSY breaking is transmitted via the gravitational interactions. The typical scale of SUSY breaking in such models is of order of $\sqrt{M_W M_{P\ell}} \simeq 10^{11}$ GeV and (b) the so called Gauge Mediated Susy Breaking (GMSB) type models [3], where SUSY breaking is mediated by the gauge interactions of the standard models via one and two loop radiative corrections (for instance squark squared masses are of order of $m^2_{\tilde{q}} \simeq \left(\frac{\alpha}{4\pi}\right)^2 \Lambda^2$). So the natural scale of SUSY breaking in these models is of order $10-100$ TeV raising the possibility that they are accessible to low energy tests at current and planned accelerators. The GMSB models have the extra advantage that the FCNC effects are naturally suppressed due to the fact that at the scale Λ, the squark masses are all degenerate due to the flavor blindness of the standard model gauge group. However they suffer from the so called μ problem since the gauge interactions being $U(1)_{PQ}$ symmetric do not generate a μ term. One can however add new interactions to the model to solve this problem [4].

In this paper, we discuss another class of models where the SUSY breaking is caused by the existence of an anomalous local $U(1)$ around the Planck scale, which due to the Green-Schwarz mechanism for anomaly cancellation leads [5] to a linear D-term which leads to supersymmetry breaking via the Fayet-Illiopoulos mechanism. Attempts have recently been made [6,7,8,9] to build realistic particle physics models using this new SUSY breaking mechanism. This SUSY breaking is fed down to the visible sector [7] both by the D-term as well as by the supergravity effects. It was shown in Ref. [7] that in the resulting theory, the gaugino masses are suppressed. It was also conjectured in Ref. [7] that the FCNC and CP violating effects in these models are suppressed. In Ref. [9], explicit models were constructed where both the FCNC effects as well as the electric dipole moment of the neutron were shown to naturally suppressed. It was further shown how one may have a suppressed μ term in these models using the Giudice-Masiero mechanism [10] and how one may hope to understand the fermion mass hierarchies. The primary features of these models that helped in solving the FCNC and the SUSY CP problems are that the relative squark mass difference (between the like-charged squarks of the first two generations) $\delta_q \equiv \Delta m_{\tilde{q}}^2/m_{\tilde{q}}^2$, the gaugino masses relative to the average squark masses $\delta_\lambda \equiv m_\lambda/m_{\tilde{q}}$ as well as $\mu/m_{\tilde{q}}$ and $A/m_{\tilde{q}}$ are all small, with the suppression characterized by a common parameter $\epsilon \simeq 10^{-2}$. The ϵ parameter is related to the magnitude of the $U(1)$ anomaly [11] which can be calculated in terms of the low energy fermion spectrum and is therefore not an arbitrary parameter.

It is the goal of this paper to elaborate on the various results of the Ref. [9] as well as to study the naturalness of the various higher dimensional non-renormalizable terms necessary in this model in terms of higher scale renormalizable theories. We find that indeed it is possible to generate only the desirable non-renormalizable terms in the effective low energy theory. We also study in more detail the implications of such underlying theories for fermion masses as well as FCNC and CP violating effects. The paper is organized as follows: in Sec. 2, we present the effective Lagrangian for the for the model and discuss the electroweak symmetry breaking; in Sec. 3, we discuss the electrweak symmetry breaking in the model; in Sec. 4, we show how the suppressions of the FCNC and SUSY CP effects arise; in Sec. 5, the implications of the model for the fermion masses is discussed. In Sec. 6, we discuss the naturalness of the low energy effective action in terms of an underlying renormalizable theory. In Sec. 7, we discuss the cancellation of the cosmological constant in the model and make some brief comments on the mass spectrum of the theory.

II. EFFECTIVE ACTION FOR THE MODEL

As already alluded to, the crucial feature of the model is the existence of a $U(1)$ gauge group, which is anomalous. The $U(1)$ group may be assumed to emerge from string theories. We will assume that the anomaly is cancelled by the Green-Schwarz mechanism. Since the $U(1)$ is anomalous, i.e. $\mathrm{Tr}\mathbf{Q} \neq 0$, a Fayet-Illiopoulos term which is a linear D-term is always generated as a quantum effect. We further assume that there is a pair of hidden sector fields denoted by ϕ_+ and ϕ_- which have $U(1)$ charges $+1$ and -1 respectively and that the fields of the standard model also carry $U(1)$ charges. It is the assignment of the $U(1)$ charges to quark superfields that help in the solution of the FCNC and CP problems and in qualitatively explaining the fermion mass hierarchy. We assume the following $U(1)$ charge assignment for the fields of the model (Table I):

Table I

Fields	ϕ_+	ϕ_-	H_u	H_d	Q_3, u_3^c, d_3^c	Q_i, u_i^c, d_i^c
$U(1)$-charge	+1	-1	0	+2	0	+1

Table Caption: The $U(1)$ quantum numbers of the various fields in the theory.

In the above table, $i = 1, 2$ for the first two generations. We have omitted the leptonic fields for simplicity; one could assign them same charges as the down quark sector. Note that both the superpotential W and the Kahler potential K of the model must be invariant under the anomalous $U(1)$ symmetry. Let us first discuss the superpotential W, which we write as $W = W_0 + W_1 + W_2 + W_3$, where

$$W_0 = m\phi_+\phi_-,$$

$$W_1 = h_u Q_3 H_u u_3^c,$$

$$W_2 = (h_{u,3i} Q_3 H_u u_i^c)\frac{\phi_-}{M_{P\ell}} + (h_{d,33} Q_3 H_d d_3^c)\frac{\phi_-^2}{M_{P\ell}^2} + (h_{u,ij} Q_i H_u u_j^c)\frac{\phi_-^2}{M_{P\ell}^2}$$

$$+ (h_{d,3i} Q_3 H_d d_i^c) \frac{\phi_-^3}{M_{P\ell}^3} + (h_{d,ij} Q_i H_d d_j^c) \frac{\phi_-^4}{M_{P\ell}^4},$$

$$W_3 = (W_1 + W_2) \frac{\phi_+ \phi_-}{M_{P\ell}^2} + \cdots. \tag{1}$$

In the above equation, the ellipses denote all other higher dimensional terms allowed by the gauge symmetry and make very small contributions to the effects isolated below. The parameter m is chosen to be of the order of the weak scale.

Let us now write down the Kahler potential $K(z_i, z_i^*)$ for the fields of the model generically indicated by z_i. It can be written as the sum of two terms: one that involves the bilinear terms of the form $z_i^* z_i$ and a second piece that involves mixed terms which are strongly constrained by the $U(1)$ symmetry.

$$K = K_0 + K_1,$$

$$K_0 = \sum_i |z_i|^2,$$

$$K_1 = \lambda H_u H_d \frac{\phi_- \phi_+^\dagger}{M_{P\ell}^2} + \text{h.c.} + \cdots. \tag{2}$$

In order to proceed further, we have to write down the potential of the model involving the scalar fields $\phi_\pm$, H_u^0, H_d^0 and determine the vacuum state. The part of the potential containing the ϕ_- and ϕ_+ fields reads

$$V = m^2(|\phi_+|^2 + |\phi_-|^2)$$

$$+ \frac{g^2}{2} \left(2|H_d^0|^2 + |\phi_+|^2 - |\phi_-|^2 + \xi \right)^2. \tag{3}$$

where we have ignored the terms of order $m/M_{P\ell}$ or less. Before discussing the minimization of the full potential, let us consider the part of V setting $H_u^0 = H_d^0 = 0$. It is easy to see that its minimum breaks supersymmetry as well as the anomalous $U(1)$ gauge symmetry with [7]

$$\langle \phi_- \rangle = \left(\xi - \frac{m^2}{g^2} \right)^{1/2}, \ \langle \phi_+ \rangle = 0 \tag{4}$$

$$\langle F_{\phi_+} \rangle = m \left(\xi - \frac{m^2}{g^2} \right)^{1/2} .$$

If we parameterize $\xi = \epsilon M_{P\ell}^2$, for $m \ll M_{P\ell}$, we have $\langle \phi_- \rangle \simeq \epsilon^{1/2} M_{P\ell}$ and $\langle F_{\phi_+} \rangle \simeq \epsilon^{1/2} m M_{P\ell}$. Assuming that ξ-term is induced by loop effects, one can estimate [5,7] $\xi = \frac{g^2 \mathrm{Tr}\mathbf{Q} M_{P\ell}^2}{192\pi^2}$, so that ϵ can be assumed to be of order 10^{-2}. It was pointed out in ref. [7] that the gaugino masses are generated in this model by superpotential terms of type $\lambda' W^\alpha W_\alpha \left(\frac{\phi_+ \phi_-}{M_{P\ell}^2} \right)$. As a result, one gets gaugino masses to be $m_{\lambda_g} = \lambda' \epsilon m$. If we choose $m \simeq 1$ TeV, then we need $\lambda' \sim 5$ to get the gluino mass at M_Z of 100 GeV as required by experiments.

From the K_1 term in the Kahler potential supergravity effects induce a μ-term by means of the Giudice-Masiero mechanism [10]. Indeed, K_1 induces at low energy the operator

$$\lambda \int d^4\theta H_u H_d \frac{\phi_- \phi_+^\dagger}{M_{P\ell}^2}, \tag{5}$$

giving rise to a μ-term, with $\mu = \lambda \epsilon m$. Notice that the corresponding B-term in the potential is induced at order ϵ^2, by the term $H_u H_d \phi_-^2 \phi_+ \phi_+^\dagger / M_{P\ell}^4$. There are bigger contributions to $B\mu$ from the renormalization group running of the parameters from the Planck scale down.

III. ELECTROWEAK SYMMETRY BREAKING

We integrate out the heavy field ϕ_- to obtain the effective potential of the light fields. Minimization with respect to ϕ_- gives

$$|\phi_-|^2 = \xi + |\phi_+|^2 + 2|H_d^0|^2 - \frac{m^2}{g^2}. \tag{6}$$

The effective potential of the fields (ϕ_+, H_d^0, H_u^0) is at the leading order in $m^2/M_{P\ell}^2$

$$
\begin{aligned}
V &= 2m^2|\phi_+|^2 + m_{H_u}^2|H_u^0|^2 + m_{H_d}^2|H_d^0|^2 \\
&\quad - m_3^2\left(H_u^0 H_d^0 + \text{ h.c.}\right) + D\text{−terms}, \\
m_{H_d}^2 &= |\mu|^2 + 2m^2 + m_0^2, \\
m_{H_u}^2 &= |\mu|^2 + m_0^2, \\
m_3^2 &= B\mu.
\end{aligned} \tag{7}
$$

where we have indicated by "D-terms" the usual D-terms coming from $SU(2)\otimes U(1)$ and m_i^2 denotes the supersymmetry soft-breaking terms coming from supergravity, $m_0^2 \sim \epsilon m^2$. Note that all the values in the above equation are at the Planck scale. They have to be extrapolated down to the weak scale, when we expect that $m_{H_u}^2 \simeq |\mu|^2 + m_2^2$ with $m_2^2 \leq 0$ so that H_u has a vacuum expectation value. However since $m_{H_d}^2$ is proportional to m^2 at the Planck scale, we expect it to remain sizable at the weak scale. This implies that our model will prefer a large $\tan\beta$. Also from the equation for eletroweak symmetry breaking:

$$
\frac{1}{2}M_Z^2 = \frac{m_{H_d}^2 - m_{H_u}^2 \tan^2\beta}{\tan^2\beta - 1} \tag{8}
$$

we see that for $m_{H_d} \simeq 500$ GeV, a value of $\tan\beta \simeq 10$ may be enough to get M_Z of the desired order. But for instance $\tan\beta \simeq 1$ is not at all adequate unless $m \simeq 100$ GeV, in which case we will get much too small a value for the gluino masses.

Let us now look at other parameters of the theory. It is clear from the the Eq. (1) that A_u as well as A_d suppressed by powers of ϵ (Table II):

Table II

A in units of m	$A_{u,33}$	$A_{u,3i}$	$A_{u,ij}$	$A_{d,33}$	$A_{d,3i}$	$A_{d,ij}$
ϵ suppression	ϵ	$\epsilon^{3/2}$	ϵ^2	ϵ^2	$\epsilon^{5/2}$	ϵ^3

Table caption: Degree of suppression of the various A-parameters in the theory.

Note however that these are the values at the Planck scale and they will evolve to higher values at the weak scale. It is however important to note that both the values of A and B remain of order ϵ at most since the value of B at weak scale is proportional to m_{λ_g} times the renormalization logarithm factor and similarly for A. For instance a crude estimate would lead to $B\mu \simeq \lambda'^2\epsilon^2 m^2$, which for $m \simeq 500$ GeV can be of order $(50\ GeV)^2$ or so, for $\epsilon \simeq 1/30$.

IV. FLAVOR CHANGING NEUTRAL CURRENT EFFECTS AND THE ELECTRIC DIPOLE MOMENT OF THE NEUTRON

Let us now discuss the FCNC effects in this model. To study this, we note that squark masses $m^2_{\tilde{q}}$ (both left and right handed types) receive two contributions: a universal contribution from the D-term which is of order m^2 and a non-universal contribution from the supergravity Kahler potential of order $F^2_{\phi_+}/M^2_{P\ell} \equiv \epsilon m^2$. As both these contributions are extrapolated from the Planck scale down to the weak scale the pattern of the first two generation squark masses remain practically unchanged whereas the masses of the stop receive significant contributions. It was noted in [2] that in order to satisfy the present observations of FCNC effects (such as $K^0 - \bar{K}^0$ mixing), the mixings between the $\tilde{s}$ and the $\tilde{d}$ squarks (i.e. $m^2_{\tilde{s}\tilde{d}}$) in the flavor basis or the squark mass differences between the first two generations in the mass basis must satisfy a stringent constraint. In the flavor basis, it is given by (see Dugan et al., in [2]), $\mathrm{Im}\left(\frac{m^4_{\tilde{s}\tilde{d}}}{m^4_{\tilde{q}}}\right) \leq 6 \times 10^{-8}\frac{m^2_{\tilde{q}}}{m^2_W}$. We have assumed the phases in our model to be arbitrary; therefore the most stringent constraint comes from the CP-violating part of the $K^0 - \bar{K}^0$ mass matrix. In our model, $m^2_{\tilde{s}\tilde{d}}$ arises purely from the supergravity effects are of order $\sim \epsilon m^2$ and the above FCNC constraint is satisfied if $\epsilon \simeq 10^{-2}$ or so. Thus our model confirms the conjecture of Ref. [7].

The electric dipole moment of the neutron d^e_n in supersymmetric models have been discussed in several papers [12] and it is by now well-known that the gluino intermediate states in the loop graph contributing to the d_n^e gives a contribution which is some three orders of magnitude larger than the present experimental upper limit for generic values of the parameters. The situation is different in our model since we see that a number of parameters of the model such as the gluino masses, the A and B are down by powers of ϵ. In order to see the impact of this on the NEDM, we will again consider the charge assignment for the first model where the Kahler potential induced mass splittings in the squark masses are of order ϵm^2. For the gluino contribution, we borrow from the calculation of Kizukuri and Oshimo [12], which gives:

$$\begin{aligned} d_n^e &= \frac{2e\alpha_s}{3\pi}\left(\sin\alpha_u A_u - \sin\theta_\mu \cot\beta|\mu|\right) \\ &\times \frac{m_u}{m_{\tilde{q}}^2}\frac{1}{m_{\lambda_3}} I\left(\frac{m_{\tilde{q}}^2}{m_{\lambda_3}^2}\right), \end{aligned} \tag{9}$$

where $\alpha_u = \theta_{A_u} - \theta_{\lambda_3}$ is the differerence between the phases of the A-term and the gluino mass. $m_{\tilde{q}}$ denotes the mass of the heavier of the two eigenstates. Since in this model, $m_{\lambda_3} \simeq \sqrt{\epsilon}m$ and $m_{\tilde{q}} \simeq m$, one finds that $I \simeq \epsilon$. This leads to $d_n^e \simeq \frac{2\alpha_s}{3\pi}\epsilon^{3/2}\frac{m_u}{m^2}$. Here we have used the fact that $A \sim \epsilon m$; $\mu \sim \epsilon m$. For $\epsilon \simeq 10^{-2}$, this gives an additional suppression of 10^{-3} over the prediction of generic parameter values of the MSSM (i.e. even for $m \simeq 100$ GeV, we get $d_n^e \simeq 10^{-25}$ e$\cdot$ cm). There is also a down quark contribution with a similar expression; but in this model $A_d \ll A_u$, only the second term in the above equation with $\cot\beta$ replaced by $\tan\beta$. By the same line of reasoning as above, this term also also naturally suppressed. We wish to point out that the above suppression depends on the fact that Q_1, u_1^c, d_1^c all have nonzero $U(1)$ charge. If on the other hand, d^c and u^c had zero charge, their dominant mass would come from the supergravity effect and, as a result, $m_{\tilde{d}^c}^2 \sim m_{\tilde{u}^c}^2 \simeq \epsilon m^2$. The above gluino contribution to d_n^e would then be less suppressed (by a factor $\sqrt{\epsilon}$ rather than $\epsilon^{3/2}$).

V. FERMION MASS MATRICES

Let us now discuss the pattern of fermion masses suggested by this model. First note that only the Yukawa coupling $Q_3 H_u u_3^c$ is allowed without any suppression from the ϵ factor explaining why the top quark has large mass [13,14]. On the other hand, the other Yukawa couplings are suppressed with powers of ϵ qualitatively explaining why their masses are so much smaller than the top quark mass.

From the superpotential in Eq.(1), we get the following kind of up and down quark mass matrices.

$$M_u = m_1 \begin{pmatrix} \epsilon & \epsilon & \sqrt{\epsilon} \\ \epsilon & \epsilon & \sqrt{\epsilon} \\ \sqrt{\epsilon} & \sqrt{\epsilon} & 1 \end{pmatrix} \tag{10}$$

and

$$M_d = m_2 \begin{pmatrix} \epsilon^2 & \epsilon^2 & \epsilon^{3/2} \\ \epsilon^2 & \epsilon^2 & \epsilon^{3/2} \\ \epsilon^{3/2} & \epsilon^{3/2} & \epsilon \end{pmatrix}. \tag{11}$$

where $m_{1,2}$ are mass parameters related to the $v_{u,d}$ and the Yukawa couplings. The first interesting prediction of this model is that $m_c \simeq \epsilon m_t$ and $m_b \sim \epsilon m_t$. Note that these are in qualitative agreement with observations. $m_s \sim \epsilon^2 m_t$ may also be acceptable if ϵ is not literally 10^{-2} but somewhat larger. Furthermore, if there is a horizontal symmetry between the first and the second generation, then we expect $m_u \sim m_d \sim 0$ which is also not unreasonable. We find this an encouraging aspect of the model that needs further study beyond the scope of this paper.

VI. UNDERLYING HIGH SCALE THEORY AND NATURALNESS OF THE KAHLER AND THE SUPERPOTENTIAL

In this Section we want to show that the superpotential as well as the Kahler potential chosen can indeed arise as an effective theory from an underlying renormalizable model which is valid around the Planck scale. The reason for such an exercise is the following: note that we show that in our model the μ term is suppressed naturally to the desired electroweak scale because it arises from the Kahler potential term $H_u H_d \phi_- \phi_+^\dagger$. However, if we look naively at the model, the gauge symmetries also allow a superpotential term $\int d^2\theta H_u H_d \phi_-^2 / M_{P\ell}$ which would lead to a $\mu \simeq M_{P\ell}$. This would of course be undesirable. We will show in this section there is an underlying theory where only the first term arises as an effective term at low energies and not the latter. Similarly all the higher dimensional superpotential terms that are responsible for the quark masses can also arise in this theory. This makes our choice of the Kahler as well as superpotential technically natural.

Let us assume that theory above the scale $\epsilon^{1/2} M_{P\ell}$ is characterised by the following fields in addition to the ones already given earlier: $SU(2)_L$ doublet vectorlike, colorless fields: $L, \bar{L}$ and color singlet and $SU(2)_L$ singlet fields N_i^c with $i = 1, 2, 3$ and color triplet or anti-triplet fields $D^c, \bar{D}^c, D^{c\prime}, \bar{D}^{c\prime}$. These particles are assumed to have masses $\sim M_{P\ell}$ and are expected to decouple below $M_{P\ell}$ so that at $\mu \sim \epsilon^{1/2} M_{P\ell}$, the theory will have the same structure as in sec.II. Naive decoupling arguments would seem to support this assumption. The $U(1)$ charge assignment for these fields are given in Table III.

Table III

Fields	$L, N_3^c, \bar{D}^{c\prime}$	$\bar{L}, N_1^c, D^{c\prime}$	N_2^c	$\bar{D}^c$	$\bar{D}^c$
$U(1)$-charge	+1	-1	0	+2	-2

Table caption: The $U(1)$ charge assignment of the fields of the underlying theory.

We also assume that there is a Z_2 symmetry under which the fields $L, \bar{L}, N_i^c$ $(i=1,2,3)$ are odd and the remaining fields are even. The allowed gauge and Z_2 invariant couplings involving the heavy and light fields can be written as a superpotential W_5

$$\begin{aligned} W_5 = & H_u L_1 N_1^c + M_1 L\bar{L} + M_2 N_1^c N_3^c + M_3 N_2^c N_2^c \\ & + N_1^c N_2^c \phi_+ + N_2^c N_3^c \phi_- + H_d N_1^c \bar{L} \\ & + Q_3 H_d D^c + M_D D^c \bar{D}^c + \bar{D}^c \phi_- D^{c\prime} + M_{D'} D^{c\prime} \bar{D}^{c\prime} + \bar{D}^{c\prime} \phi_- d_3^c \end{aligned} \tag{12}$$

It is the easy to see that μ, $B\mu$ and $Q_3 H_d d_3^c \phi_-^2$ terms are generated by the diagrams in Fig. 1, 2 and 3 respectively. On the other hand a term of the form $H_u H_d \phi_-^2$ is never generated in the effective low energy theory. It is possible to add to the theory extra D^c and $\bar{D}^c$ type fields with appropriate quantum numbers so that the other higher dimensional terms that lead to quark masses for lower generations can emerge.

One may have hoped that this underlying theory could be used to completely eliminate the R-parity violation from the effective low energy theory in a natural manner. It however turns out that in this particular example, it does not happen. There are however suppressions by powers of ϵ in front of the various R-parity violating couplings.

VII. COSMOLOGICAL CONSTANT AND SPARTICLE SPECTRUM

The model chosen so far has a cosmological constant of order $V_0 \sim \epsilon m^2 M_{P\ell}^2$. It is however easy to set it to zero by adding to the superpotential of the model (Eq. (1)) a constant term denoted by β^3, where β has dimension of mass. Requiring the cosmological constant to vanish implies that

$$\beta^6 = \frac{\epsilon m^2 M_{P\ell}^4 - g^{-2} m^4 M_{P\ell}^2}{3 - \epsilon + g^{-2} m^2 M_{P\ell}^{-2}} \tag{13}$$

The squark mass splittings in this case become of order $\Delta m_{\tilde{Q}}^2 \sim \frac{\epsilon}{3} m^2$. This implies that for $\epsilon \sim 1/30$, we get a suppression in the squark mass splittings of order 10^{-2}. The gravitino mass can be estimated to be $m_{3/2} \sim (\frac{\epsilon}{3})^{1/2} m$.

Let us also briefly comment on the expected masses of the superpartners in this model. The two key parameters are the superpotential mass m and the anomaly factor ϵ. We will assume $m \simeq 500 - 1000$ GeV. The value of ϵ is taken to be $\frac{1}{30}$. We then find that at $M_{P\ell}$, the gluino mass $M_{\tilde{G}} \simeq \lambda' \times (17 - 34)$ GeV. At the scale M_Z, one gets $M_{\tilde{G}} \simeq \lambda' \times (51 - 68)$ GeV. So if chose $\lambda' \geq 2$, we would be in compliance with the present experimental constraints. If on the other hand we chose $m = 100$ GeV as an extreme example, one would be driven to the light gluino scenario (which though not favored, is perhaps not excluded [15]). It thus appears that the gluino mass could be a potential embarrasment for these models if either the light gluino is definitively ruled out or one is unable to add a new source for the gluino mass to the model. At the moment we find these models to have so many attractive features that we wish to pursue them as serious candidates hoping that this issue will find a resolution. As far as the chargino masses are concerned, if the corresponding λ' is also chosen to be around two, the charginos states appear nearly degenerate since the μ-term in this model is also likely to be small. A detailed investigation of the expected sparticle spectra for plausible parameter ranges of the theory is presently under way.

Another point that distinguishes these models from the gauge mediated SUSY breaking scenarios is that we expect the squarks of the first and the second generation and the sleptons (assuming the leptons have the same charge +1 as the corresponding quarks) to have nearly the same mass. Note that in the GMSB models the sleptons are considerably lighter.

VIII. CONCLUSION

In conclusion, we have studied ways to construct interesting realistic supersymmetric models of quarks and leptons using the idea that an anomalous $U(1)$ gauge symmetry is responsible for generating supersymmetry breakdown. These models have the attractive feature that they solve several fine tuning problems of the MSSM associated with FCNC effects and electric dipole moment of the neutron. They also give desirable values for the A, the B and the μ parameters and also have the potential to qualitatively explain fermion mass hierarchies. We also show how the effective higher dimensional terms used in making the model realistic can emerge from an underlying renormalizable theory in a natural manner.

The work of R. N. M. is supported by the NSF grant no. PHY-9421385 and the work of A. R. is supported by the DOE and NASA under Grant NAG5–2788. R. N. M. would like to thank B. Brahmachari for discussions and A. R. would like to thank E. Poppitz, S. Trivedi and E. Dudas for comments and discussions.

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Figure caption

Figure 1 The one-loop diagram that leads to the effective operator $H_u H_d \phi_- \phi_+^\dagger$ that gives the μ-term at low energies.

Figure 2 The one-loop diagram that leads to the $B\mu$ term at low energies.

Figure 3 The tree diagram that gives the operator $Q H_d d^c \phi_-^2$ at low energies.