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FROM HIGH ENERGY SUPERSYMMETRY BREAKING TO LOW ENERGY PHYSICS
THROUGH DECOUPLING

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A B S T R A C T

We reconsider a realistic model of electroweak and strong interactions with calculable mass spectrum at the tree level in which supersymmetry and an extra gauge group factor $\tilde{U}(1)$ beyond $SU(3) \times SU(2) \times U(1)$ are both broken at very high energies: $M_{\text{SUSY}} \approx \sqrt{M_W M}$, $M_{\tilde{U}(1)} \approx M$ with $M \gg M_W$. In spite of these high-energy scales, especially the large scale of supersymmetry breaking, the low energy spectrum - including the relevant Higgs boson - is decoupled from the heavy degrees of freedom. Due to the "non-renormalization" theorems this decoupling persists to all orders in perturbation theory.

Models of electroweak and strong interactions based on spontaneously broken global supersymmetry are faced with a general difficulty which was first pointed out by Fayet ¹⁾. For a gauge group $G = SU(3)_C \times SU(2) \times U(1) \times \tilde{G}$ from the general form of the mass matrices ^{1), 2)} in the quark and scalar quark sector one readily finds that the contribution to the scalar mass matrix

$$\sum_{\alpha} g_{\alpha} T_{ab}^{\alpha} \langle D^{\alpha} \rangle \quad (1)$$

where D^{α} are auxiliary field components of the vector multiplets and T_{ab}^{α} are the matrix generators, needs to be positive definite over all quarks of given charge and helicity. Otherwise, one scalar partner of the quarks will be degenerate or lighter than the down quark itself, at least at the level of tree masses. One way out is to use the so-called O'Raifeartaigh breaking ³⁾ and to have the mass splitting, as well as, in many cases, the breaking of $SU(2) \times U(1)$, arising mainly as a quantum effect through radiative corrections ⁴⁾. The only solution to the problem of having a realistic mass spectrum at the tree level is to enlarge the gauge group, as proposed by Fayet ¹⁾, to include at least an extra generator $\tilde{Y}$ with charges of definite signs over all matter multiplets and a non-vanishing $\langle D_{\tilde{Y}} \rangle$ term contributing to (1) and actually dominating over all other possible contributions from T_{3L} and Y . The minimal gauge group is $SU(3) \times SU(2) \times U(1) \times \tilde{U}(1)$ with the ξ term ⁵⁾ associated with the new factor $\tilde{U}(1)$ used to trigger supersymmetry breaking.

The main point of this article is to show that in order to get a realistic low energy mass spectrum for quarks, leptons, Higgses and their superpartners, the $\tilde{U}(1)$ need not be a low energy phenomenon, but rather it can be broken at super-large energies M relative to the breaking scale M_W of the electroweak group. In spite of the fact that the members of the $\tilde{Y}$ vector multiplet become superheavy ($\sim M$), they decouple from the low energy physics, controlled by the much smaller scale M_W . The vacuum expectation value of $\langle D_{\tilde{Y}} \rangle \approx M_W^2$ can contribute with a positive sign to Eq. (1), which is what is needed to get the right mass splitting between quarks, leptons and their scalar superpartners.

An important feature of models in which this situation occurs is that supersymmetry is broken at the intermediate energy scale $M_{SUSY} = \sqrt{M_W M}$. This in turn induces splittings among bosons and fermion masses of order M_{SUSY} in the superheavy sector of the model, but not in the light sector because of the decoupling of the Goldstino from the light sector itself. In fact, the typical mass splitting in the light sector is, to all orders in perturbation theory,

$$\sim (M_{SUSY}^2/M) = M_W.$$

This can be summarized as follows

$$(m_B^2 - m_f^2)_{\text{light}} = F a_{\text{light}} \quad (2)$$

and $F \sim M_W M$ and $a_{\text{light}} = M_W/M$. F is the goldstino decay constant and a is the coupling of the goldstino to a supermultiplet. This is the mechanism which keeps the appropriate $[SU(2) \text{ doublets}]$ Higgs fields light ($\sim M_W$). The decoupling mechanism alluded to above occurs because of the non-renormalization theorems⁶⁾ of global supersymmetry. An intriguing point is that the breaking of the $\tilde{U}(1)$ group at a superlarge scale, rather than at low energy as originally proposed by Fayet¹⁾, seems to be unavoidable in order to get a consistent model. The first reason for this has to do with the fact that it appears difficult to avoid supersymmetric solutions in models with D supersymmetry breaking⁵⁾ and no Yukawa couplings. In order to prevent these solutions from having lower vacuum energy than the desired minimum, the Yukawa couplings g_Y present in the superpotential must satisfy a lower bound $g_Y \gtrsim M_W/M$ where M is the $\tilde{U}(1)$ symmetry breaking scale. A second theoretical reason could be related to the impossibility of having a model with supersymmetry breaking, the absence of A-B-J triangular anomalies introduced by the $\tilde{U}(1)$ group and at the same time cancellation of quadratic divergences. Finally, a further point is related to R invariance as a continuous global symmetry. The flat space Lagrangians with a D supersymmetry breaking need such a symmetry to allow the inclusion of gravity⁷⁾. On the other hand, the unwanted effects of this symmetry may indeed be avoided by the couplings of supergravity, but only if supersymmetry itself is broken at sufficiently high energy. Phenomenological problems which are naturally avoided with a $\tilde{U}(1)$ group broken at M , have to do with neutral current phenomenology. In particular, a $\tilde{U}(1)$ group naturally suppresses any flavour changing neutral current⁸⁾ and strong CP violation⁹⁾.

In order to illustrate a model which exhibits the feature discussed above, we reconsider the multiplet structure of Ref. 10).

Beyond the 15 chiral multiplets which describe a normal family of quarks and leptons, at least four extra fields, R, R^c, S and P , only transformation under $\tilde{U}(1)$ must be introduced¹⁰⁾.

The most general superpotential compatible with gauge invariance is

$$f = \alpha Q u^c H^c + \beta Q d^c H + \gamma L e^c H + \delta L P H^c + \lambda H H^c S + \mu R R^c + m S R^c \quad (3)$$

with obvious contractions of group indices. The model therefore has, as accidental global symmetries, baryon and lepton number as well as an R type symmetry ¹⁰⁾. As in Ref. 10), we drop for a moment the term mSR^C controlling another $U(1)$ symmetry called X ($R-R^C$ conservation). The only dimensionful parameter in f is μ . The potential coming from the D^α terms of the 13 group generators is fixed by the charge assignments given in the Table. We introduce a Fayet-Iliopoulos term ⁴⁾ only for the $\tilde{U}(1)$ factor. The related D term is

$$D = g \left(-2 |H|^2 - 2 |H^C|^2 - 4 |R^C|^2 + 4 |R|^2 + 4 |S|^2 + |P|^2 + |Q|^2 + |u^C|^2 + |d^C|^2 + |L|^2 + |e^C|^2 + \xi \right) \quad (4)$$

ξ , which is taken to be positive, receives a quadratically divergent renormalization ¹¹⁾

$$\xi \rightarrow \xi + g \frac{\Lambda^2}{(4\pi)^2} T_2 \tilde{Y} = \xi + 12 g \frac{\Lambda^2}{(4\pi)^2} \quad (5)$$

Its natural value is therefore of order Λ^2 itself, where Λ is the energy cut-off of the model ($\Lambda = M_{Pl}$?). In order to understand the features of the potential of our model, it is instructive to consider first the case with unbroken supersymmetry.

It can be shown that for $\mu = 0$ supersymmetry is unbroken and all supersymmetric minima are given by non-vanishing expectation values of R, R^C, S, P and satisfying Eq. (4)

$$-4 |R^C|^2 + 4 |R|^2 + 4 |S|^2 + |P|^2 + \xi = 0 \quad (6)$$

Any of these minima gives rise to the breaking of $\tilde{U}(1)$ only, occurring at a mass $M = g\sqrt{\xi}$ with the formation of a massive vector multiplet. It is important to notice at this point that by letting some of the Yukawa couplings $\alpha, \beta, \gamma, \delta$ be zero, new supersymmetric minima other than the ones given by Eq. (6) would be introduced. This would be a disaster because they would not be removed by the parameter μ introduced in f as given in Eq. (3). It is therefore necessary to have $\alpha, \beta, \gamma, \delta$ all different from zero, and actually satisfying a lower bound given below.

Let us now discuss the full potential where the term $\mu R R^C$ is present. Unbroken supersymmetry now requires $\langle R \rangle, \langle R^C \rangle$ to be zero, which is incompatible with Eq. (6). Supersymmetry is broken and the minimization of the full potential forces the fields with a negative $\tilde{U}(1)$ charge R^C, H, H^C to acquire a non-vanishing vacuum expectation value, whereas the positively charged fields, at least in a certain range of parameters, all remain at zero vacuum expectation value.

If we consider the potential $V(H, H^C, R^C)$ for the range of values of parameters

$$\begin{aligned} g_2^2 &> 2\lambda^2 \\ 4g^2\xi &> \mu^2 \left(1 + \frac{8g^2}{\lambda^2}\right) \end{aligned} \quad (7)$$

it reaches the minimum at

$$H = \begin{pmatrix} \frac{\mu}{\sqrt{2}\lambda} \\ 0 \end{pmatrix}, \quad H^C = \begin{pmatrix} 0 \\ \frac{\mu}{\sqrt{2}\lambda} \end{pmatrix}, \quad R^{C2} = \frac{\xi}{4} - \mu^2 \left(\frac{8g^2 + \lambda^2}{16g^2\lambda^2} \right) \quad (8)$$

At this point

$$D_{SU(2)} = D_Y = 0, \quad D = \frac{\mu^2}{4g}, \quad \frac{\partial F}{\partial R} = f_R = \mu \left(\frac{\xi}{4} - \mu^2 \left(\frac{8g^2 + \lambda^2}{16g^2\lambda^2} \right) \right)^{1/2} \quad (9)$$

and f_R is the only non-vanishing f auxiliary field. By working out the quadratic part of the over-all V , one sees that all other scalar fields are stabilized at zero. The positive D term as given by Eq. (9) gives the correct mass splitting between the quarks, leptons and their scalar superpartners. On the other hand, the supersymmetry breaking, although also controlled by μ , takes place at a mass scale of the order of the geometrical mean of the two scales $\mu, \sqrt{\xi}$. In fact, at the minimum

$$V = F^2 = \frac{1}{2} D^2 + |f_R|^2 \simeq |f_R|^2 \simeq \frac{\xi \mu^2}{4} = g \xi D \quad (10)$$

where F is the goldstino decay constant. As already pointed out in Ref. 10), this is the first example of high energy supersymmetry breaking which coexists with a low energy scale for the $SU(2) \times U(1)$ breaking. Remarkably, a very similar

phenomenon occurs in the class of models based on O'Raifeartaigh breaking which have been recently discussed ¹²⁾.

Several comments are now in order. The first has to do with the would-be supersymmetric minima if we make one of $\alpha, \beta, \gamma, \delta$ vanish in Eq. (3). Around these points, to first order in these couplings, generally denoted by α , the potential takes the value

$$V \approx \alpha^2 \bar{f}^2 \quad (11)$$

to be compared with Eq. (10).

One ends up with a lower bound of the type

$$\alpha \geq \mathcal{O}\left(\frac{\mu}{\sqrt{3}}\right) = \mathcal{O}\left(\frac{\lambda g}{g_2} \cdot \frac{M_W}{M_{\tilde{\gamma}}}\right) \quad (12)$$

In turn this gives a lower bound on all fermion masses

$$m_f \geq \mathcal{O}\left(\frac{\lambda g}{g_2} \cdot \frac{M_W}{M_{\tilde{\gamma}}}\right) \quad (13)$$

including the neutrino. Specifically the neutrino receives a mass through the coupling in Eq. (3) given by δL_{PH}^C , of the Dirac type, with no breaking of lepton number. The fermion component of the P multiplet can be identified with the right-handed neutrino. Another comment concerns the decoupling of the fields with a tree-level mass of order μ from the superheavy sector of the theory. This is particularly relevant in view of the fact that supersymmetry is broken at a mass scale $M_{SUSY} \approx \sqrt{M_W M}$.

One may be tempted to argue that all the splitting of the theory, including particles of mass of order μ , will be of order M_{SUSY} , which would be a disaster for the low energy theory since $M_{SUSY} \approx \sqrt{M_W M}$. However, this does not happen because the auxiliary field f_R , whose vacuum expectation value is responsible for the large supersymmetry breaking, is only coupled to superheavy particles. Radiative corrections will induce couplings of f_R to light multiplets but all at most of order $\mu/\sqrt{\xi}$. In fact, having integrated the superheavy degrees of freedom, one can write an effective Lagrangian for the light fields which is supersymmetric up to mass terms of order μ , and interaction weighted by at least one power of μ^2/ξ .

This decoupling prevents the Higgs doublets from receiving a mass of order M_{SUSY} . Splitting of order $\sqrt{M_W M}$ only occurs among the members of super-heavy multiplets to which the goldstino ⁹⁾

$$\psi_g \sim R_{\text{fermion}} \quad (14)$$

is mainly coupled.

Another comment concerns the problem of A-B-J anomalies. The group $\tilde{U}(1)$ introduces new anomalies into the model. If one introduces new fields in order to remove them, for example a triplet and an antitriplet T, T^c , supersymmetry is unbroken ¹⁰⁾. We do not know of any example without A-B-J anomalies, traceless preserving $SU(3)_C \times U(1)_{\text{em}}$, and employing the Fayet-Iliopoulos mechanism for the breaking of supersymmetry. In the absence of such a model one could dispose of the $\tilde{U}(1)$ anomalies by letting $\xi \rightarrow \infty$.

In this way one would have a supersymmetric Lagrangian for the light fields with an explicit soft supersymmetry breaking ¹³⁾ coming from mass terms. More reasonably if we fix ξ to be the cut-off scale (M_{Pl} ?), then the anomaly problem is not a severe one for the low energy spectrum. In any case, gravitational effects, for example, will play a rôle ⁷⁾.

As a final point, we would like to emphasize that the Lagrangian described in this paper predicts a detailed spectrum for all the particle states we have introduced to implement supersymmetry. The weak vector bosons Z^0, W form almost massive supersymmetric vector multiplets. They are degenerate in mass with their Dirac fermionic partners, of charge 0 and ± 1 , respectively, and split from their scalar partners (a neutral and a charged scalar from the H, H_C multiplets) by an amount $\Delta m_{\text{gauge}}^2 = \mu^2$. The matter chiral multiplets have scalar partners of quarks and leptons with universal mass splitting given by $\Delta m_{\text{matter}}^2 = \frac{1}{4}\mu^2$. As pointed out in Ref. 1), these two mass formulae give an upper bound to the mass of the scalar partners of quarks and leptons of about 40 GeV. Other than the particles discussed so far, there are the massless gluinos (because of R symmetry) and an almost massless Peccei-Quinn axion (because of the breaking at M_W of the additional global symmetry X) which are probably in conflict with experiments.

Both of these problems could be solved by introducing ¹⁰⁾ an extra field N , without gauge interactions, with a new superpotential term

$$f_N = \eta N R^c S + f(N) \quad (15)$$

A suitably chosen polynomial $f(N)$ gives rise to an explicit breaking of both R and X symmetries.

A perhaps more interesting mechanism for breaking the X symmetry alluded to above is to include in the Lagrangian the only other allowed term

$$f_m = m S R^c \quad (16)$$

This term breaks the X symmetry but not R invariance. R invariance can be broken by gravitational interactions which can give a sizeable effect if $\xi \approx O(M_{Pl}^2)^{7)}$. The superpotential (3) with the term (16) included fulfils the decoupling theorem only if

$$m \leq O\left(\frac{\mu^2}{\sqrt{\xi}}\right) \quad (17)$$

while the bound given by Eq. (7) remains almost unaffected. For instance, charge conservation now requires

$$g_2^2 \left(\frac{1}{2\lambda^2} + \frac{m\sqrt{\xi}}{2\mu^2} + \frac{m^2\xi}{2\mu^4} \right) > 1$$

instead of $g_2^2/2\lambda^2 > 1$ as given by Eq. (7).

With the term f_m given by (16) added to the previous Lagrangian, the vacuum expectation value of the Higgs doublets becomes

$$|H|^2 = |H^c|^2 = \frac{\mu^2}{2\lambda^2} + \frac{m\sqrt{\xi}}{2\lambda} \quad (18)$$

while the vacuum expectation value of the D auxiliary field is still given by Eq. (9).

The would-be Goldstone boson associated with X invariance becomes massive with mass proportional to $m\sqrt{\xi}$. The goldstino particle is now given by the following combination of fermions

$$\frac{\mu\sqrt{\xi}}{2} \left(R_f + \frac{m}{\mu} S_f \right) + \frac{\mu^2\lambda}{8} \quad (19)$$

where λ is the gauge fermion of the $\tilde{U}(1)$ group. From (18) and (19) it is evident that the inequality (17) must be satisfied ^{*)}. A more interesting possibility is that these phenomenological problems so far discussed could be solved by a suitable embedding of the present model in a GUT and/or by the inclusion of local supersymmetry ⁷⁾. The simplest GUT which comes to mind is an $SU(5) \times \tilde{U}(1)$ gauge theory with a $\tilde{U}(1)$ group broken at superlarge energies. Whether such a model can be successfully constructed is an open question.

Finally, we would like to mention that the coupling to supergravity of models such as the one discussed in the present paper, could give a substantial modification of the mass spectrum and provide new interactions ⁷⁾. This is because local supersymmetry requires sizeable effects, which are no longer negligible when the supersymmetry breaking occurs at a high energy scale

$$M_{\text{susy}} \sim \sqrt{M_W M_{\text{Pl}}} \simeq (10^{10} - 10^{11}) \text{ GeV}$$

The effect of supergravity corrections on the global supersymmetric potential will be discussed elsewhere ¹⁴⁾.

*) The smallness of the parameter m , as well as that of the neutrino Yukawa coupling δ , is "natural" in the technical sense that both of them control symmetries and are only renormalized by logarithms.

G	$=$	$SU(3)$	$\otimes$	$SU(2)$	$\otimes$	$U(1)$	$\otimes$	$\tilde{U}(1)$
Q_i^α		3		2		$\frac{1}{6}$		1
$u_i^{c\alpha}$		$\bar{3}$		1		$-\frac{2}{3}$		1
$d_i^{c\alpha}$		$\bar{3}$		1		$\frac{1}{3}$		1
L^α		1		2		$-\frac{1}{2}$		1
e^c		1		1		1		1
H^α		1		2		$-\frac{1}{2}$		-2
$H^{c\alpha}$		1		2		$\frac{1}{2}$		-2
S		1		1		0		4
R		1		1		0		4
R^c		1		1		0		-4
P		1		1		0		1

TABLE - The chiral supermultiplet content of the model, with the relative transformation properties under G .

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