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Exotic Signatures From Supersymmetry *†

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Abstract

Minor changes to the standard supersymmetric model, such as soft flavor violation and R parity violation, cause large changes in the signatures. The origin of these changes and the resulting signatures are discussed.

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INTRODUCTION

Physicists crave simple frameworks and elegant models which describe a wide variety of phenomena. In the world of supersymmetry this has led to a standard picture: the minimal low energy supergravity model, which will be described in the next section. The vast majority of super-phenomenology is done within this particular model. I find this quite troublesome. Supersymmetry at the TeV scale may well be completely wrong; that does not bother me at all, it is just a basic assumption which we have to make to get started. What troubles me is our nearly blind adherence to what has become the standard supersymmetric model. Our only reason for this particular model is that, to the theorists eye, it seems to be the most economical framework to describe the plethora of new particles and interactions which supersymmetry requires. Economy is a great thing, and I do not have a substitute for this model, however, the crucial point is that apparently innocuous changes in the theory can cause enormous changes in the experimental signatures.

On the other hand it is not a good idea to throw out the standard supersymmetric model and give equal weight to all formulations. One notorious problem of supersymmetry is that without some constraints from model-building you can arrange to get almost any signature you like. In this lecture I would like to start from the standard supersymmetric model, and consider changes in the structure of the model which are quite mild but which I find quite plausible and which have crucial phenomenological consequences. The "exotic" signatures of the title should be understood to be these consequences of changing the assumptions behind the standard model, and should not be taken to be random exotica pulled from a hat.

I will discuss only two such changes, and both have to do with the symmetry structure of the model. It is well known that an $SU(3) \times SU(2) \times U(1)$ gauge symmetry is insufficient symmetry to guarantee proton stability at the weak scale in a supersymmetric model. The usual convention is to add a matter parity symmetry. I will investigate alternative possibilities and find that the most important result is that missing energy signatures at colliders are replaced by events with multi-jets and/or multi-charged isolated leptons.

The second topic is that of flavor physics, which I will deal with only briefly. The decoupling theorem means that heavy particles influence physics at low energies only via the effects they have on renormalizing coupling constants of interactions of the light fields. In the standard model flavor violation occurs only via the Kobayashi–Maskawa matrix. In supersymmetry there are other flavor matrices, but in the standard supersymmetric model it is assumed that these extra matrices are given in terms of the usual Kobayashi–Maskawa matrix. This assumption may be incorrect even if the field content of the low energy theory is unchanged. Extra flavor violation may occur due to the effects of very heavy particles renormalizing the flavor matrices away from the standard form. I will try to convince you that these effects are generic and have important consequences for signatures.

THE MINIMAL LOW ENERGY SUPERGRAVITY MODEL

We must at least define the MLES model¹ which we will be extending. This is a supersymmetric $SU(3) \times SU(2) \times U(1)$ gauge theory with three 15-plets of chiral superfields for “matter”

$$Q(3, 2, 1/6) \quad U^c(\bar{3}, 1, -2/3) \quad D^c(\bar{3}, 1, 1/3) \quad L(1, 2, -1/2) \quad E^c(1, 1, 1)$$

and two for “Higgs”

$$H(1, 2, 1/2) \quad H'(1, 2, -1/2).$$

Supersymmetry itself does not allow for a distinction between matter and Higgs fields, so we impose one by hand: we require the MLES model to be invariant under matter parity under which the matter superfields change sign but the Higgs superfields do not. The most general gauge invariant, renormalizable superpotential is then

$$f = Q\lambda_U U^c H + Q\lambda_D D^c H' + L\lambda_E E^c H' + \mu H H' \quad (1)$$

where λ_U, λ_D and λ_E are 3×3 matrices in generation space and μ is a dimensionful parameter which ensures that the theory does not possess a Peccei-Quinn symmetry.

A useful way of remembering how to get the vertices of the supersymmetric interactions of Equation (1) in terms of component fields is to write down the

usual Yukawa couplings and replace external lines in pairs by their superpartners as in Figure 1. Finally scalar trilinear and quartic interactions are generated by differentiating f , considered as a function of the scalar components of the superfields A_i , as shown in the lower part of Figure 1.

The model that I am describing, despite its name, doesn't have much to do with supergravity. However, a crucial aspect of the model is the structure of the soft supersymmetry breaking operators. These operators can be obtained in a very plausible fashion from supergravity theories (for a detailed discussion and review of the possibilities see reference 2), but we will not need to know anything about supergravity. In the simplest scheme there are four types of soft operators:

$$m^2 A_i^* A_i = m^2 (H^* H + \tilde{q}^* \tilde{q} + \dots) \quad (2a)$$

$$Bm [f_2]_A = Bm \mu H H' \quad (2b)$$

$$Am [f_3]_A = Am \tilde{q} \lambda_U \tilde{u}^c H + \dots \quad (2c)$$

$$-\frac{\tilde{m}}{2} (\tilde{g}\tilde{g} + \tilde{\omega}\tilde{\omega} + \tilde{b}\tilde{b}) \quad (2d)$$

$[f_{2,3}]_A$ are the bilinear, trilinear parts of the superpotential as functions of the scalar components of the superfields A_i . A and B are complex constants with magnitudes of order unity, and $\tilde{g}, \tilde{\omega}, \tilde{b}$ are the gaugino fields for $SU(3), SU(2), U(1)$ gauge groups. Note that H sometimes refers to a superfield and sometimes to its scalar component. Superpartners are differentiated from the standard model particles by a tilde.

Is the MLES just an irrelevant extension of the standard model which has introduced five new parameters and a host of new particles for naught? Experiment must decide. It is theoretically attractive because, unlike the standard model, the theory has no quadratic divergences. All parameters scale according to well behaved renormalization group equations (RGE). If we write a supersymmetric theory with a supersymmetry breaking scale much above the weak scale, then on integrating out the superpartners we will recover the standard model as a low energy effective theory with its quadratically divergent Higgs mass. To prevent this, the supersymmetry breaking and weak scales should be comparable. The MLES model incorporates this automatically: if m and \tilde{m} are made very large the Higgs boson will decouple, hence m and \tilde{m} cannot be

made larger than the weak scale v : m and \tilde{m} are taken to be $0(v)$. Similarly if $\mu \gg v$ the Higgs decouples. The most puzzling feature of MLES is why the mass parameters in the supersymmetric and supersymmetry breaking vertices have comparable sizes.

How does electroweak symmetry breaking occur in this model? Apparently the Higgs mass squared is positive as given in equation 2a. However, RG scaling of the H mass squared parameter due to the large top Yukawa coupling makes it negative by the weak scale. The term in (2b) then induces a linear term in the H' field which consequently also gets a vev. In the MLES theory all component vertices have an even number of superpartner fields. This means the theory possesses a symmetry under which the sign of these fields is reversed: R parity. It has the consequence that the lightest superpartner (LSP) is stable.

R parity is a discrete subgroup of a continuous $U(1)$ R symmetry which rotates the coordinate θ of superspace. The superfields can be expanded in terms of θ and component fields, for example

$$Q = \tilde{q} + \theta q + \dots, \quad (3a)$$

$$H = \tilde{H} + \theta \tilde{H} + \dots, \quad (3b)$$

$$Z = \dots + \theta \sigma^\mu \bar{\theta} Z_\mu + \bar{\theta} \bar{\theta} \tilde{Z} + \dots \quad (3c)$$

You can now check very simply that R parity in MLES, which reverses the sign of θ and all superpartner component fields, is exactly the same thing as matter parity, which reverses the sign of matter superfields, but not Higgs or vector superfields.

R parity plays a central role in the phenomenology of the MLES model, and largely determines the nature of its experimental signatures. This is for two reasons:

- i) R parity implies that direct production of superpartners will occur in pairs.
- ii) Once a superpartner has been made you can never get rid of it (except for the possibility that it might come across another superpartner to annihilate). This is important cosmologically since relic superpartners from the big bang will decay to products which include the LSP, and since the LSP is stable it could be the dark matter. To avoid cosmological problems the LSP should therefore be

neutral: a neutralino or sneutrino. The stability of the LSP is also crucial for lab. searches for supersymmetry. Once produced in a high energy collision, a neutral, stable LSP will escape the apparatus and leave a missing energy signature. The vast majority of searches for supersymmetry, and limits on superpartner masses, have used this signature³.

R PARITY BREAKING

Since the conservation of *R* parity plays such a central role in present thinking in supersymmetric models, a central theme of this set of lectures will be to challenge this standard viewpoint and to explore possibilities for *R* parity breaking. I will restrict myself to discussing models with minimal field content (i.e. as in MLES) and which have explicit *R* parity violation in the renormalizable superpotential at the weak scale.* The only possible gauge invariant, supersymmetry, *R* violating operators in these models are those which violate lepton number

$$f_{\Delta L} = \lambda L E^c L + \lambda' Q D^c L + \mu' L H \quad (4)$$

and those which violate baryon number

$$f_{\Delta B} = \lambda'' U^c D^c D^c. \quad (5)$$

Hence there are only four logical possibilities for models, as shown in Table I.

Table 1

	Excluded	MLES	$\Delta L \neq 0$	$\Delta B \neq 0$
$LE^c L, QD^c L$	✓	X	✓	X
$U^c D^c D^c$	✓	X	X	✓

Most theorists would probably have opted for the model in the first column of Table I, since it contains all possible gauge invariant interactions, and this seems most natural. However, this possibility is excluded because the proton decays with a weak decay rate. The next simplest version is to assume that neither is present; this produces the MLES model shown in the second column.

*This excludes the cases of *R* parity violation via sneutrino vevs⁴ at or beneath the weak scale, and via higher dimension operators.⁵

It might be argued that since unification generally treats quarks and leptons on equal footings this MLES model is to be preferred to either the “ $\Delta L \neq 0$ ” or “ $\Delta B \neq 0$ ” models. In the rest of this section I will show that this argument is false. Those who are interested in the signatures of the “ $\Delta L \neq 0$ ” and “ $\Delta B \neq 0$ ” models can skip to the next section.

The argument that unification conflicts with the “ $\Delta L \neq 0$ ” and “ $\Delta B \neq 0$ ” models can be phrased in $SU(5)$ notation. The matter representations occur in a ten $T(Q, U^c, E^c)$ and five-bar $\bar{F}(D^c, L)$. The problem is that the interaction $T\bar{F}\bar{F}$ contains both B and L violating terms and hence would lead to the first column in Table 1, which is excluded. There are two perfectly acceptable ways to evade this. The first is to try a different gauge group, for example in flipped $SU(5)$ the representations are $T(Q, D^c, N^c), \bar{F}(U^c, L)$ and E^c where N^c is a $SU(3) \times SU(2) \times U(1)$ singlet. Introducing a 10-plet of Higgs, H_{10} , which acquires a vev, the interaction $H_{10}\bar{F}\bar{F}E^c$ contains LLE^c but not $U^cD^cD^c$.⁶

Even simpler is to arrange for the grand unified theory to possess a discrete symmetry (other than R parity) which allows $\bar{F}H_5$ (where H_5 is a 5-plet of Higgs containing the doublet H) and the usual Yukawas but forbids everything else.⁷ The low energy superpotential then contains LH as well as the usual $H'H$ term. Note that the field which actually acquires the weak vev is the linear combination of L and H' which couples to H (recall that the H mass squared is driven negative by the large top quark Yukawa coupling in the RGE, and the bilinear term in f then determines which combination of L and H' acquires a linear term and a vev). Identifying the true lepton fields by rotating to a new doublet basis, in which L no longer have bilinear terms in f , induces $Q\lambda_D D^c L$ and $L\lambda_E E^c L$ terms. This is a very simple way of generating the “ $\Delta L \neq 0$ ” model, and furthermore the L violating operators have a flavor structure which is related to the usual Yukawa interactions.

A simple variant of this scheme follows from actually having the LH term produced by a spontaneous breaking of matter parity at the weak scale. Consider an $SU(5)$ theory which has a gauge singlet matter multiplet N in addition to three generations of T, \bar{F} . The most general matter parity invariant superpo-

tential is

$$f = T\lambda_1 TH_5 + T\lambda_2 \bar{F}H'_5 + \mu H_5 H'_5 + mNN + \lambda_3 \bar{F}NH_5 \quad (6)$$

where m is to be taken comparable to μ . The low energy theory is now that of the minimal model given in equation (1) together with the interactions of N : $mNN + \lambda_3 LNH$. If λ_3 is quite large it will appear in the RGE for the scalar mass-squared parameter for \tilde{N} and induce $\langle \tilde{N} \rangle \neq 0$. This spontaneously breaks R parity and induces the LH term, which after rotation, gives LLE^c and QD^cL .

This shows, in perhaps as clear a way as possible, that unification does not really favor the MLES theory from the " $\Delta L \neq 0$ " model. The reason for the different behavior of B and L , or of quarks and leptons, can be traced to the fact that H_5 has been split by the $SU(5)$ breaking into superheavy triplets and light doublets. Had the triplets been light (which of course leads to disastrous B and L violation via λ_1 and λ_2) the $\langle \tilde{N} \rangle$ vev would cause mass mixing of Higgs triplets and quarks generating B violation. In supersymmetry the missing partners mechanism can split the triplet from H and can therefore be expected to allow L violation but not B violation.

COLLIDER SIGNATURES OF R PARITY VIOLATION

A general discussion of experimental signatures of R parity violation is impossible; there are simply too many parameters to keep track of. As usual there are the supersymmetry breaking parameters of equation (2) which determine the LSP and the spectrum of heavier superpartners. Usually one arranges for the LSP to be neutral, either a gaugino-Higgsino combination $\tilde{\chi}(\tilde{\gamma}, \tilde{Z}, \tilde{H}^0, \tilde{H}^{0'})$ or a sneutrino ($\tilde{\nu}$). This is because it is believed that a charged stable LSP is cosmologically excluded. With R parity violation the LSP is unstable so that the cosmological argument no longer applies; it is necessary to rethink the likelihood of the various LSP candidates. QCD radiative corrections tend to make colored particles heavier than those without color; hence I would expect the gluino to be the heaviest gaugino and squarks to be heavier than sleptons. The LSP is therefore most likely to be neutral as before, $\tilde{\chi}$ or $\tilde{\nu}$, or the charged versions $\tilde{\chi}^\pm$ or $\tilde{\ell}^\pm$.

In addition to the uncertainty in the superpartner spectrum, there is the question of the size and flavor structure of the Yukawa parameters which de-

scribe the ΔB and ΔL violation. For the “ $\Delta B \neq 0$ ” model there are six such parameters in λ'' , and for “ $\Delta L \neq 0$ ” model there are fifteen in λ and λ' . Infact, experiments provide quite severe constraints: $\Delta B \neq 2$ processes such as neutron oscillation and ^{16}O decay implies $\lambda''_{112} \lesssim 10^{-6}$, and lepton number violating process such as $\mu \rightarrow e\gamma$ lead to severe bounds on the λ and λ' as well. Indeed you might guess that all these parameters must be very small. This is incorrect, for example λ'_{333} and λ''_{223} can be $O(1)$. In figuring out how large the various coefficients can be, the following rules of thumb are useful.

- i) B violation amongst quarks of higher generation is fairly harmless, while that amongst light quarks is deadly.
- ii) If just one element of λ or λ' is large (with all other small) then it can be very large ($\simeq .1$). This is true for any element except $\lambda_{331}, \lambda'_{331}$. This is because the resulting four light fermion operators conserve lepton number. The limit of about .1 applies to many¹⁰ but certainly not all coefficients.
- iii) If more than one element of λ and λ' is large then the constraints may be extremely powerful if they induce processes such as $\mu \rightarrow e\gamma$. This gives a strong limit on the product $\lambda_{112}\lambda_{221}$, for example.
- iv) It may be possible to arrange for many λ, λ' to be non-negligible providing they violate only one individual lepton number. For example, suppose that R parity violation has its origin in the operator $L_3 H$. Electron and muon number are conserved, and the L_3/H' rotation induces $L_i E_i^c L_3$ and $Q_i D_i^c L_3$.

In the rest of this section I will illustrate the signatures to be expected in e^+e^- and hadron colliders in the “ $\Delta L \neq 0$ ” model.⁸ There are other signatures of R parity violation that I will not discuss.⁹ There is a great variety of signatures, depending on which elements are large and the superpartner spectrum. My examples will illustrate how spectacular the events can be, and will not be exhaustive. Infact, for simplicity I will restrict my attention to the case where the LSP is either $\tilde{\nu}$ or $\tilde{\chi}$ (which I will think of as having roughly equal $\tilde{\gamma}, \tilde{Z}, \tilde{H}$ and \tilde{H}' components). It will also be clear to you that in several cases existing data

places limits on the masses and couplings. I will not try to give present bounds since I expect the picture to change enormously over the next year, and my main aim is to alert experimentalists that their data may reveal supersymmetry in an unexpected way.

At e^+e^- colliders superpartners can be created singly ($e^+e^- \rightarrow \tilde{\nu}, e^+e^- \rightarrow \tilde{\chi}\nu$) in pairs or via Z decay ($Z \rightarrow \tilde{\chi}\tilde{\chi}, \tilde{\nu}\tilde{\nu}, \tilde{\nu}\ell^+\ell^-, \dots$). For the sneutrino resonance the signature depends on whether $\tilde{\nu}$ is the LSP so $\tilde{\nu} \rightarrow e^+e^-, \mu^+\mu^-$ or if $\tilde{\chi}$ is whence $\tilde{\nu} \rightarrow \tilde{\chi}\nu, \tilde{\chi} \rightarrow \ell^+\ell^-\nu$. In the former case you could see a peak in Bhabha scattering more spectacular than the Z

$$\left(\frac{e^+e^- \text{ event rate at } \tilde{\nu} \text{ peak}}{e^+e^- \text{ event rate at } Z \text{ peak}} \right) \simeq 25 \left(\frac{100 \text{ GeV}}{m_{\tilde{\nu}}} \right) \left(\frac{250 \text{ MeV}}{\Delta E} \right) \left(\frac{\lambda}{.1} \right)^2. \quad (7)$$

The latter case gives two charged leptons with significant missing energy. The cross-section is again $\simeq 10^3 \left(\frac{\lambda}{.1} \right)^2$ units of R .⁸ A similar signature occurs even if the $\tilde{\nu}$ is very heavy since $e^+e^- \rightarrow \tilde{\chi}\nu$ can occur directly.

Direct $\tilde{\nu}$ or $\tilde{\chi}$ production in e^+e^- requires a large $LE^c L_i$ operator. It may be that this is suppressed by the same chiral symmetry that makes the electron light. In this case the most interesting possibilities at e^+e^- machines occur in Z decays (or perhaps via direct double superpartner production $e^+e^- \rightarrow \tilde{\chi}\tilde{\chi}, \tilde{\nu}\tilde{\nu}^*$). If kinematically allowed, $Z \rightarrow \tilde{\chi}\tilde{\chi}^{10}$ and $Z \rightarrow \tilde{\nu}\tilde{\nu}^*$ could have $\simeq 1\%$ branching ratios (in the case of $\tilde{\chi}$ via its \tilde{H} component). The production rate is independent of the size of λ or λ' which now effect the signature via the decay:

$$\tilde{\nu} \rightarrow \ell^+\ell^-, \bar{q}q \text{ or } \tilde{\chi} \rightarrow \ell^+\ell^-\nu, \bar{q}q\ell^\pm, \bar{q}q\nu.$$

giving many interesting signatures.

A very important question in these signatures is the lifetime of $\tilde{\nu}$ or $\tilde{\chi}$. If λ, λ' were extremely small they would escape the detector before decay and these models become similar in their signatures to the MLES. An order of magnitude estimate of the LSP decay rates is

$$\Gamma_{\tilde{\nu}} \simeq \frac{\lambda^2}{8\pi} m_{\tilde{\nu}} \quad (8a)$$

$$\Gamma_{\tilde{\chi}} \simeq \frac{\lambda^2 e^2}{192\pi^3} \frac{m_{\tilde{\chi}}^5}{m^4} \quad (8b)$$

where m is the mass of the relevant exchanged scalar. Thus the decay vertices should be separated from the production vertex by distance

$$d_{\tilde{\nu}} \simeq 10^{-10} \gamma \beta \text{ cm} \left(\frac{10^{-2}}{\lambda} \right)^2 \left(\frac{50 \text{ GeV}}{m_{\tilde{\nu}}} \right) \quad (9a)$$

$$d_{\tilde{\chi}} \simeq 10^{-4} \gamma \beta \text{ cm} \left(\frac{10^{-2}}{\lambda} \right)^2 \left(\frac{50 \text{ GeV}}{m_{\tilde{\chi}}} \right)^5 \left(\frac{m}{100 \text{ GeV}} \right)^4 \quad (9b)$$

where β is the LSP speed. Over most of parameter space the $\tilde{\nu}$ will not give a gap; however $\tilde{\chi}$ decays will give gaps as λ becomes small and m large.

The character of the signals at hadron colliders is similar.¹¹ They fall into the same three groups: W/Z decays, continuum pair production and single superpartner production. These are shown in Table 2 together with the signals at e^+e^- colliders. Clearly there are a very large number of signatures. This is especially true when cascade decays of one superpartner to a lighter one are considered. For example, in the resonance production of a slepton there is the possibility that $\tilde{\ell} \rightarrow \bar{q}q$, giving a bump in the two jet cross-section, and there is also the possibility of a cascade decay $\tilde{\ell} \rightarrow \ell \tilde{\chi}$ followed by $\tilde{\chi} \rightarrow \bar{q}q\ell$ (via QD^cL) or $\tilde{\chi} \rightarrow \ell\ell\bar{e}$ (via LE^cL). Rather than discuss all signatures (which are best figured out from the table) I choose to discuss three possibilities which seem to me especially probable and significant. More details on these and other hadron collider signatures can be found in Reference 11. The greatest hope is for the single production since it gives the possibility of probing large masses. However, if the relevant λ' is small the rate will be too low to observe, since the cross-section is proportional to λ'^2 . However in this regard high energy hadron colliders are more promising than e^+e^- machines. If λ'_{11i} is too small, then it is still possible to use sea quarks and have a rate proportional to $|\lambda'_{22i}|^2$ or $|\lambda'_{33i}|^2$. The price paid for using sea quarks is not large at high energies, and the rate is large anyway. In Figure 2 the cross-section for $\bar{p}p \rightarrow \tilde{\nu}$ is plotted for $\lambda'_{11i} = 1$. For $\sqrt{s} = 2 \text{ TeV}$ and $m_{\tilde{\nu}} = 100 \text{ GeV}$, $10^5 \lambda'^2$ events would result from a 10 pb^{-1} dataset.

Although λ'_{111} and λ'_{112} are constrained to be less than .1, λ'_{113} could be as large as 1 so that one could expect up to 10^5 events in such a run. If the $\tilde{\nu}$ decays back into $\bar{q}q$ then the signature is a two jet event with invariant mass

$m_{\tilde{\nu}}$ the process is illustrated in Figure 3. The crucial question is: can this be seen above the QCD two jet background? To get a feel for this the QCD two jet differential cross-section for $\sqrt{s} = 2$ TeV has been plotted in Figure 4, together with the peak of the resonant sneutrino production cross-section $d\sigma_{\tilde{\nu}}/dM$ (at $M = m_{\tilde{\nu}}$). The signal is roughly .1 of the background. What luminosity L would be required to see a bump in an energy bin of size $m_{\tilde{\nu}}/10$ which is five times the statistical uncertainty in the background? The number of background events in this bin is

$$B = \frac{m_{\tilde{\nu}}}{10} \frac{d\sigma_{jj}}{dM} L \quad (10)$$

whereas the signal is

$$S = \frac{\Gamma_{\tilde{\nu}}}{2} \frac{d\sigma_{\tilde{\nu}}}{dM}(M = m_{\tilde{\nu}})L. \quad (11)$$

Using the result that $\frac{d\sigma_{\tilde{\nu}}}{dM} \simeq .1 \frac{d\sigma_{jj}}{dM}$ over the range of interest we find that $S/\sqrt{B} > 5$ implies

$$L > .1 pb^{-1} \left(\frac{1}{\lambda'}\right)^4 \left(\frac{m_{\tilde{\nu}}}{100 GeV}\right)^{2.7} \quad (12)$$

where we have used $\Gamma(\tilde{\nu} \rightarrow \bar{q}q) = 3\lambda'^2 m_{\tilde{\nu}}/16\pi$ and $\sigma_{\tilde{\nu}}(\sqrt{s} = 2TeV) \simeq 8nb \lambda'^2 \left(\frac{100 GeV}{m_{\tilde{\nu}}}\right)^{2.7}$ for the range $50 GeV < m_{\tilde{\nu}} < 250 GeV$. We conclude that this is only a viable signature if λ' is close to unity, but in this case it may be feasible to search up to quite high $m_{\tilde{\nu}}$.

Much easier is the case when the $\tilde{\nu}$ has a cascade decay via a gaugino: $\tilde{\nu} \rightarrow \tilde{\chi}^0 \nu$ or $\tilde{\chi}^{\pm} \ell^{\mp}$. In this case the gaugino could decay $\tilde{\chi} \rightarrow \bar{q}q\nu, \bar{q}q\ell^{\pm}$ (if QD^cL dominates) or $\tilde{\chi} \rightarrow \ell^{\pm}\nu\nu, \ell^{\pm}\ell^{\mp}\nu, \ell^{\pm}\ell^{\mp}\ell^{\pm}$ (if LE^cL dominates) giving events with up to four isolated charge leptons. We simply do not know if the $\tilde{\chi}$ are lighter or heavier than the $\tilde{\nu}$. If they are lighter, then for λ' smaller than the electroweak gauge couplings, this cascade will be the dominant decay. Thus for $\lambda' \simeq .1$, a $10 pb^{-1}$ run at $\sqrt{s} = 2$ TeV will yield 700 events for $m_{\tilde{\nu}} \simeq 250$ GeV. This is clearly a very powerful probe!

If the $\tilde{\chi}$ is the LSP why not simply produce it directly? This can certainly be done, but because it is not a resonance production the cross-section is not so large. Consider, for example, the t and u channel squark exchange diagrams for $u\bar{d} \rightarrow \ell^+ \tilde{\chi}^0$ in the QD^cL model. Suppose $\tilde{\chi}^0 = \beta\tilde{\gamma} + \dots$ then the contribution

to this process via the photino component of the state is

$$\hat{\sigma}(u\bar{d} \rightarrow \ell^+ \tilde{\chi}^0) \simeq \frac{5\pi}{324} \frac{\alpha\alpha_\lambda}{m_{\tilde{q}}^4} \beta^2 \left(1 - \frac{m_{\tilde{\chi}}^2}{\hat{s}}\right)^2 \left(\hat{s} + \frac{m_{\tilde{\chi}}^2}{2}\right) \quad (13)$$

where $\alpha_\lambda = \lambda^2/4\pi$ and we have taken $m_{\tilde{q}}^2 \gg \hat{s}, m_{\tilde{\chi}}^2$, where \hat{s} is the parton center of mass energy squared. This is a reasonable limit to study: here we are taking $\tilde{\chi}$ to be the LSP, it may be very much lighter than the scalar superpartners if there is an approximate continuous R symmetry.

The parton cross-section of (13) is now folded with the $u\bar{d}$ luminosities to get a $\bar{p}p \rightarrow \tilde{\chi}^0 \ell^+ \dots$ cross section.

$$\sigma(\bar{p}p \rightarrow \tilde{\chi}^0 \ell^+ \dots) = \int \frac{d\hat{s}}{\hat{s}} \frac{5\pi}{324} \alpha\alpha_\lambda \beta^2 \left(\frac{\hat{s}}{m_{\tilde{q}}^2}\right)^2 \left(1 - \frac{m_{\tilde{\chi}}^2}{\hat{s}}\right)^2 \left(1 + \frac{m_{\tilde{\chi}}^2}{2\hat{s}}\right) \left(\frac{\tau}{\hat{s}} \frac{dL}{d\tau}\right)_{u\bar{d}} \quad (14)$$

Instead of doing this numerically, I will do a very rough, but useful and simple, estimate. As an example I'll take: $m_{\tilde{\chi}} = 100$ GeV, $m_{\tilde{q}} = 300$ GeV and assume the region around $\sqrt{\hat{s}}$ of 300 GeV dominates the integral, at which point $\left(\frac{\tau}{\hat{s}} \frac{dL}{d\tau}\right)_{u\bar{d}} = 1nb$ for $\sqrt{s} = 2$ TeV $\bar{p}p$ collisions. Hence I estimate

$$\sigma(\bar{p}p \rightarrow \tilde{\chi}^0 \ell^+ \dots) \simeq 10^{-4} \lambda^2 \beta^2 nb \quad (15)$$

giving $\simeq \lambda^2 \beta^2$ events in a run of $10 pb^{-1}$. The best hope is if λ, β are both close to unity and $\tilde{\chi}^0$ decays via the LE^cL operator giving an event with three isolated charged leptons and some missing transverse energy. However, the main point is the low event rate compared with resonant scalar production.

As a final example I discuss gluino pair production via QCD. The cross-section is shown in Figure 5 for $\bar{p}p$ collisions. It is clearly very large giving 10^3 events for $\sqrt{s} = 2$ TeV, $m_{\tilde{g}} = 100$ GeV and a $10pb^{-1}$ run. The gluino is not expected to be the LSP, hence we expect cascade decays to dominate: $\tilde{g} \rightarrow q\bar{q}\tilde{\chi}$ followed by $\tilde{\chi}$ decay. Now the important point is that the exotic signatures from $\tilde{\chi}$ decay are not dependent on λ being large; the production was $O(\alpha_s^2)$. Thus even if there are only very small LE^cL coefficients the events will have 2–6 isolated charged leptons (depending on how many $\tilde{\chi}$ are $\tilde{\chi}^0$ and how many $\tilde{\chi}^\pm$).

SOFT FLAVOR VIOLATION

If you study equations (1) and (2) you will discover that in the standard supersymmetric model individual lepton numbers are conserved, and the only quark flavor violation occurs via the usual Kobayashi–Maskawa matrix K . Since this is true at tree level, all radiative flavor breaking will be proportional to powers of K . For example renormalization group scaling of the down squark mass matrix via the diagram of Figure 6 introduces $\Delta m_{dL}^2 \propto K^+ m_u^2 K$. These effects are well-known and have been exhaustively studied.

Suppose we add to the minimal model some extra fields X which have trilinear couplings in the superpotential to some of the matter fields $M(Q, U^c, D^c, L, E^c)$ such as $\zeta X X M$ or $\eta X M M$.¹³ In this case diagrams such as the one shown in Figure 7 induce flavor changing scalar masses for the field M proportional to the flavor parameters $\zeta^+ \zeta$ and $\eta^+ \eta$. Most important of all: these soft flavor violations of the low energy theory result even if the X fields are superheavy. This suddenly makes it extremely plausible that no matter what the ultimate high energy theory is, some non-standard soft flavor violation is likely to creep into the low energy theory. If these effects are ever discovered, it is possible that they will give us a window into physics of superheavy mass scales.

As an example, consider $X = H_3$ the superheavy Higgs triplets of $SU(5)$. In this case the superpotential contains

$$f = T \lambda_1 T H_3 + T \lambda_2 \bar{F} \bar{H}_3.$$

Surprisingly this results in individual lepton number violation, unsuppressed by powers of the grand unified scale. To see this note that the charged lepton mass matrix is proportional to λ_2 , and work in a basis where this is diagonal. The matrix λ_1 is non-diagonal and as well as leading to up quark masses contains $U^c \lambda_1 E^c H_3$, which leads to $\Delta m_{e\tau}^2 \propto K^+ m_u^2 K$. Resulting signatures are unfortunately too small to see in this case: $B(\mu \rightarrow e\gamma) \simeq 10^{-15}$.¹³ However, in other models there is no reason why the flavor matrix which appears will be K (in this case it is because the basis which diagonalizes the charged lepton also diagonalizes the down quarks). It has been pointed out¹⁴ that in flipped $SU(5)$ one gets $\Delta m_{e\tau}^2 \propto K'^+ m_u^2 K'$ where K' is a completely independent flavor matrix. This occurs because in flipped $SU(5)$ the charged lepton masses and down

quark masses arise from completely different operators. Similar effects are to be expected for hadronic flavor mixing, and this could be most important for K and B physics.

It could be argued that if non-standard model flavor violation in $\mu \rightarrow e\gamma$, $B^0 - \bar{B}^0$ mixing etc, are observed it is hardly a unique signature of these soft flavor violations in supersymmetry. This is absolutely true; it isn't a unique signature for anything. To discover and confirm the existence of supersymmetry itself will require very many separate measurements to explore the spectrum and couplings. Nevertheless, these flavor violations are a generic effect in supersymmetry, and they could become a significant probe for interactions at very high energies.

CONCLUSIONS

Let me just reiterate my conclusions about R parity violation at the weak scale. The TeV scale theory must have some symmetry to forbid proton decay. This could equally well be B, L or R . Much work has focussed on R invariant theories. There is no real justification for this. The " $\Delta L \neq 0$ " model is also very easy to obtain from more unified schemes. My list of best signatures for the discovery of the " $\Delta L \neq 0$ " model over the next year or two is

i) At e^+e^- colliders:

$$e^+e^- \rightarrow \tilde{\nu}$$

$$Z \rightarrow \tilde{\nu}^*\tilde{\nu}, \tilde{\chi}\tilde{\chi}$$

ii) At $\bar{p}p$ colliders

$$\bar{p}p \rightarrow \tilde{\nu}, \tilde{\ell} \dots$$

$$\bar{p}p \rightarrow \tilde{\chi} \dots$$

$$\bar{p}p \rightarrow \tilde{g}\tilde{g}, \tilde{g} \rightarrow q\bar{q}\tilde{\chi}$$

and in all cases the possibilities

$$\tilde{\nu} \rightarrow \ell^+\ell^-, \bar{q}q, \nu\tilde{\chi}^0, \ell\tilde{\chi}^+$$

$$\tilde{\chi}^0 \rightarrow \nu\ell^+\ell^-$$

$$\tilde{\chi}^\pm \rightarrow \ell^\pm\nu\nu \text{ or } \ell^\pm\ell^\mp\ell^\pm$$

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REFERENCES

1. For an elementary review see M. B. Wise, Proceedings of TASI, Santa Fe. World Ed. R. Slansky.
2. L. J. Hall, Proceedings of the Winter School in Theoretical Physics, Mahabaleshwar, India Lecture Notes In Physics, Springer Verlag Ed. P. Roy and V. Singh 208 (1984) 197.
3. I. Hinchliffe, Proceeding of the TASI in Elementary Particle Physics, Santa Cruz, CA, June 1986, World Scientific (1987). Ed. H. Haber p 409.
4. J. Ellis, G. Gelmini, C. Jarlskog, G. G. Ross and J. W. F. Valle, *Phys. Lett.* 150B 142 (1985).
5. M.C. Bento, L. J. Hall and G. G. Ross, *Nucl. Phys.* **B292** 400 (1987).
6. Brahm and L. J. Hall, LBL preprint, LBL 27165 (1989).
7. L.J. Hall and M. Suzuki, *Nucl. Phys.* **B231** 491 (1984).
8. This is based on: S. Dimopoulos and L. J. Hall, *Phys. Lett.* **B207**, 210 (1987). There are many other signatures, such as rare decays of mesons, which have been considered in the literature^{7,9} which I will not cover.
9. C. Aulakh and R. Mohapatra, *Phys. Lett.* **B119**, 136 (1983); I.H. Lee, *Nucl. Phys.* **B246**, 120 (1984); G.G. Ross and J.W.F. Valle, *Phys. Lett.* **B151**, 375 (1985); J.Ellis, G.Gelmini, C. Jarlskog, G.G. Ross and J.W.F. Valle, *Phys. Lett.* **B150** 142, (1985). S. Dawson, *Nucl. Phys.* **B261**, 297 (1985). F. Zwirner, *Phys. Lett.* **B132** 103, (1983); R. Barbieri and A. Masiero, *Nucl. Phys.* **B267**, 679 (1986).

10. V. Barger, G.F. Guidice and T. Han, University of Wisconsin preprint, MAD/PH/ 465 (1989).
11. S. Dimopoulos, R. Esmailzadeh, L.J. Hall, J. Merlo and G. Starkman, SLAC Pub. 4797 (1988).
12. E. Eichten, I. Hinchliffe, K. Lane and C. Quigg, *Rev. Mod. Phys.* **56** 579 (1984).
13. L.J. Hall, V.A. Kostelecky and S. Raby, *Nucl. Phys.* **B267**, 415 (1986).
14. F. Gabbiani and A. Masiero, *Phys. Lett.* **209**, 289 (1988).
15. S. Dawson, E. Eichten and C. Quigg, *Phys. Rev.* **D31**, 1581 (1985).

Figure Captions

1. Mnemonics for writing the vertices of the minimal low energy supersymmetric model.
2. Total cross-section for $\bar{p}p \rightarrow \tilde{\nu} + \text{anything}$, for $\lambda'_{111} = 1$. The three curves are for $m_{\tilde{\nu}} = 50, 100$ and 250 GeV.
3. The parton diagram for $\bar{p}p \rightarrow \tilde{\nu} \rightarrow \text{jet jet}$.
4. The solid curve is the two jet invariant mass spectrum for $\bar{p}p$ collisions at $\sqrt{s} = 2$ TeV. Both jets must satisfy the rapidity cut $|y| < 0.85$. The dashed line shows the peak of the signal for sneutrino decay into two jets. Hence for this curve $M = m_{\tilde{\nu}}$.
5. The total cross-section for $\bar{p}p \rightarrow \tilde{g}\tilde{g} + \text{anything}$. The masses for the gluino and squarks are 3,20; 50, 50 and 100, 100 GeV. The plot is reproduced from reference.¹⁵
6. One loop diagram for the anomalous dimension of the left-handed down squark scalar mass: $m_{\tilde{d}}^2$.
7. One loop diagram for the anomalous dimension of squark or slepton mass matrices from diagrams involving exotic fields X .

Table 2 [Signatures in the " $\Delta L \neq 0$ " Model]

		LLE^c	QD^cL
<i>LSP</i> Decay* Modes		$\tilde{\nu} \rightarrow \ell^+ \ell^-, \tilde{\chi} \nu$	$\tilde{\nu} \rightarrow \bar{q} q, \tilde{\chi} \nu$
		$\tilde{\chi} \rightarrow \ell^+ \ell^-, \tilde{\nu} \nu$	$\tilde{\chi} \rightarrow \bar{q} q \ell^\pm, \bar{q} q \nu, \tilde{\nu} \nu$
e^+e^- Colliders	Single Superpartner Production	$e^+e^- \rightarrow \tilde{\nu}$ $\tilde{\chi} \nu$	—
	Continuum Pair Production	$e^+e^- \rightarrow \tilde{\nu}$ $\tilde{\chi} \tilde{\chi}$	—
$\bar{p}p$ Colliders	<i>Z</i> Decays	$Z \rightarrow \tilde{\nu}^* \tilde{\nu}$ $\tilde{\chi} \tilde{\chi}$ $\tilde{\nu} \ell^+ \ell^- \dots$	
	<i>W</i> Decays	$W^\pm \rightarrow \tilde{\ell}^\pm \tilde{\nu}$ $\tilde{\chi}^\pm \tilde{\chi} \dots$	
	Continuum pair Production	$\bar{p}p \rightarrow \tilde{\nu}^* \tilde{\nu}, \ell^+ \ell^-, \tilde{\nu} \ell^\pm$ $\bar{p}p \rightarrow \tilde{\nu}, \tilde{\ell}^+ \tilde{\ell}^-, \tilde{\nu} \tilde{\ell}^\pm$	
	Single Superpartner Production	$\bar{p}p \rightarrow \tilde{\nu}, \tilde{\ell}^\pm$ $\tilde{\chi}, \tilde{\chi}^\pm$	

*If LSP is $\tilde{\chi}^\pm$ or $\tilde{\ell}^\pm$ then similar decays occur, but occasionally $\nu \leftrightarrow \ell^\pm$ as required by charge conservation.

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