Ref.TH.3253-CERN

LOW ENERGY SOFTLY BROKEN SUPERSYMMETRY

A.B. Lahanas and T.R. Taylor^{*)}

CERN -- Geneva

ABSTRACT

We consider supersymmetric low energy Lagrangians based on the gauge group $SU(3) \times SU(2) \times U(1) \times \tilde{U}(1)$, where $\tilde{U}(1)$ is an additional axial group à la Fayet. Models with explicit soft breaking of supersymmetry are considered as an alternative to models with axial anomalies. We argue that soft breaking can be very efficient and illustrate this with a simple example in which the introduction of a negative mass squared term for only one scalar field leads to a phenomenologically interesting model.

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^{*)} On leave of absence from the Institute of Theoretical Physics, Warsaw University, Warsaw, Poland.

Supersymmetry might solve the naturalness problem¹⁾ in theories with elementary Higgs scalars. The absence of quadratic divergences in supersymmetric theories prevents Higgs boson masses from getting huge radiative corrections. Many attempts have been made to build a realistic supersymmetric model of low energy effective interactions, with the supersymmetry broken at energies accessible to accelerators presently under construction. The most remarkable class of models, developed by Fayet, is based on the gauge group $SU(3) \times SU(2) \times U(1) \times \tilde{U}(1)$, where $\tilde{U}(1)$ is an additional local axial symmetry²⁾⁻⁴⁾. Supersymmetry can then be spontaneously broken by the Fayet-Iliopoulos D term⁵⁾. The extra $\tilde{U}(1)$ factor may generate large masses of squarks and sleptons - the spin zero superpartners of quarks and leptons - provided that one assigns a positive $\tilde{U}(1)$ charge $\tilde{Y}$ to the three families of left-chiral supermultiplets involving all observed fermions. Moreover, this type of breaking is favoured over other ones by the natural suppression of flavour changing neutral electroweak interactions⁶⁾. Unfortunately, the aforementioned $\tilde{Y}$ assignment leads to anomalies in triangle diagrams with $\tilde{U}(1)$ axial vector current insertions. In order to cancel axial anomalies one is forced to introduce, in particular, new strongly interacting supermultiplets with a negative $\tilde{Y}$. The scalar components of such supermultiplets generally acquire vacuum expectation values (v.e.v.s), breaking the $SU(3)$ colour group^{4),7)}. This is the reason why, as far as we know, nobody has succeeded in constructing an anomaly-free model with a realistic pattern of symmetry breaking in tree approximation. In such a situation we prefer to sacrifice exact global supersymmetry than to accept an axial anomaly. By insisting on the absence of quadratic divergences we are led to softly broken supersymmetry⁸⁾. The purpose of this paper is to demonstrate that one can use the explicit breaking of supersymmetry very efficiently. Our proposal is to introduce negative mass squared terms for some spin zero particles in order to trigger the desired gauge symmetry breaking. We will illustrate this attitude with a simple example of a supersymmetric model in which the introduction of an explicit breaking controlled by only one extra parameter leads to a phenomenologically interesting model.

Let us consider one family of leptons and quarks which form fermionic components of left-chiral scalar superfields with the following $SU(3) \times SU(2) \times U(1) \times \tilde{U}(1)$ quantum numbers:

$$\begin{aligned}
 L(1, 2, -\frac{1}{2}, 1) & \quad ; \quad \bar{E}(1, 1, 1, 1) \\
 Q(3, 2, \frac{1}{6}, 1) & \quad ; \quad \bar{D}(\bar{3}, 1, \frac{1}{3}, 1) \quad ; \quad \bar{U}(\bar{3}, 1, -\frac{2}{3}, 1)
 \end{aligned}
 \tag{1}$$

In addition, we need two Higgs superfields

$$H(1, 2, -\frac{1}{2}, -2) \quad ; \quad \bar{H}(1, 2, \frac{1}{2}, -2) \quad (2)$$

In order to cancel anomalies we choose the following set of additional superfields

$$\begin{aligned} \Phi(3, 1, \frac{1}{3}, -2) \quad ; \quad \bar{\Phi}(\bar{3}, 1, -\frac{1}{3}, -2) \\ S(1, 1, 0, 4) \quad ; \quad \bar{N}(1, 1, 0, 1) \end{aligned} \quad (3)$$

The most general renormalizable interaction among the chiral scalar superfields has the form of the superpotential

$$\begin{aligned} W = a(L\epsilon H)\bar{E} + b(Q\epsilon H)\bar{D} + c(Q\epsilon\bar{H})\bar{U} \\ + d(H\epsilon\bar{H})S + e(\Phi\bar{\Phi})S + f(L\epsilon\bar{H})\bar{N} \end{aligned} \quad (4)$$

where the totally antisymmetry tensor ϵ contracts the $SU(2)$ indices, and the $SU(3)$ indices are contracted in a self-evident way^{*)}.

The superfield content of our model has been motivated by the $SO(10) \times \tilde{U}(1)$ decomposition of the exceptional group $E_6^{(9,3,4)}$. However, our case differs from the E_6 one by a reversed $U(1)$ hypercharge assignment for Φ and $\bar{\Phi}$, i.e., $1/3$ and $-1/3$ instead of $-1/3$ and $1/3$, respectively. In this way, while preserving the anomaly cancellation and tracelessness of $U(1)$ and $\tilde{U}(1)$, we avoid baryon and lepton number violation in the interactions $\bar{U}\Phi\bar{E}$, $(Q\epsilon L)\bar{\Phi}$ and $\epsilon_{ijk}\bar{U}^i\bar{D}^j\bar{\Phi}^k$. The tracelessness of $U(1)$ and $\tilde{U}(1)$ protects D_Y and $\tilde{D}_Y$ from quadratic divergences¹⁰⁾ and leaves open the possibility of embedding our model in a grand unified one. The superpotential W automatically conserves baryon and lepton number. Moreover, violation of these quantum numbers does not appear in non-renormalizable interactions of dimension five; however, it can occur in higher dimensions as a possible manifestation of grand unification. Therefore the proton decay rate is suppressed by at least two inverse powers of the unification scale as in conventional grand unified schemes.

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^{*)} By an appropriate redefinition of superfields we make all couplings real and positive.

The superfield $\bar{N}$ represents a right-handed neutrino and its spin zero superpartner. A vacuum expectation value for the Higgs scalar would generate a Dirac mass for the neutrino proportional to the coupling constant f . A very small value of f is natural in the sense of 't Hooft¹¹⁾, since putting $f = 0$ increases the symmetry of the Lagrangian by an additional R symmetry under which only the superfield $\bar{N}$ transforms non-trivially, i.e., by a non-vanishing phase. This symmetry will remain unbroken. Therefore, if the neutrino is massless in the tree approximation, it will remain so to all orders of perturbation theory. Otherwise its mass would have to be renormalized by counterterms which are at most logarithmically divergent. Hereafter we put $f = 0$, i.e. we choose the case of a strictly massless neutrino. Nevertheless, our analysis obviously applies to the massive neutrino case ($f \neq 0$) as well.

The potential of the scalar field is given by

$$\begin{aligned}
 V = & a^2 |H\bar{E}|^2 + a^2 |L\epsilon H|^2 + |aL\bar{E} + bQ\bar{D} + d\bar{H}S|^2 \\
 & + b^2 |Q\epsilon H|^2 + c^2 |Q\epsilon \bar{H}|^2 + |bH\bar{D} + c\bar{H}\bar{U}|^2 \\
 & + |cQ\bar{U} + dHS|^2 + |dH\epsilon\bar{H} + e\Phi\bar{\Phi}|^2 + e^2 |\Phi S|^2 + e^2 |\bar{\Phi}S|^2 \\
 & + \frac{1}{2} D_Y^2 + \frac{1}{2} D_{\bar{Y}}^2 + \frac{1}{2} \vec{D}_W^2 + \frac{1}{2} \vec{D}_S^2
 \end{aligned} \tag{5}$$

where

$$\begin{aligned}
 D_Y = & g_Y \left(-\frac{1}{2} |L|^2 + |\bar{E}|^2 + \frac{1}{6} |Q|^2 + \frac{1}{3} |\bar{D}|^2 - \frac{2}{3} |\bar{U}|^2 \right. \\
 & \left. - \frac{1}{2} |H|^2 + \frac{1}{2} |\bar{H}|^2 + \frac{1}{3} |\Phi|^2 - \frac{1}{3} |\bar{\Phi}|^2 \right) + \xi_Y
 \end{aligned} \tag{6}$$

$$\begin{aligned}
 D_{\bar{Y}} = & g_{\bar{Y}} \left(|L|^2 + |\bar{E}|^2 + |Q|^2 + |\bar{D}|^2 + |\bar{U}|^2 - 2|H|^2 - 2|\bar{H}|^2 \right. \\
 & \left. - 2|\Phi|^2 - 2|\bar{\Phi}|^2 + 4|S|^2 + |\bar{N}|^2 \right) + \xi_{\bar{Y}}
 \end{aligned} \tag{7}$$

$$\vec{D}_W = g_W (L^\dagger \vec{t} L + Q^\dagger \vec{t} Q + H^\dagger \vec{t} H + \bar{H}^\dagger \vec{t} \bar{H}) \tag{8}$$

$$\vec{D}_s = g_s (\bar{Q} \vec{T} Q - \bar{D} \vec{T} D - \bar{U} \vec{T} U + \Phi \vec{T} \Phi - \bar{\Phi} \vec{T} \bar{\Phi}) \quad (9)$$

The capital letters now denote the scalar components of the corresponding supermultiplets. The parameters ξ_Y and $\xi_{\tilde{Y}}$ are the coefficients in the Lagrangian of the D components of the U(1) and $\tilde{U}(1)$ gauge vector supermultiplets, respectively. The SU(3), SU(2), U(1) and $\tilde{U}(1)$ coupling constants are denoted by g_S, g_W, g_Y and $g_{\tilde{Y}}$, whereas $\vec{T}$ and $\vec{\tilde{T}}$ are the SU(2) and SU(3) generators in their fundamental representations, respectively.

Direct inspection shows that the potential (5) has no minimum with the desired properties. The easiest way to see this is to consider the special case of $\xi_Y = 0$ and $\xi_{\tilde{Y}} > 0$. Then the potential has an absolute minimum equal to zero at the following v.e.v.s

$$\langle H^\alpha \rangle = \langle \bar{H}_\beta^* \rangle \varepsilon^{\beta\alpha} = h^\alpha \quad (10)$$

$$-\langle \Phi^i \rangle = \langle \bar{\Phi}^{*i} \rangle = f^i \quad (11)$$

with the constraints

$$|h|^2 + |f|^2 = \frac{\xi_{\tilde{Y}}}{4g_{\tilde{Y}}} ; \quad d |h|^2 = e |f|^2 \quad (12)$$

and zero v.e.v.s of all other fields. The vacuum is supersymmetric and breaks colour. This is an example of the aforementioned general difficulty encountered when trying to construct a realistic model.

We remedy this problem by modifying the potential to

$$V' = V - \mu^2 |S|^2 \quad (13)$$

expecting that a negative mass squared for the scalar field S might trigger its vacuum expectation value at the minimum and generate a positive mass squared for the Φ and $\bar{\Phi}$ scalars. This in turn may save the colour group from spontaneous breaking. The potential V' explicitly breaks supersymmetry at the scale μ , but only by one operator of dimension two^{*}). Though the general analysis of such a potential as a function of many complex variables is rather complicated, we are able to prove the following statements:

*) There are some indications that this type of explicit breaking may be suggested by arguments involving gravity; see J. Wess, talk given at the Erice Workshop (1981).

i) For a suitable choice of parameters in the Lagrangian the point

$$\langle H^a \rangle = \langle H^1 \rangle \delta_1^a ; \quad \langle \bar{H}^a \rangle = \langle \bar{H}^2 \rangle \delta_2^a ; \quad \langle S \rangle \quad (14)$$

with

$$\left. \begin{aligned} |\langle H^1 \rangle|^2 \\ |\langle \bar{H}^2 \rangle|^2 \end{aligned} \right\} = \frac{2}{d^2(24g_Y^2 - d^2)} [d^2 g_Y^2 \xi_Y - (d^2 - 8g_Y^2) \frac{\mu^2}{4}] \mp \frac{g_Y \xi_Y}{g_W^2 + g_Y^2 - 2d^2} \quad (15)$$

$$|\langle S \rangle|^2 = 2 \frac{d^2 + 8g_Y^2}{d^2(24g_Y^2 - d^2)} \left[\frac{\mu^2}{4} - \frac{2d^2}{d^2 + 8g_Y^2} g_Y \xi_Y \right] \quad (16)$$

and zero v.e.v.s of the other fields is a local minimum of the potential V' (13). To be explicit we quote an appropriate range of parameters which yields this minimum, and for which the minimum possesses the additional property

$$\left| \frac{g_Y \xi_Y}{g_W^2 + g_Y^2 - 2d^2} \right| \ll \frac{1}{2} (|\langle H^1 \rangle|^2 + |\langle \bar{H}^2 \rangle|^2) < |\langle S \rangle|^2 \quad (17)$$

This range is determined by the following self-consistent set of inequalities:

$$\begin{aligned} \xi_Y &> 0 \\ 8g_Y^2 &< d^2 < \min \left\{ \frac{1}{2} g_W^2, 24 g_Y^2 \right\} \\ 6g_Y^2 \xi_Y &< \mu^2 < \frac{4d^2}{(d^2 - 8g_Y^2)} g_Y \xi_Y \\ e^2 &> 2(24g_Y^2 - d^2) \frac{g_Y \xi_Y}{\mu^2 - 6g_Y^2 \xi_Y} + 8g_Y^2 \end{aligned} \quad (18)$$

which ensures positivity of the mass matrices for all physical scalar fields.

ii) Within the same range of parameters and for vanishing Yukawa coupling constants a , b and c , the minimum (14) is a global one. In fact, the realistic Yukawa coupling constants are much smaller than all the other couplings involved in the Lagrangian (we assume g_Y to be of order g_Y) and we therefore expect our minimum to correspond to the true vacuum.

The v.e.v.s (14) spontaneously break the gauge group $SU(3) \times SU(2) \times U(1) \times \tilde{U}(1)$ down to $SU(3) \times U(1)_{EM}$. As in the standard Weinberg-Salam theory, the photon turns out to be a combination of the neutral $SU(2)$ gauge boson W_3 and the $U(1)$ hypercharge boson B :

$$A = B \cos \theta_w + W_3 \sin \theta_w \quad (19)$$

where

$$\cos \theta_w = \frac{g_w}{(g_w^2 + g_Y^2)^{1/2}}$$

The charged $SU(2)$ vector bosons get a mass

$$m_w^2 = \frac{1}{2} g_w^2 (|\langle H^1 \rangle|^2 + |\langle \bar{H}^2 \rangle|^2) \quad (20)$$

The $\tilde{U}(1)$ boson $\tilde{Z}$ mixes with the state Z orthogonal to the photon through the following mass matrix elements:

$$m_Z^2 = \frac{m_w^2}{\cos^2 \theta_w} \quad (21)$$

$$m_Z^2 = 16 m_w^2 \frac{g_Y^2}{g_w^2} + 32 g_Y^2 |\langle S \rangle|^2 = 16 g_Y^2 \left(\frac{\mu^2}{d^2} - 3 \frac{m_w^2}{g_w^2} \right) \quad (22)$$

$$m_{Z\tilde{Z}}^2 = \frac{2 g_w g_Y}{\cos \theta_w} \frac{g_Y \tilde{\xi}_Y}{(g_w^2 + g_Y^2 - 2d^2)} \quad (23)$$

From the physical point of view, the most attractive situation is described by Eq. (17). By an appropriate choice of parameters (in particular by increasing μ sufficiently high above m_w) we can achieve the $\tilde{U}(1)$ breaking scale $\langle S \rangle$ large enough to ensure that the $\tilde{U}(1)$ current does not substantially modify the low-energy neutral current phenomenology. The $\tilde{Z}$ boson will be much heavier than all the others ($m_{\tilde{Z}} \gtrsim 1$ TeV) and will not mix with the electroweak Z boson.

A very important feature of our model is that it predicts the following relation between the mass of a fermion and its spin zero superpartner (if it corresponds to a scalar field which does not acquire a v.e.v.):

$$m_{sf}^2 - m_f^2 = g_w I_f \langle D_3 \rangle + g_Y Y_f \langle D_Y \rangle + g_{\tilde{Y}} \tilde{Y}_f \langle D_{\tilde{Y}} \rangle \quad (24)$$

where I_f , Y_f and $\tilde{Y}_f$ denote the third component of weak isospin, hypercharge and the $\tilde{U}(1)$ charge of a chiral multiplet, respectively. The relation (24) is characteristic of theories with the supersymmetry spontaneously broken by a D term¹²⁾ and it explains why the superpartners of light fermions are not observed at present energies. The chiral multiplets for all of the light fermions carry positive $\tilde{Y}_f$ charge, so that in the large μ limit we obtain for the relation (24) [using Eqs. (7), (14)-(18) and (24)]:

$$m_{sf}^2 (\tilde{Y}_f > 0) \simeq m_f^2 + \frac{1}{4} \tilde{Y}_f \left(\mu^2 - \frac{2d^2}{g_w^2} m_w^2 \right) > m_f^2 + \frac{1}{4} \tilde{Y}_f (\mu^2 - m_w^2) \quad (25)$$

This means that squarks and sleptons will be much heavier than quarks and leptons provided that μ is large enough, which is consistent with previous considerations. A new feature when compared with other models is the appearance of the mixing between superpartners of left-handed and right-handed fermions, as a consequence of the following term in the Lagrangian:

$$ad \langle S \rangle \langle \bar{H}^2 \rangle (j_{eL})(j_{eR})^* + cd \langle S \rangle \langle H^1 \rangle (j_{uL})(j_{uR})^* \quad (26)$$

$$+ bd \langle S \rangle \langle \bar{H}^2 \rangle (j_{dL})(j_{dR})^* + \dots + c.c.$$

Again, in the large μ limit, these couplings are suppressed by a factor $O(m_f/\mu)$ relative to the diagonal mass terms, but can still play a significant role in the radiative generation of the gluino and photino masses. We will come to this point later.

The masses of all particles can now be easily calculated. Before describing the spin 1/2 and 0 particle content in the tree approximation let us define the physical neutral scalar field

$$H^{1'} = e^{-iA_g \langle H^1 \rangle} H^1 - |\langle H^1 \rangle| \quad (27)$$

and similarly $\bar{H}^{2'}$ and S' . We restrict ourselves to the general characteristics of the spectrum:

i) spin 1/2:

- $H, \bar{H}, S, \Phi, \bar{\Phi}$ with masses of order $m_{\tilde{Z}} \gtrsim 1 \text{ TeV}$ and the $\tilde{U}(1)$ gauge fermion with mass $m_{\tilde{Z}}$;
- fermionic superpartners of the W_+, W_- and Z electroweak bosons with masses m_W and m_Z , respectively;
- ordinary quarks and leptons (including the right-handed neutrino);
- massless gluinos and photino.

ii) spin 0:

- squarks, sleptons, $\Phi, \bar{\Phi}$ and $\text{Re}S'$ with masses of order $m_{\tilde{Z}}^*$;
- $\text{Re}H^{1'}$, $\text{Re}\bar{H}^{2'}$, one charged Higgs boson and its antiparticle, with masses of order m_W ;
- massless axion:

$$\left[1 + \left|\frac{\langle H^1 \rangle}{\langle \bar{H}^2 \rangle}\right|^2 + \left|\frac{\langle H^1 \rangle}{\langle S \rangle}\right|^2\right]^{-1/2} \text{Im} H^{1'} + \left[1 + \left|\frac{\langle \bar{H}^2 \rangle}{\langle H^1 \rangle}\right|^2 + \left|\frac{\langle \bar{H}^2 \rangle}{\langle S \rangle}\right|^2\right]^{-1/2} \text{Im} \bar{H}^{2'} + O\left(\frac{m_W}{m_{\tilde{Z}}}\right) \text{Im} S' \quad (28)$$

orthogonal to the plane spanned by the Goldstone bosons absorbed by the Z and $\tilde{Z}$ gauge bosons through the Higgs mechanism.

The presence of an axion is not a surprise. Our Lagrangian possesses a Peccei-Quinn $U(1)_{PQ}$ global symmetry¹³⁾, corresponding to an R symmetry under which all superfields transform by a common phase. The spontaneous breakdown of $U(1)_{PQ}$ gives rise to the Weinberg-Wilczek axion¹⁴⁾. From instanton effects the axion may get a mass of order 100 keV, probably inconsistent with present experimental limits. It is worth mentioning that a modified version of our model, with some extra superfields including a total gauge singlet, could break $U(1)_{PQ}$ explicitly and solve the problem of the axion by giving it a large mass in a similar way to how it was done in Ref. 4).

The same R symmetry, if unbroken, would protect the masslessness of the photino and gluinos to all orders of perturbation theory. Since this R is spontaneously broken, the photino and gluinos will get masses from radiative

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*) We neglect here mixings between $\text{Re}S'$ and $\text{Re}H^{1'}$, $\text{Re}\bar{H}^{2'}$, which are of order $O(m_W/m_{\tilde{Z}})$.

corrections^{15),16)}. The dominant contribution to the gluino mass comes from the diagram shown in Fig. 1¹⁵⁾, which involves the off-diagonal mass term of the scalar components of Φ and $\bar{\Phi}$, analogous to (26). It gives a mass correction

$$m_g \simeq \alpha_s \frac{m_w^2}{m_{\tilde{Z}}^2} O(1) \quad (29)$$

where the factor $O(1)$ depends on the parameters of the theory. If we exclude an unnatural adjustment of parameters, Eq. (29) gives us an upper bound on the gluino mass of order 10 GeV. For the photino, the α_s coupling constant should be replaced by α_{EM} , hence its mass will be of order 500 MeV.

The generalization of our model to more than one family of quarks and leptons is straightforward. Equation (24) implies that the mass squared differences between squarks (sleptons) belonging to different families are naturally of the same order as those between their corresponding quarks (leptons), which according to Refs. 6) seems to be a necessary condition for the suppression of neutral flavour-changing electroweak interactions.

ACKNOWLEDGMENTS

We wish to thank H.P. Nilles and K. Tamvakis for many useful discussions and important comments. We have also benefited from discussions with J.-M. Gerard and D.V. Nanopoulos. We bow down to the lotus feet of J.S. Hagelin for kindly correcting the manuscript.

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FIGURE CAPTION

The dominant contribution to the radiatively induced gluino mass. f_Φ, f_Φ^- denote fermions and s_Φ, s_Φ^- their scalar superpartners.

Fig. 1