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Lawrence J. Hall and Ian Hinchliffe

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SIMPLE VIABLE MODELS OF LOW ENERGY SUPERSYMMETRY[†]

Lawrence J. Hall[*]

Lawrence Berkeley Laboratory
Department of Physics
University of California
Berkeley, California 94720

and

Ian Hinchliffe

Lawrence Berkeley Laboratory
University of California
Berkeley, California 94720

ABSTRACT

Necessary conditions are established which must be satisfied by models of low energy supersymmetry if they are to be phenomenologically acceptable at the tree level. These conditions enable one to solve the longstanding problem of giving the scalar partners of quarks and leptons large masses in such renormalizable models. A specific model is written down.

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There has been considerable renewed interest in low energy supersymmetric models, partly with a view to solving the hierarchy problem [1-10]. A number of authors have addressed the problem of constructing viable renormalizable models incorporating supersymmetry breaking at a scale of order 1 TeV. There are three types of such models: supercolor models [10] in which assumptions are made about non-perturbative effects in field theory, models in which the supersymmetry is broken softly[3, 4], and models in which the supersymmetry is broken spontaneously in perturbation theory [2, 5, 8, 9]. No model of the latter type in which the correct vacuum is uniquely determined at the tree level has been constructed. It is the purpose of this paper to construct such models.

An N = 1 supersymmetric gauge theory [11] with gauge group $G = \prod_\alpha G^\alpha$ describes the interaction of vector superfields (with scalar, fermion and vector components D^α, λ^α and v_μ^α) with chiral superfields ϕ_a (containing scalar and fermion components F_a, A_a and ψ_a). It is often convenient to use the same name for a chiral superfield and its A component. Self interactions of the chiral fields are described by a superpotential $W(\phi)$ which is at most cubic in ϕ, and does not contain hermitian conjugate fields ϕ^+. The tree level effective potential for the dynamical A components is given in terms of the auxiliary fields D^α and F_{ai} by

$$V(\phi) = \frac{1}{2}\sum_\alpha (D^\alpha)^2 + \sum_a |F_{ai}|^2$$

$$= \sum_\alpha \frac{g_\alpha^2}{2}\left(\sum_a \phi_{ai}^+ T_{ij}^{\alpha a}\phi_{aj} + \xi^\alpha\right)^2 + \sum_a \left|\frac{\partial W}{\partial \phi_{ai}}\right|^2 \qquad (1)$$

where α labels different factors of the gauge group, a the various chiral superfield representations, and i (which is always summed) the components of each representation. The gauge couplings for G^α are g^α, and $T^{\alpha a}$ are the generators of G^α in the representation of ϕ_a. The Fayet-Iliopoulos terms [12] are present only when α labels a U(1) factor. The spontaneous breaking of supersymmetry leads to a tree level relationship between scalar, fermion and vector masses [13]

$$\sum_J (-1)^{2J}(2J + 1)\,\mathrm{Tr}\, m_J^2 = \sum_\alpha \frac{g_\alpha^2 d^\alpha}{2}\,\mathrm{Tr}\, q^\alpha \qquad (2)$$

where m_J is the mass matrix for spin J fields. The sum on α now runs only over U(1) factors. The charges of the chiral multiplets are $q^\alpha = \delta_{ab}\delta_{ij} q^{\alpha a}$, and

$$d^\alpha = \left(\sum_a \langle\phi_{ai}\rangle^\dagger T_{ij}^{\alpha a}\langle\phi_{aj}\rangle + \xi^\alpha\right).$$

Only U(1) factors with $\mathrm{Tr}\, q^\alpha \neq 0$ contribute to the right hand side of Equation 2. In the standard SU(3) × SU(2) × $U(1)_y$ model, $\mathrm{Tr}\, q^y = 0$. A supersymmetric extension of this model could only be made viable at the tree level (i.e. have large masses for scalar partners of quarks and leptons, called squarks and sleptons) by having other heavy (order 100 GeV) chiral fields with fermions heavier than scalars. It is simple to see that this will not work; the trace of q^y over quarks vanishes. It can be seen from equation 1 that the sum of the squared masses of squarks will vanish as the quark mass goes to zero. Solutions to this problem of squark and slepton masses have been sought in two directions.

If the gauge group of the standard model is extended by $\tilde{U}(1)$ [14] with $\tilde{d} \neq 0$ and the trace of $\tilde{q}$ over quarks and leptons separately non-zero, chiral representations can be split in mass. This has been the basis for much attempted model building [2,5] and phenomenological

speculation. Squarks and sleptons are all given the same sign changes under $\tilde{U}(1)$ so they all receive large masses. No renormalizable realistic model of this type exists. In anomaly free models fields tend to acquire vacuum expectation values (vevs) so that $\tilde{d} = 0$ or color breaks or both [5, 8].

These failures have encouraged investigation of other mechanisms in which squark and slepton masses arise from radiative corrections once supersymmetry has been broken. This has been done in both supercolor [10] and O'Raifeartaigh [8, 9] models. Supercolor models must resort to dynamical speculation about the structure of the vacuum. In existing O'Raifeartaigh type models the tree level effective potential is not sufficient to constrain all the squark and slepton fields to have zero vevs. In the model of reference 9, while the unbroken gauge group is guaranteed to be SU(3) × $U(1)_{e.m.}$, the vevs of the sneutrinos are not fixed. Non-zero values would lead to a phenomenologically unacceptable model. In the model of Reference 8, both color and electromagnetism could be broken. Both groups show that a local minimum of the effective potential occurs when all squarks and sleptons have zero vevs; in this vacuum two or three loop radiative corrections to masses make these models viable. The tree level vacuum degeneracy necessitates the calculation of the effective potential beyond leading order to determine whether these local minima are indeed the true vacuum.† We will avoid such difficulties at the expense of enlarging the gauge group, considering models of the former type which have sufficient squark and slepton masses at the tree level. We will establish necessary conditions which the superpotential must satisfy in such models, and illustrate one way of fulfilling these requirements in a specific model.

† We understand that the authors of Reference [8] are undertaking such a calculation.

One constraint comes from the requirement of anomaly freedom. The general conditions for anomaly freedom are rather cumbersome, but they can be usefully reorganised into mass relations for scalar fields. For any particular supermultiplet ϕ_x transforming as an irreducible represention R of G, it follows from Equation 1 that the trace of the mass matrix for ϕ_x from D^2 is given by

$$4 \; \mathrm{Tr} \; \left(\frac{\partial^2 V}{\partial\phi^{\dagger}_{xi}\partial\phi_{xj}}\right) = 4\sum_\alpha g_\alpha^2 d^\alpha \mathrm{Tr} \; q^{\alpha x}_{ij} \tag{3}$$

where the equality holds only if $<\phi_x> = 0$. Since α now runs only over Abelian factors $q^{\alpha x}_{ij} = \delta_{ij} \; q^\alpha(R)$. The contribution from D^2 to the mean square mass of representation R is $m_D^2(R) = 4\sum_\alpha g_\alpha^2 d^\alpha q^\alpha(R)$. If the model is to be free of $SU(3)^2 \times U(1)$ anomalies, then

$$\sum_{ab} \mathrm{Tr}(T^{3a}T^{3b}q^\alpha) = \sum_R C(R)q^\alpha(R) = 0 \tag{4}$$

where T^{3a} are the color generators in representation ϕ_a, and the irreducible representations R have color Casimirs $C(R)$ and charges $q^\alpha(R)$. Equation 4 can be converted into a mass relation by multiplying by $g_\alpha^2 d^\alpha$ and summing over α

$$\sum_R C(R)m_D^2(R) = 0 \; . \tag{5}$$

Different anomaly constraints lead to different mass formulae, but for our purposes the relation among colored representations is the most powerful.

The chiral representations ϕ_a may be divided into two classes, those in which the fermions are light ϕ_ℓ (for example less than 2 GeV), and the rest ϕ_h. The ϕ_ℓ will contain light quark and lepton superfields which can appear in $W(\phi)$ only in cubic interactions. The contributions from W to light squark and slepton masses will be comparable to the fermion masses, and if we neglect such small masses

$m^2(R_\ell) = m_D^2(R_\ell)$. Equation 5 now reads

$$\sum_{R_\ell} C(R_\ell)m^2(R_\ell) + \sum_{R_h} C(R_h)m_D^2(R_h) = 0. \tag{6}$$

In a realistic model $m^2(R_\ell)$ must be large and positive. The mass relation then implies that there must exist at least one colored field (K) which has negative contribution to its mass squared from D^2: $m_D^2(K) < 0$. Any realistic model must therefore have a superpotential which gives K a large positive mass squared preventing the spontaneous breakdown of color: $m^2(K) = m_D^2(K) + m_W^2(K) > 0$. The utility of such mass relations is that they sum up the effects of all U(1) factors, and show that a negative contribution to $m_D^2(K)$ from one U(1) cannot be cancelled by a positive contribution from another. From considerations of these mass relations it seems that little will be gained by having additional U(1) factors beyond $\tilde{U}(1)$.

Giving K a large positive $m_W^2(K)$ is not trivial. If K is in a real representation of color a term K^2 in W implies $m_D^2(K) = 0$ which contradicts the definition of K. If K is charged under some of the U(1)s, a term $K\bar{K}$ (where $\bar{K}$ is an independent chiral superfield with conjugate charges and color) in W can certainly give $m_W^2(K) > 0$. As $C(K)m_D^2(K) + C(\bar{K})m_D^2(\bar{K}) = 0$, this pair of fields does not provide a cancellation of the positive squark contribution to Equation 6. The only possibility is for K to appear in a cubic coupling in W in which two of the three fields are colored. The color singlet field (J) must acquire a vev thus giving $m_W^2(K) > 0$. We will now show that in order to guarantee this, it is necessary to have additional terms in W involving J. Assume J only couples to U(1) factors and

$$W = \lambda K^2 J (\mathrm{or} \; \lambda K\bar{K}J) + f(\mathrm{anything \; but} \; J).$$

Consider the equation for a stationary point in V

$$\frac{\partial V}{\partial J^\dagger} = 4\langle J\rangle \sum_\alpha g_\alpha^2 q_\alpha(J) d^\alpha = 0$$

where we have inserted vevs and required that $\langle K\rangle = 0$. Since $m_D^2(K) < 0$ (or $m_D^2(K) + m_D^2(\bar K) < 0$), $\sum_\alpha g_\alpha^2 q_\alpha(J) d^\alpha > 0$ so that a W of this form gives the unsatisfactory result $\langle J\rangle = 0$. If J and K transform non-trivially under some other broken non-Abelian factor a similar argument leads to the same conclusion. Therefore additional terms must be present in W which force $\langle J\rangle \neq 0$, ie $W = \lambda K^2 J$(or $\lambda K\bar K J$) + (Driving terms for J). All previous models failed because they did not satisfy this criterion. A superpotential of the above form necessarily has a mass scale, so that models with trilinear superpotentials [5] either have light squarks and sleptons at tree level or broken color.

The simplest superpotential of this type is

$$W_1 = \lambda K K' J + f X (J J_1 - \mu^2)$$

where J_1 has conjugate quantum numbers to J, and X is neutral. One can easily see that this form is unlikely to be successful since only the product JJ_1 is fixed. As $m_D^2(J_1) < 0, J_1$ can get a much larger vev than J such that $\xi^\alpha + q^\alpha(\langle J\rangle^+\langle J\rangle - \langle J_1\rangle^+\langle J_1\rangle) < 0$ for all α. Squark and slepton fields are then driven to get vevs such that d^α tries to relax to zero. This difficulty can be avoided by adding mass terms for J and J_1

$$W_2 = \lambda K K' J + f X (J J_1 - \mu^2) + M J_1 J_2 + \bar M J J_3 .$$

We now have an O'Raifeartaigh model [15] so that supersymmetry will be broken in this sector. For $f^2\mu^2 > M\bar M$, and in the limit of vanishing gauge coupling constants the minimum of this potential is

$$\langle J\rangle^2 = \frac{M}{\bar M}\left(\mu^2 - \frac{M\bar M}{f^2}\right), \quad \langle J_1\rangle^2 = \frac{\bar M}{M}\left(\mu^2 - \frac{M\bar M}{f^2}\right), \quad \langle K\rangle = 0, \langle X\rangle \text{ undetermined},$$

$$\langle J_2\rangle = -\frac{f}{M}\langle J\rangle\langle X\rangle \text{ and } \langle J_3\rangle = -\frac{f}{M}\langle J_1\rangle\langle X\rangle .$$

The indeterminacy will be removed by D^2 terms. X, J_2, J_3 acquire infinitely large vevs such that $J_2^2 - J_3^2$ is finite; $\langle J\rangle$ and $\langle J_1\rangle$ are changed infinitesimally. The d^α can then relax to zero. This difficulty can be avoided by removing J_3 from the theory. With $g_\alpha = 0$, $\langle J\rangle$ would be infinite but the D^2 terms prevent this.

Before writing down a realistic model the problem of weak interaction breakdown must be solved while maintaining $d^\alpha \neq 0$. The Higgs superfields (H) which couple to quarks and leptons must be constrained to have non-zero vevs. Since squarks and sleptons have positive mass squared from D^2 these Higgs fields generally have $m_D^2(H) < 0$, and easily get vevs which allow d^α to relax toward zero. This can be prevented by using a Fayet-Iliopoulos [12] type model. For example, a U(1) gauge theory with $W = mH_+H_-$ leads to a potential

$$V = m^2(H_+^\dagger H_+ + H_-^\dagger H_-) + \frac{g^2}{2}(H_+^\dagger H_+ - H_-^\dagger H_- + \xi)^2$$

which for $\xi g^2 > m^2$ has $\langle H_+\rangle = 0$ and $\langle H_-\rangle = \xi - \frac{m^2}{g^2}$ so that $d = \frac{m^2}{g^2} > 0$. The field H_- will play the role of the Higgs, while H_+ is an additional field. In the remainder of this paper an explicit realistic model is written down which incorporates the general features discussed above.

The gauge group is chosen to be SU(3) × SU(2) × U(1) × $\tilde{U}(1)$, and the anomaly free representations of the chiral fields are shown in Table 1. The superpotential is

$$\begin{aligned} W = {} & \lambda_K K^2 J + \lambda_T T^2 J + \lambda_S \bar S S J + f X (J J_1 - \mu^2) \\ & + M J_1 J_2 + m J J_1 \\ & + m_H \bar H H + m_H' \bar H' H' + g_E L \bar E H + g_D Q \bar D H + g_U Q \bar U H' . \end{aligned} \tag{7}$$

K, T, S and $\bar S$ are all prevented from getting vevs by the same

O'Raifeartaigh sector. Three generations of quarks and leptons are understood. In the following we will set $m_H = m'_H$ for convenience. This is not essential. The physics of the model is not changed provided m'_H/m_H is order one. The potential is:

$$V = \left|\frac{\partial W}{\partial \phi_a}\right|^2 + \frac{1}{2}(D^2)^2 + \frac{1}{2}(D^3)^2$$
$$+ \frac{g_y^2}{2}\Big(\frac{1}{6}Q^+Q - \frac{2}{3}\bar{U}^+\bar{U} + \frac{1}{3}\bar{D}^+\bar{D} - \frac{1}{2}L^+L^- + \bar{E}^+\bar{E} + \frac{\sqrt{5}}{2}S^+S - \frac{\sqrt{5}}{2}\bar{S}^+\bar{S}$$
$$- \frac{1}{2}H^+H^- + \frac{1}{2}\bar{H}^+\bar{H} + \frac{1}{2}H'^+H' - \frac{1}{2}\bar{H}'^+\bar{H}'\Big)^2$$
$$+ \frac{\tilde{g}^2}{2}\Big(Q^+Q + \bar{U}^+\bar{U} + \bar{D}^+\bar{D} + L^+L + \bar{E}^+\bar{E} - 2K^+K - 2T^+T - 2S^+S - 2\bar{S}^+\bar{S}$$
$$+ 4J^+J - 4J_1^+J_1 + 4J_2^+J_2 + (11)^{1/3}Y^+Y - 2H^+H + 2\bar{H}^+\bar{H} - 2H'^+H' + 2\bar{H}'^+\bar{H}' + \tilde{\xi}\Big)^2$$

ξ^Y has been set to zero. This is weakly natural because renormalization (infinite) of ξ^Y vanishes when the trace of q^Y over each set of degenerate scalars vanishes [16]. For a suitable range of parameters the global minimum is at

$$\langle X\rangle = -m/f$$
$$\langle J\rangle^2 = \frac{m_1}{m_2}\left(\mu^2 - \frac{m_1 m_2}{f^2}\right) \quad \langle J_1\rangle^2 = \frac{m_2}{m_1}\left(\mu^2 - \frac{m_1 m_2}{f^2}\right)$$
$$\langle H\rangle = \binom{h}{0}, \ \langle H'\rangle = \binom{0}{h} \quad \text{with } m_1^2 = M^2 - 2m_H^2, \ m_2^2 = 2m_H^2 \ \text{and}$$
$$h^2 = \left(\frac{\tilde{\xi}}{4} - \frac{m_H^2}{8\tilde{g}^2}\right) + \left(\mu^2 - \frac{m_1 m_2}{f^2}\right)\left(\frac{m_1^2 - m_2^2}{m_1 m_2}\right).$$

All other vevs vanish. In this vacuum $\tilde{d} = \frac{m_H^2}{2\tilde{g}^2}$ so that the squark and slepton masses are $\frac{m_H}{\sqrt{2}}$. In order to obtain the correct value for M_W $4h^2 = (240\ \text{GeV})^2$. To ensure that $\langle K\rangle, \langle T\rangle, \langle S\rangle$ and $\langle\bar{S}\rangle$ vanish, $\lambda_{K,T,S}$ should not be too small. A suitable set of parameters is

$$f = \lambda_K = \lambda_T = \lambda_S = 1/2, \ \tilde{g} = 1/14$$

$$\mu = 1, \ M = \frac{1}{2}, \ m_H = \frac{1}{5}, \ \tilde{\xi} = 2.34, \ m = \frac{1}{3}$$

with TeV the unit of mass. The Yukawa couplings (g_E, g_D, g_U) have the same values as in the standard model. For these particular values the heaviest particles are K and T with mass 860 GeV, while S and $\bar{S}$ weigh 400 GeV and squarks and sleptons weigh 140 GeV.

A model is natural in the strong sense if W contains all terms consistent with the symmetries of the model. The only symmetry under which X tranforms and which does not eliminate $\mu^2 X$ is an R symmetry. Therefore, any strongly natural O'Raifeartaigh model must have an R symmetry. As vector superfields are real they necessarily have zero R character, which means that the fermionic component λ^α transforms non-trivially. For unbroken gauge factors λ^α can not get a Dirac mass, and a Majorana mass term of the type $\lambda\lambda$ is forbidden by R unless R is spontaneously broken. Therefore, in a strongly natural model either the gluinos are massless or there is a massless Goldstone boson (or axion if R has a color anomaly). To avoid there difficulties we will not use R symmetries. It might appear that the term mJJ_1 ensures that W has no R invariance. This is incorrect; a non-linear R symmetry exists in W. When X is displaced by its vev and the component Lagrangian inspected a linear R symmetry is found to exist. We must modify W to remove this problem. This is easily but inelegantly done by adding a term

$$\lambda_Z Z(\bar{N}N + \bar{P}P - \mu'^2) + m_3(\bar{N}N' + \bar{N}'N)$$
$$+ m_4(\bar{P}P + \bar{P}'P) + m_5\bar{N}N + m_6\bar{P}P$$

where Z is a singlet field, N, P, P', N' are colored but singlets under other gauge factors, and $\bar{N}$, $\bar{P}$, $\bar{P}'$, $\bar{N}'$ are in conjugate representations to N, P, P', N'. The parameters μ, λ, m_i are chosen such that vevs vanish, and $m_5 \neq m_6$ so that the R symmetry does not reappear. Our model will be natural in the weak sense that terms not present in V at the tree level will be induced with small finite order g^2 coefficients.

The large ratio of squark to K masses gives rise to the possibility that the vacuum may be unstable at one loop. For example, a negative contribution to the squared squark mass of order $\frac{g^2}{4\pi} M_K^2$ could be disastrous. Fortunately the only such contribution is for $g = \tilde{g}$ and these terms are much smaller than the tree level squared squark mass. Other one loop corrections to squark and slepton masses are harmless.

The model has three neutral gauge bosons. The photon and its couplings have been constructed to be identical to the standard model. Four fermion coupling between quarks and leptons mediated by the $\tilde{Z}$ boson (Zum) have strength $\sim G_F/1000$. $M_W = M_Z \cos\theta(1+\rho)$ and $\rho < 0.01$ if $m_H - m_H' < 100$ GeV.

Apart from quarks and leptons the only light fermions at tree level are the photino, gluinos, ψ_Z, ψ_Y and the Goldstino which are all massless. The winos, zino and Zumino have Dirac masses of order 100 GeV. The photino and gluino acquire Majorana masses from radiative corrections.

The model has an axion. Independent $\tilde{U}(1)$ rotations can be performed on the set of fields J, J_1, J_2, K, T, S and $\bar{S}$, and on the remainder. Only one of these currents is gauged, the other has an anomaly giving mass to the corresponding Goldstone boson. This, and other aspects of this model are under investigation. [17]

We have demonstrated that it is straightforward to construct models of low energy supersymmetry with a well defined vacuum. Criteria have been established which any such model must satisfy. An explicit example has been given. There are several problems to be overcome before such models can be perturbatively grand unified. For example Tr $\tilde{q} \neq 0$, and also the SU(3) and SU(2) coupling constants are not asymptotically free above 1 TeV. This latter problem is probably more serious, and afflicts

all existing supersymmetric models other than those of Ref.[3,4]. Work in this direction is in progress.

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TABLE I: Left Handed Chiral Superfields

	SU(3)	SU(2)	$U(1)_y$	U(1)
Q_x	3	2	1/6	1
$\bar{U}_x$	$\bar{3}$	1	-2/3	1
$\bar{D}_x$	$\bar{3}$	1	1/3	1
L_x	1	2	-1/2	1
E_x	1	1	1	1
K	8	1	0	-2
T	1	3	0	-2
S	1	2	$\sqrt{5}/2$	-2
$\bar{S}$	1	2	$-\sqrt{5}/2$	-2
J	1	1	0	4
J_1	1	1	0	-4
J_2	1	1	0	4
X	1	1	0	0
Y	1	1	0	$(11)^{1/3}$
H	1	2	-1/2	-2
$\bar{H}$	1	2	+1/2	+2
H'	1	2	+1/2	-2
$\bar{H}'$	1	2	-1/2	+2

Quantum numbers of fields appearing in the superpotential Eq. 7. x = 1, 2, 3 label generations.

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