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John Ellis, G. Gelmini, C. Jarlskog*, G.G. Ross** and J.W.F. Valle+

CERN-Geneva

Abstract

In some phenomenological supersymmetric models R-parity (+1 for particles, -1 for sparticles) is spontaneously broken along with tau lepton number L_τ by a vacuum expectation value v_τ of the tau sneutrino ν_τ. To avoid excess stellar energy loss through majorons, there should also be explicit L_τ violation through right-handed neutrinos. To have a sufficiently light ν_τ, either v_τ is very small which is unnatural and boring, and /or the Higgs mixing parameter ε is very small. We find that in the limit $\varepsilon \to 0$:

both the forward-backward asymmetry in $e^+e^- \to \tau^+\tau^-$ and the τ lifetime are unchanged,

$Z^0 \to \tilde{\gamma}\nu_\perp$ decays are possible where $\nu_\perp$ is an extra neutrino

squarks and gluinos may decay into τ or ν_τ

the photino $\tilde{\gamma}$ can decay into $\nu_\perp f\bar{f}$ with a detectable secondary vertex.

Single production of (R-odd) sparticles may occur.

*Present address: Department of Physics, University of Stockholm, Stockholm Sweden.

On leave from: Department of Physics, University of Bergen, Bergen, Norway.

**Present address: Dept of Theoretical Physics, University of Oxford, 1 Keble Road, Oxford, England.

+On leave from: Theory Group, Rutherford Appleton Laboratory, Chilton, Didcot, Oxon OX11 0QX, UK and also University of Brasilia, 70910 Brasilia, DF, Brazil.

Most studies of supersymmetric phenomenology have concentrated on models in which there is an exactly conserved multiplicative quantum number called R-parity, which is + 1 for all conventional particles and -1 for all their supersymmetric partners. In such models sparticles can only be pair-produced, they always decay into a lighter sparticle, and the lighter sparticle (probably the photino $\tilde{\gamma}$) is absolutely stable. However, it is easy to construct models in which R-parity is broken, either explicitly together with lepton number, e.g. by a Higgs-lepton coupling H'L in the superpotential P,[1] or else spontaneously by a sneutrino vacuum expectation value $\langle 0|\tilde{\nu}_{e,\mu,\tau}|0\rangle \equiv v_{e,\mu,\tau} \neq 0$ [2]. Indeed, it has recently been pointed out that $v_\tau = O(\langle 0|H|0\rangle \equiv v)$ is a generic feature of many phenomenological supergravity models,[3] while $v_{e,\mu} \neq 0$ is possible but less likely. In view of the stringent upper limits on L_e and L_μ violation, we concentrate on models with only $v_\tau \neq 0$, in which L_τ is spontaneously broken together with R-parity.

In this paper we explore the phenomenological constraints on models with spontaneous violation of R-parity and L_τ, and propose some experimental signatures.

We start by writing down the mass matrix for the charged fermions mixing with the τ:

$$(\tilde{W}^+, \tilde{H}'^+, \tau^+)_L \begin{bmatrix} M_2 & g_2 v & g_2 v_\tau \\ g_2 v' & \varepsilon & 0 \\ 0 & h_\tau v_\tau & h_\tau v \end{bmatrix} \begin{bmatrix} \tilde{W}^- \\ \tilde{H}^- \\ \tau^- \end{bmatrix}_L \qquad (1)$$

where M_i i=1,2 are the U(1), SU(2)(supersymmetry violating) contributions to gaugino masses, g_i i=1,2 the gauge couplings, and h_τ is the tau Yukawa coupling to the H Higgs. For reasons which will become apparent when we discuss the mixing of neutral fermions, we are interested in the limit $\varepsilon \to 0$. Because eq (1) has zero determinant as $h_\tau \to 0$ there exists a light state, the physical "τ". Expressed in terms of left-handed two-component spinors, we have

$$``\tau" = \left(\tau^+ - \frac{2h_\tau v v_\tau}{g_2 r^2}\tilde{W}^+ , \quad \frac{v\tau^- - v_\tau \tilde{H}^-}{r}\right) \tag{2a}$$

of mass

$$m_\tau = \frac{h_\tau (v_\tau^2 - v^2)}{r} \tag{2b}$$

where

$$r = (v^2 + v_\tau^2)^{\frac{1}{2}} \tag{2c}$$

We will return later to the phenomenological implications of the decomposition (2a) of what we call the physical τ.

The heavier eigenstates can easily be found in two interesting limits: (a) $M_2 << g_2 V$ and (b) $M_2 >> g_2 V$ where

$$V = (r^2 + v'^2)^{\frac{1}{2}}$$

In case (a) we find

$$\tilde{W} \equiv \left(\tilde{W}^+ + \frac{v' M_2}{g_2 (r^2 - v'^2)}\tilde{H}'^+ + \frac{2vh_\tau v_\tau}{g_2 r^2}\tau^+ , \quad \frac{v\tilde{H}^- + v_\tau \tau^-}{r} + \frac{rM_2}{g_2 (r^2 - v'^2)}\tilde{W}^-\right) \tag{3a}$$

$$M_{\tilde{W}} \approx g_2 \sqrt{v^2 + v_\tau^2}$$

$$``\tilde{H}" \equiv \left(\tilde{H}'^+ - \frac{v' M_2}{g_2 (r^2 - v'^2)}\tilde{W}^+ , \quad \tilde{W}^- - \frac{M_2}{g_2 (r^2 - v'^2)}(v\tilde{H}^- + v_\tau \tau^-)\right) \tag{4a}$$

$$m_{\tilde{H}} \approx g_2 v'$$

while case (b) yields

$$``\tilde{W}" \cong \left(\tilde{W}^+ + \frac{v'}{M_2}\tilde{H}'^+ , \quad \tilde{W}^- + \frac{v}{M_2}\tilde{H}^- + \frac{v_\tau}{M_2}\tau^-\right) \tag{3b}$$

$$``\tilde{H}" \cong \left(\tilde{H}'^+ - \frac{v'}{M_2}\tilde{W}^+ , \quad - \frac{v\tilde{H}^- + v_\tau \tau^-}{r} + \frac{r}{M_2}\tilde{W}^-\right) \tag{4b}$$

$$m_{\tilde{H}} \approx g_2^2 \frac{rv'}{M_2}$$

Notice that in case (b) there is a light charged state which would have already been observed at PETRA unless M does not exceed 1 Tev or so. In any event, if it is light enough, it could also be observed in W-decay at the CERN $p\bar{p}$ collider.

The mass matrix for the neutral supersymmetry fermions mixing with the ν_τ is

$$(\tilde{W}^3, \tilde{B}, \tilde{H}^o, \tilde{H}'^o, \nu_\tau) \begin{bmatrix} M_2 & 0 & \frac{-g_2 v}{\sqrt{2}} & \frac{g_2 v'}{\sqrt{2}} & \frac{-g_2 v_\tau}{\sqrt{2}} \\ 0 & M_1 & \frac{g_1 v}{\sqrt{2}} & \frac{-g_1 v'}{\sqrt{2}} & \frac{g_1 v_\tau}{\sqrt{2}} \\ \frac{-g_2 v}{\sqrt{2}} & \frac{g_1 v}{\sqrt{2}} & 0 & -\varepsilon & 0 \\ \frac{g_2 v'}{\sqrt{2}} & \frac{-g_1 v'}{\sqrt{2}} & -\varepsilon & 0 & 0 \\ \frac{-g_2 v_\tau}{\sqrt{2}} & \frac{g_1 v_\tau}{\sqrt{2}} & 0 & 0 & 0 \end{bmatrix} \begin{bmatrix} \tilde{W}^3 \\ \tilde{B} \\ \tilde{H}^o \\ \tilde{H}'^o \\ \nu_\tau \end{bmatrix} \tag{5}$$

For a viable theory there must be a light tau-neutrino ν_τ. It has been recently argued that m_{ν_τ} must be <O(1 MeV) to avoid $e^+e^-\nu$ decay [4]. In this case it is expected that the decay $\nu_\tau \to 3\nu$ will typically proceed only with lifetime greater than the age of the universe [5]. Thus conventional limits on stable neutral relics should apply, unless there are new fast decay modes.

The determinant of this mass matrix is $\varepsilon^2 v_\tau^2(g_2^2M_2+g_1^2M_1)$. There are various limits in which there will be a light neutral state (needed for the ν_τ); namely, $v_\tau \to 0$, or $\varepsilon \to 0$, or $M_2, M_1 \to \infty$. In the latter limit there are two states with mass M_2 and M_1 so the determinant of the remaining masses is $\frac{\varepsilon^2 v_\tau^2(g_2^2M_2+g_1^2M_1)}{M_1M_2}$

The case $v_\tau \to 0$ restores the usual supersymmetric phenomenology, with the exception that the lightest supersymmetric state may be unstable. We will not consider this (boring) limit, but turn to the possibilities if $\frac{v_\tau}{v'} = O(1)$. Since M_1 and M_2 can not be $>>>O(m_W)$ in order to keep the τ neutrino mass $<O(100\text{eV})$, ε must be small [F1]

$$\varepsilon < 100 \text{ eV} \tag{6}$$

In the limit $\varepsilon \to 0$, there are two massless states given by

$$"\nu_\tau" = \frac{1}{r} \, (v\nu_\tau - v_\tau \tilde{H}^o) \tag{7}$$

and

$$\nu_\perp = \frac{1}{V} \, \left\{\frac{v'}{r} \, (v_\tau \nu_\tau + v\tilde{H}^o) + r\tilde{H}'^o\right\} \tag{8}$$

where $V^2 = v^2 + v'^2 + v_\tau^2$ and $r^2 = v^2 + v_\tau^2$.

Here we have used the freedom to choose a basis amongst the massless states so that ν_τ is the SU(2) current eigenstate partner of the τ^- given in eq(2a). This shows that the τ lifetime is unchanged in this limit. For ε small, but non-zero, the mass eigenstates will be mixtures of ν_τ and $\nu_\perp$, but the τ decay rate will be only changed at most at $O(\frac{\varepsilon}{m_\tau})^2$, a negligibly small correction.

The remaining neutral eigenstates are orthogonal to ν_τ and $\nu_\perp$. They are given by

$$\psi_i = N_i^{-1}\left[a_i\tilde{W}^3 + b_i\tilde{B} + v\tilde{H}^o - v'\tilde{H}'^o + v_\tau \nu_\tau\right] \tag{9}$$

where N_i are normalisation factors and a_i and b_i are given by

$$a_i = -\frac{g_2 V^2}{\sqrt{2}(\lambda - M_2)}$$

$$b_i = \frac{g_1 V^2}{\sqrt{2}(\lambda - M_1)} \tag{10}$$

and λ_i is the mass given by the solutions of

$$\lambda(\lambda-M_1)(\lambda-M_2) = \frac{V^2}{2}\left[g_2^2(\lambda-M_1) + g_1^2(\lambda-M_2)\right] \tag{11}$$

Let us consider the solution of eq (10) in various limits. As discussed above the most interesting new possibility is $(\frac{v_\tau}{v'}) = O(1)$ and, indeed, in many models this is its natural value. What is the expectation for the other parameters M_i, v? The ino masses M_i are supersymmetry breaking masses. In supergravity models with non-minimal kinetic energy terms these masses appear at tree level and should be $O(m_W)$. Radiative corrections will split these masses at low scales, so we expect $M_{\tilde{g}} > M_2 > M_1$, where $M_{\tilde{g}}$ is the gluino mass. Even if absent at tree level, in theories with a broken continuous R symmetry M_i will be $O(\alpha_i m_{3/2})$ through radiative corrections. In standard supersymmetric models the vev of H^o is expected to be of $O(v')$. In R parity breaking models this no longer is true and $(\frac{v}{v'})$ can be arbitrarily small. However in order to preserve this pattern of electroweak symmetry breakdown we should have $h_t > h_b$ and so $\frac{v}{v'} > \frac{m_b}{m_t}$.

In table 1 we give the mass eigenstates ψ_i for two interesting limits (a) $\frac{M_i}{g_2 V} \ll 1$ [F2] and (b) $\frac{M_i}{g_2 V} \gg 1$. Before we discuss their phenomenology we must first consider the scalar sector.

If the only violation of τ lepton number is spontaneous ($\langle \nu_\tau \rangle \neq 0$), there exists a corresponding Goldstone boson, the Majoron J [7]. The process $\gamma + e \to J + e$ inside Red Giant stars would lead to excessive energy loss unless $\nu_\tau < 100$ KeV [8]. Such a small though nonzero value of ν_τ is unnatural from the point of view of existing models and boring phenomenologically. Therefore we assume that lepton number is also violated explicitly. One way of doing this is via L-violating superpotential terms such as H'L or $QQ^c L$, or via analogous soft SUSY breaking terms. These possibilities have been considered in the literature [1,2,9], and we prefer an alternative with a long phenomenological pedigree, namely L violation by right-handed neutrino mass terms

$$P \ni m_N NN + \frac{m_f}{v'} NLH' \qquad (12)$$

where m_N and m_f are general matrices in flavor space. In these models it is necessary that the Majoron be heavier than 10 MeV or so, so as not to have been produced in red giant stars whose characteristic temperature is $T < O(1)$ MeV. The contribution from eq (12) to the Majoron mass is $O(\sqrt{m_{3/2} m_\nu})$ where the neutrino mass has the usual form $m_{\nu_\tau} = \frac{m_f^2}{m_N}$. (Thus neutrinos would not decay through Majoron emission). For $m_{\nu_\tau} < 100$ eV, (ie $m_n > 10^9$ GeV for $m_f \sim 10$ GeV) consistent with the cosmological bounds, and $m_{3/2} = 1$ TeV, the majoron mass would be O(10MeV) as required. We have found that in many models there are additional contributions to the Majoron mass coming from induced soft terms of the form $[LH']_A$. For example the superpotential of eq(12) plus an interaction term λN^3 leads, via the graphs of Fig 1, to a term $\propto \frac{\lambda^2}{4\pi^2} m_{3/2}^2 LH'$, even in the limit $m_N \to \infty$. This would generate a Majoron mass $m_J^2 = O(\frac{\lambda^2}{4\pi^2} m_{3/2}^2)$ which is easily of several GeV in magnitude.

In this way one can live with vevs for $\tilde{\nu}_\tau$ of order of v'. In this case J has significant components along the Higgs directions and can be seen in axion searches in beam dump experiments. These exclude J with mass < 10MeV, the same limit as we found from red giant stability.

A second Goldstone boson will arise in models with a Peccei-Quinn U(1) symmetry, under which H and H' transform differently. This axion-like field a will acquire a mass through an $\varepsilon H'H$ superpotential term,

$$m_a = O(\sqrt{\varepsilon\, m_{3/2}}) \qquad (13)$$

With the bound on ε derived above, for $m_{3/2} = 1$TeV, this gives $m_a = 10$MeV, just consistent with beam dump bounds. Thus there is an interesting narrow window for models with an approximate Peccei-Quinn symmetry, in which the pseudo-axion would soon be found while the ν_τ (or ν_1) could play an important role in galaxy formation.

Alternatively, one can construct models which do not possess a Peccei-Quinn U(1) symmetry when $\varepsilon \to 0$, for example by introducing a heavy gauge singlet chiral superfield Y with superpotential couplings

$$P \ni H'H\, Y,\ Y^3 \qquad (14)$$

In this case the pseudo-axion acquires a mass,

$$m_a = O(\frac{\lambda}{2\pi} m_{3/2}) \qquad (15)$$

which can easily be large enough to avoid the bounds discussed previously.

We turn now to the phenomenology of these R parity breaking models. The main difference (cf eq (2a)) is that the physical τ^- is a mixture mostly of the original τ^-, a current eigenstate, and $\tilde{H}^-$, the Higgsino carrying the same weak hypercharge. There are (in addition to ν_e and ν_μ) two light "neutrino" states ν_τ (the SU(2) doublet partner of the τ^-) and the orthogonal combination, ν_J (cf eq(8)). As discussed above, the τ lifetime is essentially unchanged.

$e^+e^-\to\tau^+\tau^-$: The τ^+ has very small mixing $O(m_\tau/m_W)$ with the $\tilde{W}^+$ while the τ^- only has substantial mixing with the $\tilde{H}^-$ (2a), which has the same weak isospin $I_3=-\frac{1}{2}$ as the τ^-. Therefore the Z^o couplings to the τ are indistinguishable from those of the standard model, and the total cross-section and the forward-backward asymmetry in $e^+e^-\to\tau^+\tau^-$ are completely canonical.

Cosmological bounds From eqs (5-8) we see there are two light states which would be in kinetic equilibrium during Big Bang Nucleosynthesis (BBN). [10] This is just consistent with the bound that there should be at most one additional neutrino-like species. The (pseudo) majoron and (pseudo) axion do not contribute to the effective neutrino number since their masses are >O(10) MeV.

W,Z decays The new light neutrino states will also be produced in Z decays potentially in conflict with the UA2 experimental result [11]

$$\Delta N^{\nu}_{UA2} < 0 \quad (90\%CL)\ , < 2 \quad (95\%CL) \tag{16}$$

However this bound is obtained assuming a conventional W width and full strength Z coupling to $\nu_\perp$. In the present model (limit (a)) [F3] this coupling may be substantially smaller than the usual coupling by a factor

$$\frac{g(Z\nu_\perp\nu_\perp)}{g(Z\nu_\tau\nu_\tau)} = \frac{v'^2-r^2}{v'^2+r^2} \tag{17}$$

Moreover, from eq (4) we see that $W\to\tilde{H}\nu_\perp$ may be kinematically possible. Because of this the limit of eq (16) does not directly apply.

Off-diagonal Z decays Since we are mixing states with different I_3 and Y it is possible in principle to have off diagonal Z_o couplings. It is easy to check using eqs (2-4) that the Z^o has no large off-diagonal couplings to charged fermions. However, equations (7-10) tell us that it can have sizeable off-diagonal couplings to neutral fermions. For example, in limit (a),

$$\frac{g(Z^o\tilde{\gamma}\nu_\perp)}{g(Z^o\bar{\nu}_e\nu_e)} = 2\sqrt{2}\,\frac{g_1g_2(M_1-M_2)rv'}{g^3V^3} = O(\frac{m_{\tilde{\gamma}}}{m_Z}) \tag{18}$$

The above result means that the Z^o could have a substantial decay rate into $\tilde{\gamma}+\nu_\perp$.

$\tilde{q}\to q\tau$, $\tilde{q}\to q\nu_\tau$ etc: We see from equation (2) that decays into τ^+ are suppressed by $O(m_\tau/m_W)^2$ coming from the $\tilde{W}^-$ admixture and into τ^- by $O(m_q/m_W)^2$ coming from the $\tilde{q}q\,\tilde{H}^-$ coupling. We see from equation (7,8) that decays into ν_τ or $\nu_\perp$ are suppressed by $O(m_{q_u}/m_W)^2$ or $O(\frac{m_{q_d}}{m_W}\,\frac{v'}{v})^2$ coming from the $\tilde{H}^o$ and $\tilde{H}'^o$ admixture, or by $(\varepsilon/m_W)^2$ coming from the $\tilde{W}^3$ and $\tilde{B}$ admixtures.

We expect $\tilde{q}\to q+\tilde{\gamma}$ (and $\tilde{g}\to q+\bar{q}+\tilde{\gamma}$) to dominate as usual, if kinematically allowed. If not, as is the case in limit (b), the new decay modes may be significant.

$\tilde{\gamma}$ decays It is easy to see from (2) and Table (1a) that there is no charged current coupling of the $\tilde{\gamma}$ to the τ. Therefore $\tilde{\gamma}\to\tau\,(f\bar{f}')$ decays do not occur. However, we saw that, for example in limit (a), there is a $\tilde{\gamma}-\nu_\perp$ neutral current coupling (see equation (18)). Therefore the $\tilde{\gamma}$ can decay into $\nu_\perp(e^+e^-$ or $\mu^+\mu^-$ or $\tau^+\tau^-$ or $q\bar{q}$ or $\bar{\nu}_e\nu_e$ or $\bar{\nu}_\mu\nu_\mu$ or $\bar{\nu}_\tau\nu_\tau)$ with a rate $\Gamma(\tilde{\gamma}\to\nu_\perp\chi)=O(\frac{G_F^2m_{\tilde{\gamma}}^5}{192\pi^3})\,(\frac{m_{\tilde{\gamma}}}{m_Z})^2$. If $m_{\tilde{\gamma}} \leq O(10)$ GeV the corresponding lifetime could easily be $\geq O(10^{-11})$ seconds, providing the useful signature of a separated decay vertex. The $\tilde{\gamma}\to\nu_\perp(\bar{\nu}\nu)$ decays would give a pure missing energy signature analogous to conventional stable photinos. For photinos lighter than O(1 Gev) (this value depends on the charged stau mass) we expect the radiative decay into neutrino plus a photon to be dominant. In this case there are strong cosmological limits on the lifetime. Typically one requires $\tau < O(10$ sec), which implies $m_{\tilde{\gamma}} > O(1\text{MeV})$ so as to avoid photo disassociation of primordially synthesized light elements. (Alternatively, one could have extremely long lifetimes, $\tau > O(10^{24}\text{sec})$ but this would require a ultra light photino, which is unlikely).

Unstable photinos may be the clearest experimental signature of models with R-parity broken in the manner discussed in this paper [F4].

Single production of squarks, gluinos etc.

If R parity is broken, the new supersymmetric (R-odd) states may be singly produced. Typical graphs, are shown in Fig 2.

These modes are suppressed by $O(\frac{m_{q_u}}{M_W})^2$, or $O(\frac{m_{q_d}}{M_W} \frac{v'}{v})^2$ as the τ, ν_τ, $\nu_\perp$ couple only through their $\tilde{H}^o$ and $\tilde{H}^{o'}$ admixture.

For a top quark of O(40 GeV) mass the suppression factor is not large and the production of $\bar{\tilde{t}}\, \tilde{t}\, \nu_\perp$ would be sizeable. If $(\frac{v'}{v})$ is large the production of $\tilde{b}\tilde{b}\, \nu_{\tau \text{ or } \perp}$, $\tilde{b}\tilde{t}\tau^-$ could also be appreciable. It is straightforward to extend this analysis to the production of other supersymmetric states.

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References

[1] L.J. Hall, M. Suzuki, Nucl. Phys. B231, 419 (1984).
[2] C.S. Aulakh, R.N. Mohapatra, Phys. Lett. 119B, 136 (1983).
[3] G.G. Ross, J.W.F. Valle, RAL preprint RAL-84-102
[4] S. Sarkar, ESO-CERN Symposium on Particle Physics and Cosmology, November 1983. S. Sarkar, A.M. Cooper, CERN TH-3976 (1984).
[5] J. Schechter, J.W.F. Valle, Phys. Rev. D25, 774 (1982).
Y. Hosotani, Nucl. Phys. B191, 411 (1981).
[6] J.W.F. Valle Phys. Lett. 131B, 87 (1983)
G.B. Gelmini, J.W.F. Valle, Phys. Lett. 142B, 181 (1984).
[7] Y. Chikashige, R.N. Mohapatra, R.D. Peccei, Phys. Lett. 98B, 265 (1980), Phys. Rev. Lett 45, 1926 (1980).
G.B. Gelmini, M. Roncadelli, Phys. Lett. 99B, 411, (1981).
[8] H. Georgi, S.L. Glashow, S. Nussinov, Nucl. Phys. B193, 297 (1981).
M. Fukugita, S. Watamura, M. Yoshimura, Phys. Rev. Lett. 48, 1522 (1982).
[9] I.H. Lee, Phys. Lett. 138B, 121 (1984).
[10] D. Schramm, G. Steigman, Phys. Lett. 141B, 337 (1984).
[11] P. Bagnaia et al (UA2 collaboration), Zeit fur Physik C, 24, 1 (1984).

Footnotes

[F1] The weaker bound $\epsilon < 160$ MeV (corresponding to the laboratory bound on the mass of the τ-neutrino) holds if neutrinos can decay within the lifetime of the universe. Models exist [6] in which the ν_τ decays into a Goldstone boson, the Majoron, and another light neutrino. In this case additional singlet Higgs fields, whose V.E.V. give large Majorana masses to the right-handed neutrinos should be included. Lepton number would then be only spontaneously broken and the Majoron would be a true Goldstone boson but invisible.

[F2] Notice that in this approximation the zino becomes a Dirac fermion. This is however, not a general feature of these models, as can be seen, eg, in Table 1(b).

[F3] In case (b), $gV << M_1$, there are additional couplings of the Z to neutral fermions.

[F4] It may also be possible to choose parameters in the R-parity broken model in such a way that the photino decays invisibly into a Goldstone boson and a neutrino. In such a case the usual experimental signature for the photino as missing energy would be recovered.

Figure Captions

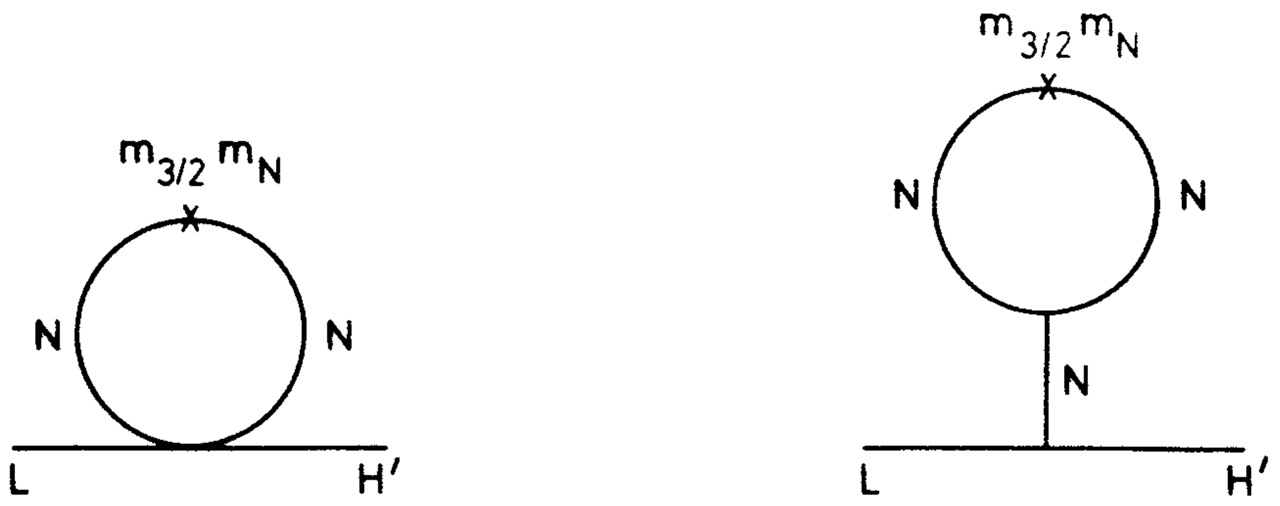
Fig. 1
Graphs giving rise to the soft LH' supersymmetry breaking term.

Fig. 2
Typical graphs leading to single production of supersymmetric (R-odd) states.

(a) gluino production
(b) squark production

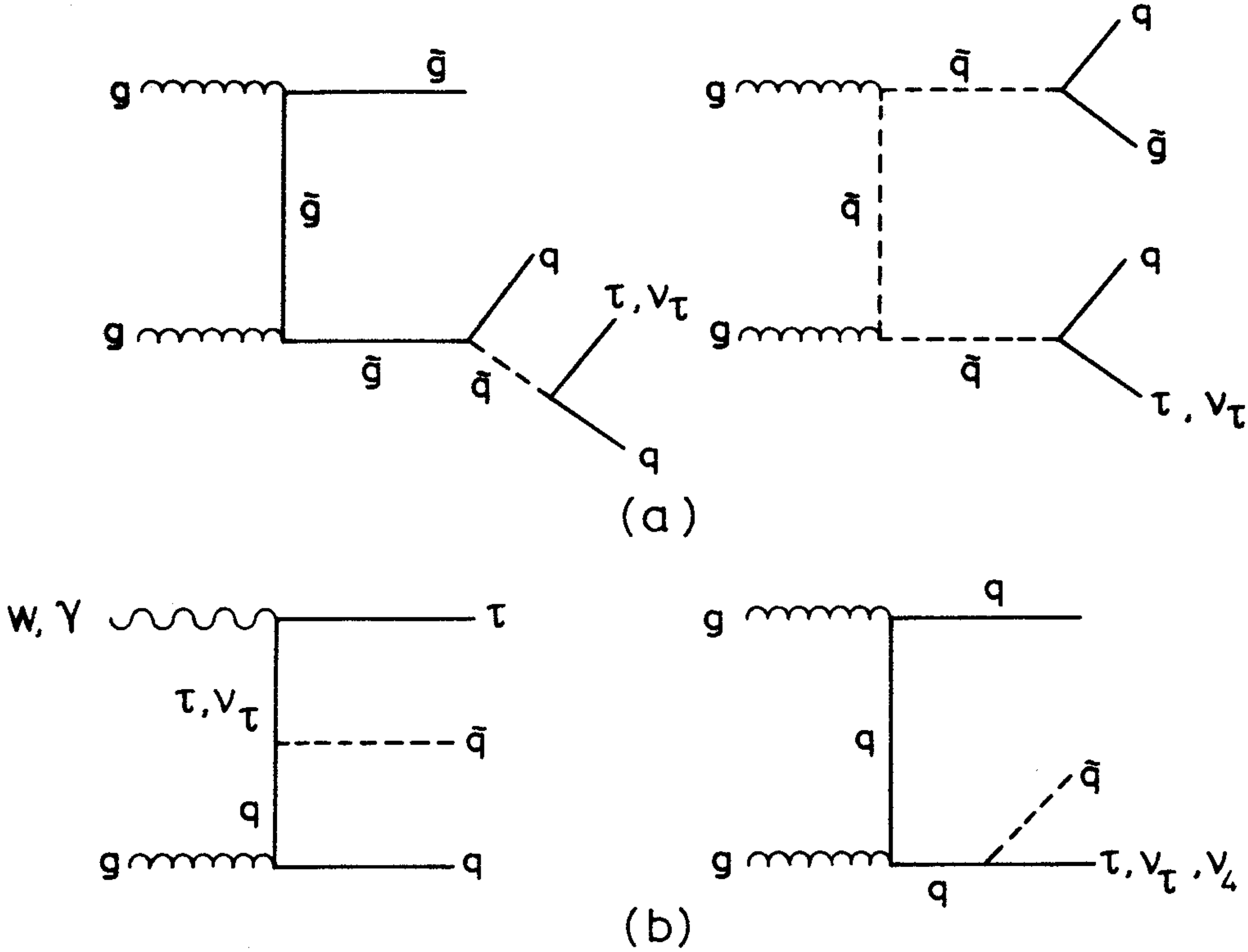
	limit (a) $g_1V >> M_1$		limit (b) $g_1V << M_1$		
component	"$\tilde{\gamma}$"	"$\tilde{Z}_\pm$"	"$\tilde{H}^0$"	"$\tilde{B}$"	"$\tilde{W}_3$"
$\tilde{W}_3$	$\frac{g_1}{g}$	$\pm\frac{g_2}{\sqrt{2}g}$	$\frac{g_2V}{\sqrt{2}M_2}$	$-\frac{g_1g_2V^2}{2M_1(M_1-M_2)}$	1
$\tilde{B}$	$\frac{g_2}{g}$	$\mp\frac{g_1}{\sqrt{2}g}$	$-\frac{g_1V}{\sqrt{2}M_1}$	1	$\frac{g_1g_2V^2}{2(M_1-M_2)M_2}$
$\tilde{\Pi}^0$	$-\frac{\sqrt{2}g_1g_2v(M_1-M_2)}{g^3V^2}$	$-\frac{v}{\sqrt{2}V}$	$\frac{v}{V}$	$\frac{g_1v}{\sqrt{2}M_1}$	$-\frac{g_2v}{\sqrt{2}M_2}$
$\tilde{\Pi}'^0$	$\frac{\sqrt{2}g_1g_2v'(M_1-M_2)}{g^3V^2}$	$\frac{v'}{\sqrt{2}V}$	$-\frac{v'}{V}$	$-\frac{g_1v'}{\sqrt{2}M_1}$	$\frac{g_2v'}{\sqrt{2}M_2}$
ν_τ	$-\frac{\sqrt{2}g_1g_2v_\tau(M_1-M_2)}{g^3V^2}$	$-\frac{v_\tau}{\sqrt{2}V}$	$\frac{v_\tau}{V}$	$\frac{g_1v_\tau}{\sqrt{2}M_1}$	$-\frac{g_2v_\tau}{\sqrt{2}M_2}$
mass	$\frac{g_2^2M_1+g_1^2M_2}{g^2}$	$\pm\frac{gV}{\sqrt{2}}$	$\frac{V^2}{2M_1M_2}(g_1^2M_2+g_2^2M_1)$	M_1	M_2

Table 1 Massive neutral mass eigenstates of eq (5) evaluated in two interesting limits ($g=\sqrt{g_1^2+g_2^2}$)



$m_{3/2} m_N$
N
N
L
H′
$m_{3/2} m_N$
N
N
N
L
H′

Fig. 1



g
g
$\tilde{g}$
$\tilde{g}$
$\tilde{g}$
$\tilde{q}$
q
τ, ν_τ
q
g
g
$\tilde{q}$
$\tilde{q}$
$\tilde{q}$
q
$\tilde{g}$
q
τ, ν_τ
(a)
W, γ
τ
τ, ν_τ
$\tilde{q}$
q
g
q
g
q
q
$\tilde{q}$
g
q
τ, ν_τ, ν_4
(b)

Fig. 2