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LOW-ENERGY PHENOMENOLOGY OF BROKEN SUPERSYMMETRY

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ABSTRACT

In a class of theories with spontaneously broken supersymmetry the lightest supersymmetric particles are expected to be neutral fermions with masses $< O(m_W)$ which are mainly supersymmetric partners of Higgses, while gauginos, squarks and sleptons are expected to be somewhat heavier. We discuss the low-energy phenomenology of such theories, as contrasted to models with light gauginos. LEP and other e^+e^- colliding ring machines offer good prospects for testing such theories.

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Nowadays there is considerable interest in the applications of supersymmetric field theories, particularly in the prospects they offer¹⁾ for alleviating the hierarchy problem in grand unified theories (GUTs). For this to be possible, the supersymmetric partners of known particles cannot be very heavy²⁾, even if the primordial scale of supersymmetry breaking is as large as the Planck mass. Since these unseen supersymmetric particles may be accessible to forthcoming experiments, the time is ripe to investigate their phenomenology more completely than has been done up to now. There have already been extensive studies of the low energy phenomenology of broken supersymmetry^{3),4)}, but these have mostly focused on a different class of models from those proposed^{5)-7),2)} recently. It is therefore important to explore the phenomenological contrasts between the two classes of models. In the original models³⁾ of supersymmetry broken by D terms in the potential⁸⁾, the lightest supersymmetric particles were expected to be the photino and goldstino (gravitino), called collectively *nuinos* and having masses $\ll 1$ MeV, followed by the *gluinos* with masses of a few GeV. The decay chain of every supersymmetric particle would therefore terminate in a *nuino*.

In the more recently developed class^{5)-7),2)} of models where *susy* breaking mass splittings are generated by radiative corrections, possibly triggered by F terms in the potential, the lightest supersymmetric particles are expected to be neutral fermionic partners of Higgs particles, called *shiggses*, which have slight admixtures of *gauginos*. For cosmological reasons, they must have masses $\leq 0(100)$ eV or least a few GeV⁹⁾, and their charged partners $\tilde{H}^\pm$ must be heavier¹⁰⁾, weighing at least 15 GeV to be compatible with limits from PETRA. The neutral weak interactions and hence the forward-backward asymmetries¹¹⁾ for $e^+e^- \rightarrow \tilde{H}^+\tilde{H}^-$ will be very different from those of a conventional sequential heavy lepton (see Table 1). The next lightest *susy* particle is expected to be the $U(1)$ *gaugino* $\tilde{B}$, followed by the *sleptons* $\tilde{l}$, the $SU(2)$ *gauginos* $\tilde{W}$, *squarks* $\tilde{q}$ and finally *gluinos* $\tilde{g}$. The decay chain of every *susy* particle would terminate in a stable neutral *shiggs*. In addition to the *susy* particles, there should be several spin-zero neutral Higgs particles, the lightest of which is likely to be coupled more strongly to charge $-1/3$ quarks than to charge $+2/3$ quarks (see Table 2).

These general results emerge naturally in a definite model²⁾ which can also be used to make more specific predictions. These include a) a lower bound on the top quark mass of about 25 GeV, with calculable corrections related to the masses of other observable particles, b) upper bounds on the *gaugino* masses, e.g., $m_{\tilde{W}} \leq 0(3)m_W$, c) a possible upper bound on the *slepton* masses $m_{\tilde{l}} \leq 0(1)m_W$ and d) the existence of two light neutral *shiggses*, one stable and the other perhaps slightly heavier and with a lifetime which could well be $> 0(10^{-13})$ sec.

We now proceed to a characterization of the class^{2),6),7)} of susy theories that we consider. Because of the structure of chiral superfield couplings, there must be two SU(2) Higgs superfields whose neutral scalar components H_1^0 and H_2^0 develop vacuum expectation values in order to give masses to the quarks of charges $+2/3$ and $-1/3$ respectively. Together with the SU(3) $\times$ SU(2) $\times$ U(1) gauge supermultiplets and the quark and lepton supermultiplets, these may comprise essentially all the particles with masses $\leq 0(10)$ TeV. We imagine that these superfields are embedded in a susy GUT such as SU(5), though the details of any such model are unimportant here, except for the possible inclusion⁶⁾ of a light gauge singlet chiral superfield Y which is expected to couple only to the Higgs supermultiplets.

We suppose that radiative corrections play the dominant role in generating the splittings between low-energy particles and their supersymmetric partners, along the lines discussed in Refs. 2),5)-7) and illustrated in Fig. 1. Low-energy susy breaking starts with masses for the SU(3) $\times$ SU(2) $\times$ U(1) gauginos of uncertain magnitude, but with the natural ratios:

$$m_{\tilde{g}} = \alpha_3 M_0, m_{\tilde{W}} = \alpha_2 M_0, m_{\tilde{B}} = \frac{5}{3} \alpha_1 M_0 \quad (1)$$

Note that in this class of model one does not expect the photino and zino to be the mass eigenstates. The gauginos feed positive $(\text{mass})^2$ to the squarks and sleptons, which in turn feed negative $(\text{mass})^2$ to the light Higgses in such a way that

$$\langle 0 | H_1^0 | 0 \rangle = v_1, \langle 0 | H_2^0 | 0 \rangle = v_2, v_1^2 + v_2^2 = v^2 : m_{W^\pm}^2 = \frac{g_2^2 v^2}{2}$$

Note also that in this class of models one expects the gauginos $\tilde{W}$ to be somewhat heavier than the $W^\pm$, and comparable with the squarks and sleptons. In such a theory, the mass matrix for left-handed charged fermions from the gauge and Higgs supermultiplets is

$$\begin{pmatrix} \tilde{W}^+, \tilde{H}_1^+ \end{pmatrix}_L \begin{pmatrix} \alpha_2 M_0 & g_2 v_2 \\ g_2 v_1 & 0 \end{pmatrix} \begin{pmatrix} \tilde{W}^- \\ \tilde{H}_2^- \end{pmatrix}_L \quad (2)$$

which yields in the limit $m_{\tilde{W}} \gg m_W$ the mass eigenstates

$$\begin{cases} \left(\tilde{W}^+ + \frac{g_2 v_2}{\alpha_2 M_0} \tilde{H}_1^+ \right)_L \\ \left(\tilde{W}^- + \frac{g_2 v_1}{\alpha_2 M_0} \tilde{H}_2^- \right)_L \end{cases} \quad \begin{cases} \left(\tilde{H}_1^+ - \frac{g_2 v_2}{\alpha_2 M_0} \tilde{W}^+ \right)_L \\ \left(\tilde{H}_2^- - \frac{g_2 v_1}{\alpha_2 M_0} \tilde{W}^- \right)_L \end{cases} \quad (3)$$

with masses

$$m_{\tilde{W}^\pm} \approx \alpha_2 M_0, \quad m_{\tilde{H}^\pm} \approx \frac{g_2^2 v_1 v_2}{m_{\tilde{W}}} \quad (4)$$

respectively. In the limit of large $m_{\tilde{W}}$ this latter charged fermion should therefore have a mass considerably less than $m_{\tilde{W}}$, though it must be heavier than 15 GeV to avoid having been seen at PETRA:

$$80 \text{ GeV} \gtrsim m_{\tilde{H}^\pm} \approx \frac{g_2^2 v_1 v_2}{m_{\tilde{W}}} \gtrsim 15 \text{ GeV} \quad (5)$$

Note that because both the positively and negatively charged left-handed fermions, or equivalently both left- and right-handed negatively charged fermions, have weak charges, the neutral weak couplings would be very different from those of a conventional sequential heavy lepton, as shown in Table 1. Hence the Z^0 contributions to their pair-production cross-section in e^+e^- annihilation, and their forward-backward production asymmetries¹¹⁾, would also be distinctively different from those of a conventional heavy lepton.

Turning now to the neutral fermions from the gauge and Higgs supermultiplets, they have the mass matrix

$$(\tilde{W}^3, \tilde{B}^0, \tilde{H}_1^0, \tilde{H}_2^0) \begin{pmatrix} \alpha_2 M_0 & 0 & \frac{1}{2\sqrt{2}} g_2 v_1 & \frac{-1}{2\sqrt{2}} g_2 v_2 \\ 0 & \frac{5}{3} \alpha_1 M_0 & \frac{1}{2\sqrt{2}} g_1 v_1 & \frac{-1}{2\sqrt{2}} g_1 v_2 \\ \frac{1}{2\sqrt{2}} g_2 v_1 & \frac{1}{2\sqrt{2}} g_1 v_1 & 0 & 0 \\ \frac{-1}{2\sqrt{2}} g_2 v_2 & \frac{-1}{2\sqrt{2}} g_1 v_2 & 0 & 0 \end{pmatrix} \begin{pmatrix} \tilde{W}^3 \\ \tilde{B}^0 \\ \tilde{H}_1^0 \\ \tilde{H}_2^0 \end{pmatrix} \quad (6)$$

in the simplest class of supersymmetric models^{2),6),7)}. The neutral gauginos mix with the combination $\tilde{A}^0 \equiv (v_1 \tilde{H}_1^0 - v_2 \tilde{H}_2^0)/v$ of neutral shiggses, and in the limit of large gaugino mass the diagonal form of this matrix is

$$\left. \begin{aligned} \tilde{W}^3 + \frac{g_2 v}{2\sqrt{2} \alpha_2 M_0} \tilde{A}^0 &: m \approx \alpha_2 M_0 \\ \tilde{B}^0 + \frac{3 g_1 v}{10\sqrt{2} \alpha_1 M_0} \tilde{A}^0 &: m \approx \frac{5}{3} \alpha_1 M_0 \\ \tilde{A}^0 \equiv \tilde{A}^0 - \frac{g_2 v}{2\sqrt{2} \alpha_2 M_0} \tilde{W}^3 - \frac{3 g_1 v}{5\sqrt{2} \alpha_1 M_0} \tilde{B}^0 &: m \approx \frac{g_2^2 v^2}{5 \alpha_2 M_0} \\ \tilde{S}^0 \equiv (v_2 \tilde{H}_1^0 + v_1 \tilde{H}_2^0)/v &: m = 0 \end{aligned} \right\} \quad (7)$$

There is no particular symmetry reason why there should be a strictly massless neutral fermion, though cosmology requires that since it is stable, the lightest supersymmetric particle be neutral with mass $< 0(100) \text{ eV}$ or $> 0(1 \text{ to } 10) \text{ GeV}$ ^{9),10)}. In models with an extra singlet superfield Y coupled by

$$P_Y = \lambda H_1 Y H_2 \quad (8)$$

one expects, after symmetry breaking, a term $\epsilon \tilde{H}_1^0 \tilde{H}_2^0$ where $\epsilon \lesssim m_W$, giving a mass matrix for the neutral fermions in the chiral superfields of the form

$$\begin{pmatrix} \tilde{A}^{0'} & \tilde{S}^0 & \tilde{Y}^0 \end{pmatrix} \begin{pmatrix} \frac{g_2^2 v^2}{5\alpha_2 M_0} - \frac{v_1 v_2 \epsilon}{v^2} & \frac{v_1^2 - v_2^2}{2v^2} \epsilon & 0 \\ \frac{v_1^2 - v_2^2}{2v^2} \epsilon & \frac{v_1 v_2 \epsilon}{v^2} & \frac{\lambda}{2} v \\ 0 & \frac{\lambda}{2} v & \mu \end{pmatrix} \begin{pmatrix} \tilde{A}^{0'} \\ \tilde{S}^0 \\ \tilde{Y}^0 \end{pmatrix} \quad (9)$$

where $\mu \lesssim 0(m_W)$. In the limit of small μ and ϵ the mass eigenstates and eigenvalues of (9) are approximately:

$$\tilde{A}^{0'} : m \approx \frac{g_2^2 v^2}{5\alpha_2 M_0} \quad (10a)$$

$$\tilde{S}_\pm^0 \equiv \frac{1}{\sqrt{2}} (\tilde{S}^0 \pm \tilde{Y}^0) : m \approx \frac{\lambda v}{2} \pm \frac{1}{2} \left(\mu + \frac{v_1 v_2 \epsilon}{v^2} \right) \quad (10b)$$

The cosmological constraint⁹⁾ then becomes

$$\frac{g_2^2 v^2}{m_{\tilde{W}}} \gtrsim \min \left(\frac{g_2^2 v^2}{5m_{\tilde{W}}}, \frac{\lambda v}{2} - \frac{1}{2} \left(\mu + \frac{v_1 v_2 \epsilon}{v^2} \right) \right) \gtrsim (1 \text{ to } 10) \text{ GeV} \quad (11)$$

In the limit of large μ the $\tilde{Y}^0$ state in the matrix (9) almost decouples and the lightest neutral fermion again becomes the $\tilde{S}^0$ with

$$m_{\tilde{S}^0} \rightarrow 0 \quad \text{as } \mu \rightarrow \infty \quad (12)$$

and we recover the form of Eq. (7).

Irrespective of how light the extra low-mass singlet chiral field may be, one expects the lightest supersymmetric particle, and hence the end-product

of every decay chain, to be a neutral shiggs particle. It may be quite light, or weigh a few GeV (11) which could show up in the end-point of a decay spectrum, e.g., for $\tilde{H}^+ \rightarrow \tilde{H}^0 + \ell\nu$. It is possible that there are two almost degenerate light shiggses (10b), in which case the slightly heavier one $\tilde{S}_+^0$ could decay into the lighter one $\tilde{S}_-^0$ and a fermion-antifermion pair through Z^0 exchange as shown in Fig. 2a. The lifetime for this decay mode depends critically on $m_{\tilde{S}_+^0}$ and could easily be longer than $O(10^{-13})$ sec. and hence detectable as in Fig. 2b. Indeed, if the mass difference is sufficiently small, the lifetime could be so long [$\gg O(10^{-10})$ sec] that the particle decays outside the apparatus and appears stable. In this case, the masses of the two lightest neutral shiggses would presumably be so close that they could not be distinguished, e.g., by looking at end-point effects in the decays of heavier supersymmetric particles such as the $\tilde{H}^\pm$, and their physics would be hard to distinguish from that of a single massive neutral shiggs $\tilde{S}^0$. The spectroscopy and properties of gauginos and shiggses are summarized in Table 2¹²⁾.

Also shown in Table 2 are the spectroscopy and properties of spin-zero Higgs bosons¹²⁾. Specification of these requires knowledge of the effective scalar potential for the H_1 and H_2 multiplets. The quartic gauge contribution to this potential is of course unambiguous, while quadratic terms of varying signs are generated by radiative corrections^{2),6),7)} as recalled in Fig. 1. A possible additional quartic piece related to the Yukawa interaction (8) has the effect⁶⁾ of stabilizing the potential in such a way as to give non-zero v.e.v.'s v_1 and v_2 if the negative quadratic term in the Higgs potential due to the Yukawa interaction of the top quark is sufficiently large. There are upper limits on the Yukawa couplings λ such as (8) if the absolute minimum of the potential is to conserve charge¹²⁾. The masses shown in the Higgs boson part of Table 2 apply in the limit $\lambda \ll g_2, g_1$, and we see that most model-dependence disappears in this limit. Assuming $v_2 \ll v_1$, which seems plausible since $m_b \lesssim (1/4)m_t$, we see that the sole physical charged Higgs is expected to have a mass $\approx m_{W^\pm}$. In addition to the light gauge-singlet Higgs bosons Y of unknown mass, there are three light neutral Higgs bosons, of which the lightest has a (mass)²:

$$m^2 \approx 4\lambda^2 v_2^2 + \dots \quad (13)$$

It is mainly H_2^0 , so that it has larger couplings to charge $-1/3$ quarks than to charge $+2/3$ quarks, and its couplings to vector bosons are suppressed by a factor (v_2/v_1) relative to those of a conventional single Higgs. There is a heavier ($\approx m_{Z^0}$) neutral scalar Higgs with dominant couplings to charge $+2/3$ quarks. In addition there is a neutral pseudoscalar Higgs which is essentially

the conventional axion a , and in our model it couples more strongly to charge $-1/3$ quarks because we expect $v_2 < v_1$. This particle is expected to acquire a mass of a few GeV as a result of interactions around the grand unification mass^{2),12)}. It is unfortunate that the neutral Higgses kinematically accessible in toponium and Z^0 decays are likely to be relatively weakly coupled to them.

In a model with an additional gauge singlet chiral superfield (8), there are two more light $[< O(m_W)]$ neutral Higgs bosons. Their masses are of order λv but they are likely to be undetectably weakly coupled to conventional matter. Their couplings to light quarks arise from loops involving H_1 and H_2 Higgs fields and are estimated to be of order

$$\left(\frac{\alpha}{\pi}\right) \lambda \left(\frac{m_q}{m_W}\right)^3 \left(\frac{v_2}{v_1}\right)^{1 \text{ or } 2} \leq O(10^{-12}) \lambda \quad (14)$$

In the above discussion we have tried to be as general as possible, but our considerations were prompted by a specific model^{2),12)} which we can also use to make other predictions. In this model the Higgs vacuum expectation values are made possible by negative contributions to the Higgs (mass)², of which the largest is that due to the t quark. For this to work, the t quark must be sufficiently heavy, and we have previously quoted²⁾ a necessary lower bound on its mass. One can in fact give¹²⁾ an explicit expression for m_t as a function of the two unknown parameters of the model $m_{\tilde{W}}$ and the coupling λ of Eq. (8). [L is a logarithmic factor¹²⁾ estimated to be $O(30)$.]

$$m_t^2 + m_b^2 = m_W^2 \left[\frac{\pi}{\alpha_2} \frac{3}{4} \left(\frac{\alpha_2}{\alpha_3} \right)^3 \frac{1}{L} \right] \left\{ 1 + \frac{4\pi m_W^2}{3\alpha_2^2 m_{\tilde{W}}^2 L} \left[2\alpha_\lambda + \left(1 - \frac{2v_2^2}{v^2} \right)^2 (\alpha_2 + \alpha_1 - 2\alpha_\lambda) \right] \right\} \quad (15)$$

where $\alpha_\lambda \equiv \lambda^2/4\pi$, which gives a lower bound of $O(25)$ GeV on m_t .

In this model v_2 is also given in terms of $m_{\tilde{W}}$ and λ : in the limit of large $m_{\tilde{W}}$ it becomes

$$v_2 \rightarrow m_b \sqrt{\frac{2}{3\pi^2} L \left(\frac{\alpha_3}{\alpha_2} \right)^3} \approx 30 \text{ GeV} \quad (16)$$

Turning now to Eq. (5) we see that it is possible to derive from the non-observation of a charged shiggs at PETRA an upper bound on $m_{\tilde{W}}$:

$$m_{\tilde{W}} \lesssim O(3) m_{W^\pm} \quad (17)$$

Using Eq. (1) we then find

$$m_{\tilde{g}} \lesssim O(9)m_{W^\pm}, \quad m_{\tilde{B}} \lesssim O\left(\frac{1}{2}\right)m_{W^\pm} \quad (18)$$

Using the relation of Ref. 2) between the gaugino, slepton and squark masses, we then obtain

$$m_{\tilde{l}} \lesssim O(1)m_{W^\pm}, \quad m_{\tilde{q}} \lesssim O(3)m_{W^\pm} \quad (19)$$

We therefore find¹²⁾ that in this model the supersymmetric partners of the known particles cannot be far beyond our present experimental reach, and may in particular be accessible to LEP.

The best prospects for testing the class of supersymmetric models discussed here seem to be:

- looking for charged fermions with masses $\leq m_{W^\pm}$ which do not have conventional neutral weak interactions;
- looking in their decays, or those of the Z^0 , for evidence of massive stable fermions;
- looking for charged Higgs scalars with masses $O(m_{W^\pm})$;
- looking for a rich and characteristic spectrum of neutral Higgs bosons with masses $\lesssim O(m_Z)$;
- looking for sleptons with masses $\leq O(m_W)$.

This class of models at least has the merit of being readily falsifiable.

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Particle	g_V	g_A
e^-, μ^-, τ^-	$-\frac{1}{2} + 2\sin^2\theta_W \approx 0$	$+\frac{1}{2}$
$\tilde{W}^- + \epsilon \tilde{H}^-$	$-2 - \frac{1}{2} \left(\frac{v^2}{\alpha_2^2 M_0^2} \right) g_2^2 + 2\sin^2\theta_W \approx -\frac{3}{2}$	$\frac{1}{2} \left(\frac{v_1^2 - v_2^2}{\alpha_2^2 M_0^2} \right) g_2^2 \approx 0$
$\tilde{H}^- - \epsilon \tilde{W}^-$	$-1 - \left(\frac{v^2}{\alpha_2^2 M_0^2} \right) g_2^2 + 2\sin^2\theta_W \approx -\frac{1}{2}$	$\left(\frac{v_1^2 - v_2^2}{\alpha_2^2 M_0^2} \right) g_2^2 \approx 0$

Table 1: Neutral Current Couplings of Charged Fermions

Fermions	Mass	Characteristics
$\tilde{W}^\pm + \dots$	$m_{\tilde{W}} + \dots$	Mainly $I = 1$, vector-like neutral weak couplings
$\tilde{H}^\pm + \dots$	$\frac{g_2^2 v_1 v_2}{m_{\tilde{W}}}$	Mainly $I = \frac{1}{2}$, vector-like neutral weak couplings decaying into $\tilde{H}^0 + (\ell\nu \text{ or } \bar{q}q)$
$\tilde{W}^3 + \dots$	$m_{\tilde{W}} + \dots$	Mainly $I = 1$ wino
$\tilde{B}^0 + \dots$	$\frac{5}{3}(\frac{\alpha_1}{\alpha_2})m_{\tilde{W}} + \dots$	Mainly $I = 0$ bino
$\tilde{A}^0 \equiv \frac{v_1 \tilde{H}_1^0 - v_2 \tilde{H}_2^0}{v} + \dots$	$\frac{g_2^2 v^2}{5m_{\tilde{W}}} + \dots$	Mainly $I = \frac{1}{2}$ shiggs decaying into $\tilde{H}^\pm + (\ell\nu \text{ or } \bar{q}q)$
$\tilde{S}^0 \equiv \frac{v_2 \tilde{H}_1^0 + v_1 \tilde{H}_2^0}{v}$ OR $\tilde{S}_\pm^0 \equiv \frac{1}{\sqrt{2}}(\tilde{S}^0 \pm \tilde{Y}^0)$	≈ 0 $\frac{\lambda v}{2} \pm \dots$	Mainly $I = \frac{1}{2}$ shiggs: end of all supersymmetric decay chains Mixture of $I = \frac{1}{2}$ and $I = 0$, lighter one stable, heavier one quasistable decaying into lighter one $+(\ell^+ \ell^- \text{ or } \bar{q}q)$
Bosons	(Mass) ²	Characteristics
$\frac{v_1 H_1^0 + v_2 H_2^0}{v} + \dots$	$4\lambda^2 v^2 + \dots$	$I = \frac{1}{2}$, mainly coupled to charge $-\frac{1}{3}$ quarks
$\frac{v_1 H_1^0 - v_2 H_2^0}{v}$	$\frac{g_2^2 + g_1^2}{2} v^2 + \dots$	$I = \frac{1}{2}$, mainly coupled to charge $+\frac{2}{3}$ quarks
$a \equiv \frac{\eta_1 v_2 - \eta_2 v_1}{v}$	few GeV?	Pseudoscalar coupling, larger for charge $-\frac{1}{3}$ quarks
$\frac{v_1 H_2^\pm - v_2 H_1^\pm}{v}$	$\frac{g_2^2}{2} v^2 + \dots$	Conventional charged Higgs boson
$Y?$	$O(\lambda^2 v^2)?$	$I = 0$ scalar and pseudoscalar, very weakly coupled to light quarks

Table 2: Broken Supersymmetric Spectroscopy

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FIG.1

(a)

(b)

FIG.2