

Report of the Supersymmetry Theory Subgroup

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ABSTRACT

We provide a mini-guide to some of the possible manifestations of weak-scale supersymmetry. For each of six scenarios we provide

- a brief description of the theoretical underpinnings,
- the adjustable parameters,
- a qualitative description of the associated phenomenology at future colliders,
- comments on how to simulate each scenario with existing event generators.

I. Introduction

The Standard Model (SM) is a theory of spin- $\frac{1}{2}$ matter fermions which interact via the exchange of spin-1 gauge bosons, where the bosons and fermions live in independent representations of the gauge symmetries. Supersymmetry (SUSY) is a symmetry which establishes a one-to-one correspondence between bosonic and fermionic degrees of freedom, and provides a relation between their couplings [1]. Relativistic quantum field theory is formulated to be consistent with the symmetries of the Lorentz/Poincaré group – a non-compact Lie algebra. Mathematically, supersymmetry is formulated as a generalization of the Lorentz/Poincaré group of space-time symmetries to include spinorial generators which obey specific anti-commutation relations; such an algebra is known as a graded Lie algebra. Representations of the SUSY algebra include both bosonic and fermionic degrees of freedom.

The hypothesis that nature is supersymmetric is very compelling to many particle physicists for several reasons.

- It can be shown that the SUSY algebra is the only non-trivial extension of the set of spacetime symmetries which forms one of the foundations of relativistic quantum field theory.
- If supersymmetry is formulated as a *local* symmetry, then one is necessarily forced into introducing a massless spin-2 (graviton) field into the theory. The resulting supergravity theory reduces to Einstein's general relativity theory in the appropriate limit.

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- Spacetime supersymmetry appears to be a fundamental ingredient of superstring theory.

These motivations say nothing about the *scale* at which nature might be supersymmetric. Indeed, there are additional motivations for *weak-scale supersymmetry*.

- Incorporation of supersymmetry into the SM leads to a solution of the gauge hierarchy problem. Namely, quadratic divergences in loop corrections to the Higgs boson mass will cancel between fermionic and bosonic loops. This mechanism works only if the superpartner particle masses are roughly of order or less than the weak scale.
- There exists an experimental hint: the three gauge couplings can unify at the Grand Unification scale if there exist weak-scale supersymmetric particles, with a desert between the weak scale and the GUT scale. This is not the case with the SM.
- Electroweak symmetry breaking is a derived consequence of supersymmetry breaking in many particle physics models with weak-scale supersymmetry, whereas electroweak symmetry breaking in the SM is put in “by hand.” The SUSY radiative electroweak symmetry-breaking mechanism works best if the top quark has mass $m_t \sim 150 - 200$ GeV. The recent discovery of the top quark with $m_t = 176 \pm 4.4$ GeV is consistent with this mechanism.
- As a bonus, many particle physics models with weak-scale supersymmetry contain an excellent candidate for cold dark matter (CDM): the lightest neutralino. Such a CDM particle seems necessary to describe many aspects of cosmology.

Finally, there is a historical precedent for supersymmetry. In 1928, P. A. M. Dirac incorporated the symmetries of the Lorentz group into quantum mechanics. He found as a natural consequence that each known particle had to have a partner particle – namely, antimatter. The matter-anti-matter symmetry wasn't revealed until high enough energy scales were reached to create a positron. In a similar manner, incorporation of supersymmetry into particle physics once again predicts partner particles for all known particles. Will nature prove to be supersymmetric at the weak scale? In this report, we try to shed light on some of the many possible ways that weak-scale supersymmetry might be revealed by colliders operating at sufficiently high energy.

	Boson fields	Fermionic partners
Gauge multiplets		
$SU(3)$	g^a	\tilde{g}^a
$SU(2)$	W^i	\tilde{W}^i
$U(1)$	B	\tilde{B}
Matter multiplets		
Scalar leptons	$\tilde{L}^j = (\tilde{\nu}, \tilde{e}_L^-)$ $\tilde{R} = \tilde{e}_R^+$	$(\nu, e^-)_L$ e_L^c
Scalar quarks	$\tilde{Q}^j = (\tilde{u}_L, \tilde{d}_L)$ $\tilde{U} = \tilde{u}_R^*$ $\tilde{D} = \tilde{d}_R^*$	$(u, d)_L$ u_L^c d_L^c
Higgs bosons	H_1^j H_2^j	$(\tilde{H}_1^0, \tilde{H}_1^-)_L$ $(\tilde{H}_2^+, \tilde{H}_2^0)_L$

Table I: Field content of the MSSM for one generation of quarks and leptons.

A. Minimal Supersymmetric Standard Model

The simplest supersymmetric model of particle physics which is consistent with the SM is called the Minimal Supersymmetric Standard Model (MSSM). The recipe for this model is to start with the SM of particle physics, but in addition add an extra Higgs doublet of opposite hypercharge. (This ensures cancellation of triangle anomalies due to Higgsino partner contributions.) Next, proceed with supersymmetrization, following well-known rules to construct supersymmetric gauge theories. At this stage one has a globally supersymmetric SM theory. Supersymmetry breaking is incorporated by adding to the Lagrangian explicit soft SUSY-breaking terms consistent with the symmetries of the SM. These consist of scalar and gaugino mass terms, as well as trilinear (A terms) and bilinear (B term) scalar interactions. The resulting theory has > 100 parameters, mainly from the various soft SUSY-breaking terms. Such a model is the most conservative approach to realistic SUSY model building, but the large parameter space leaves little predictivity. What is needed as well is a theory of how the soft SUSY-breaking terms arise. The fundamental field content of the MSSM is listed in Table 1, for one generation of quark and lepton (squark and slepton) fields. Mixings and symmetry breaking lead to the actual physical mass eigenstates.

The goal of this report is to create a mini-guide to some of the possible supersymmetric models that occur in the literature, and to provide a bridge between SUSY model builders and their experimental colleagues. The following sections each contain a brief survey of six classes of SUSY-breaking models studied at this workshop; contributing group members are listed in *italics*. We start with the most popular framework for experimental searches, the paradigm

- minimal supergravity model (mSUGRA) (*M. Drees and M. Nojiri*),

and follow with

- models with additional D-term contributions to scalar masses, (*C. Kolda, S. Martin and S. Mrenna*)

- models with non-universal GUT-scale soft SUSY-breaking terms, (*G. Anderson, R. M. Barnett, C. H. Chen, J. Gunion, J. Lykken, T. Moroi and Y. Yamada*)
- two MSSM scenarios which use the large parameter freedom of the MSSM to fit to various collider zoo events, (*G. Kane and S. Mrenna*)
- models with R parity violation, (*H. Baer, B. Kayser and X. Tata*) and
- models with gauge-mediated low energy SUSY breaking (GMLESB), (*J. Amundson, C. Kolda, S. Martin, T. Moroi, S. Mrenna, D. Pierce, S. Thomas, J. Wells and B. Wright*).

Each section contains a brief description of the model, qualitative discussion of some of the associated phenomenology, and finally some comments on event generation for the model under discussion. In this way, it is hoped that this report will be a starting point for future experimental SUSY searches, and that it will provide a flavor for the diversity of ways that weak-scale supersymmetry might manifest itself at colliding beam experiments. We note that a survey of some additional models is contained in Ref. [2], although under a somewhat different format.

II. Minimal Supergravity Model

The currently most popular SUSY model is the minimal supergravity (mSUGRA) model [3, 4]. Here one assumes that SUSY is broken spontaneously in a “hidden sector,” so that some auxiliary field(s) get vev(s) of order $M_Z \cdot M_{Pl} \simeq (10^{10} \text{ GeV})^2$. Gravitational – strength interactions then *automatically* transmit SUSY breaking to the “visible sector,” which contains all the SM fields and their superpartners; the effective mass splitting in the visible sector is by construction of order of the weak scale, as needed to stabilize the gauge hierarchy. In *minimal* supergravity one further assumes that the kinetic terms for the gauge and matter fields take the canonical form: as a result, all scalar fields (sfermions and Higgs bosons) get the same contribution m_0^2 to their squared scalar masses, and that all trilinear A parameters have the same value A_0 , by virtue of an approximate global $U(n)$ symmetry of the SUGRA Lagrangian [4]. Finally, motivated by the apparent unification of the measured gauge couplings within the MSSM [5] at scale $M_{\text{GUT}} \simeq 2 \cdot 10^{16} \text{ GeV}$, one assumes that SUSY-breaking gaugino masses have a common value $m_{1/2}$ at scale M_{GUT} . In practice, since little is known about physics between the scales M_{GUT} and M_{Planck} , one often uses M_{GUT} as the scale at which the scalar masses and A parameters unify. We note that R parity is assumed to be conserved within the mSUGRA framework.

This ansatz has several advantages. First, it is very economical; the entire spectrum can be described with a small number of free parameters. Second, degeneracy of scalar masses at scale M_{GUT} leads to small flavor-changing neutral currents. Finally, this model predicts radiative breaking of the electroweak gauge symmetry [6] because of the large top-quark mass.

Radiative symmetry breaking together with the precisely known value of M_Z allows one to trade two free parameters, usually taken to be the absolute value of the supersymmetric

Higgsino mass parameter $|\mu|$ and the B parameter appearing in the scalar Higgs potential, for the ratio of vevs, $\tan\beta$. The model then has four continuous and one discrete free parameter not present in the SM:

$$m_0, m_{1/2}, A_0, \tan\beta, \text{sign}(\mu). \quad (1)$$

This model is now incorporated in several publicly available MC codes, in particular ISAJET [7]. An approximate version is incorporated into Spythia [8], which reproduces ISAJET results to 10%. Most SUSY spectra studied at this workshop have been generated within mSUGRA; we refer to the various accelerator subgroup reports for the corresponding spectra. One “generically” finds the following features:

- $|\mu|$ is large, well above the masses of the $SU(2)$ and $U(1)$ gauginos. The lightest neutralino is therefore mostly a Bino (and an excellent candidate for cosmological CDM – for related constraints, see *e.g.* Ref. [9]), and the second neutralino and lighter chargino are dominantly $SU(2)$ gauginos. The heavier neutralinos and charginos are only rarely produced in the decays of gluinos and sfermions (except possibly for stop decays). Small regions of parameter space with $|\mu| \simeq M_W$ are possible.
- If $m_0^2 \gg m_{1/2}^2$, all sfermions of the first two generations are close in mass. Otherwise, squarks are significantly heavier than sleptons, and $SU(2)$ doublet sleptons are heavier than singlet sleptons. Either way, the lighter stop and sbottom eigenstates are well below the first generation squarks; gluinos therefore have large branching ratios into b or t quarks.
- The heavier Higgs bosons (pseudoscalar A , scalar H^0 , and charged H^\pm) are usually heavier than $|\mu|$ unless $\tan\beta \gg 1$. This also implies that the light scalar h^0 behaves like the SM Higgs.

These features have already become something like folklore. We want to emphasize here that even within this restrictive framework, quite different spectra are also possible, as illustrated by the following examples.

Example A is for $m_0 = 750$ GeV, $m_{1/2} = 150$ GeV, $A_0 = -300$ GeV, $\tan\beta = 5.5$, $\mu < 0$, and $m_t = 165$ GeV (pole mass). This yields $|\mu| = 120$ GeV, very similar to the $SU(2)$ gaugino mass M_2 at the weak scale, leading to strong Higgsino – gaugino mixing. The neutralino masses are 60, 91, 143 and 180 GeV, while charginos are at 93 and 185 GeV. They are all considerably lighter than the gluino (at 435 GeV), which in turn lies well below the squarks (at $\simeq 815$ GeV) and sleptons (at 750-760 GeV). Due to the strong gaugino – Higgsino mixing, all chargino and neutralino states will be produced with significant rates in the decays of gluinos and $SU(2)$ doublet sfermions, leading to complicated decay chains. For example, the $\ell^+\ell^-$ invariant mass spectrum in gluino pair events will have many thresholds due to $\tilde{\chi}_i^0 \rightarrow \tilde{\chi}_j^0 \ell^+ \ell^-$ decays. Since first and second generation squarks are almost twice as heavy as the gluino, there might be a significant gluino “background” to squark production at the LHC. A 500 GeV e^+e^- collider

will produce all six chargino and neutralino states. Information about \tilde{e}_L , \tilde{e}_R and $\tilde{\nu}_e$ masses can be gleaned from studies of neutralino and chargino production, respectively; however, $\sqrt{s} > 1.5$ TeV is required to study sleptons directly. Spectra of this type can already be modelled reliably using ISAJET: the above parameter space set can be entered via the *SUGRA* keyword.

As example B, we have chosen $m_0 = m_{1/2} = 200$ GeV, $A_0 = 0$, $\tan\beta = 48$, $\mu < 0$ and $m_t = 175$ GeV. Note the large value of $\tan\beta$, which leads to large b and τ Yukawa couplings, as required in models where all third generation Yukawa couplings are unified at scale M_{GUT} . Here the gluino (at 517 GeV) lies slightly above first generation squarks (at 480-500 GeV), which in turn lie well above first generation sleptons (at 220-250 GeV). The light neutralinos (at 83 and 151 GeV) and light chargino (at 151 GeV) are mostly gauginos, while the heavy states (at 287, 304 and 307 GeV) are mostly Higgsinos, because $|\mu| = 275$ GeV $\gg m_{1/2}$.

The masses of \tilde{t}_1 (355 GeV), \tilde{b}_1 (371 GeV) and $\tilde{\tau}_1$ (132 GeV) are all significantly below those of the corresponding first or second generation sfermions. As a result, more than 2/3 of all gluinos decay into a b quark and a \tilde{b} squark. Since (s)bottoms have large Yukawa couplings, \tilde{b} decays will often produce the heavier, Higgsino-like chargino and neutralinos. Further, all neutralinos (except for the lightest one, which is the LSP) have two-body decays into $\tilde{\tau}_1 + \tau$; in case of $\tilde{\chi}_2^0$ this is the only two-body mode, and for the Higgsino-like states this mode will be enhanced by the large τ Yukawa coupling. Chargino decays will also often produce real $\tilde{\tau}_1$. Study of the $\ell^+\ell^-$ invariant mass spectrum will not allow direct determination of neutralino mass differences, as the ℓ^\pm are secondaries from tau decays. Even \tilde{e}_L pair events at e^+e^- colliders will contain up to four tau leptons! Further, unless the e^- beam is almost purely right-handed, it might be difficult to distinguish between $\tilde{\tau}_1$ pair production and $\tilde{\chi}_1^\pm$ pair production. Finally, the heavier Higgs bosons are quite light in this case, *e.g.* $m_A = 126$ GeV. There will be a large number of $A \rightarrow \tau^+\tau^-$ events at the LHC. However, because most SUSY events will contain τ pairs in this scenario, it is not clear whether the Higgs signal will remain visible. At present, scenarios with $\tan\beta \gg 1$ can not be simulated with ISAJET, since the b and τ Yukawa couplings have not been included in all relevant decays. This situation should be remedied soon.

III. D -term Contributions to Scalar Masses

We have seen that the standard mSUGRA framework predicts a testable pattern of squark and slepton masses. In this section we describe a class of models in which a quite distinctive modification of the mSUGRA predictions can arise, namely contributions to scalar masses associated with the D -terms of extra spontaneously broken gauge symmetries [10]. As we will see, the modification of squark, slepton and Higgs masses can have a profound effect on phenomenology.

In general, D -term contributions to scalar masses will arise in supersymmetric models whenever a gauge symmetry is spontaneously broken with a reduction of rank. Suppose, for example, that the SM gauge group $SU(3) \times SU(2) \times U(1)_Y$ is

supplemented by an additional $U(1)_X$ factor broken far above the electroweak scale. Naively, one might suppose that if the breaking scale is sufficiently large, all direct effects of $U(1)_X$ on TeV-scale physics are negligible. However, a simple toy model shows that this is not so. Assume that ordinary MSSM scalar fields, denoted generically by φ_i , carry $U(1)_X$ charges X_i which are not all 0. In order to break $U(1)_X$, we also assume the existence of a pair of additional chiral superfields Φ and $\bar{\Phi}$ which are SM singlets, but carry $U(1)_X$ charges which are normalized (without loss of generality) to be +1 and -1 respectively. Then VEV's for Φ and $\bar{\Phi}$ will spontaneously break $U(1)_X$ while leaving the SM gauge group intact. The scalar potential whose minimum determines $\langle\Phi\rangle, \langle\bar{\Phi}\rangle$ then has the form

$$V = V_0 + m^2|\Phi|^2 + \bar{m}^2|\bar{\Phi}|^2 + \frac{g_X^2}{2} [|\Phi|^2 - |\bar{\Phi}|^2 + X_i|\varphi_i|^2]^2. \quad (2)$$

Here V_0 comes from the superpotential and involves only Φ and $\bar{\Phi}$; it is symmetric under $\Phi \leftrightarrow \bar{\Phi}$, but otherwise its precise form need not concern us. The pieces involving m^2 and \bar{m}^2 are soft breaking terms; m^2 and \bar{m}^2 are of order M_Z^2 and in general unequal. The remaining piece is the square of the D -term associated with $U(1)_X$, which forces the minimum of the potential to occur along a nearly D -flat direction $\langle\Phi\rangle \approx \langle\bar{\Phi}\rangle$. This scale can be much larger than 1 TeV with natural choices of V_0 , so that the $U(1)_X$ gauge boson is very heavy and plays no role in collider physics.

However, there is also a deviation from D -flatness given by $\langle\Phi\rangle^2 - \langle\bar{\Phi}\rangle^2 \approx D_X/g_X^2$, with $D_X = (\bar{m}^2 - m^2)/2$, which directly affects the masses of the remaining light MSSM fields. After integrating out Φ and $\bar{\Phi}$, one finds that each MSSM scalar (mass)² receives a correction given by

$$\Delta m_i^2 = X_i D_X \quad (3)$$

where D_X is again typically of order M_Z^2 and may have either sign. This result does not depend on the scale at which $U(1)_X$ breaks; this turns out to be a general feature, independent of assumptions about the precise mechanism of symmetry breaking. Thus $U(1)_X$ manages to leave its “fingerprint” on the masses of the squarks, sleptons, and Higgs bosons, even if it is broken at an arbitrarily high energy. From a TeV-scale point of view, the parameter D_X might as well be taken as a parameter of our ignorance regarding physics at very high energies. The important point is that D_X is universal, so that each MSSM scalar (mass)² obtains a contribution simply proportional to X_i , its charge under $U(1)_X$. Typically the X_i are rational numbers and do not all have the same sign, so that a particular candidate $U(1)_X$ can leave a quite distinctive pattern of mass splittings on the squark and slepton spectrum.

The extra $U(1)_X$ in this discussion may stand alone, or may be embedded in a larger non-abelian gauge group, perhaps together with the SM gauge group (for example in an $SO(10)$ or E_6 GUT). If the gauge group contains more than one $U(1)$ in addition to $U(1)_Y$, then each $U(1)$ factor can contribute a set of corrections exactly analogous to (3). Additional $U(1)$ groups are endemic in superstring models, so at least from that point of view one may be optimistic about the existence of corresponding D -terms and their potential importance in the study

of the squark and slepton mass spectrum at future colliders. It should be noted that once one assumes the existence of additional gauged $U(1)$'s at very high energies, it is quite unnatural to assume that D -term contributions to scalar masses can be avoided altogether. (This would require an exact symmetry enforcing $m^2 = \bar{m}^2$ in the example above.) The only question is whether or not the magnitude of the D -term contributions is significant compared to the usual mSUGRA contributions. Note also that as long as the charges X_i are family-independent, then from (3) squarks and sleptons with the same electroweak quantum numbers remain degenerate, maintaining the natural suppression of flavor changing neutral currents.

It is not difficult to implement the effects of D -terms in simulations, by imposing the corrections (3) to a particular “template” mSUGRA model. After choosing the $U(1)_X$ charges of the MSSM fields, our remaining ignorance of the mechanism of $U(1)_X$ breaking is parameterized by D_X (roughly of order M_Z^2). The Δm_i^2 corrections should be imposed at the scale M_X where one chooses to assume that $U(1)_X$ breaks. (If $M_X < M_{\text{Planck}}$ or M_{GUT} , one should also in principle incorporate renormalization group effects due to $U(1)_X$ above M_X , but these can often be shown to be small.) The other parameters of the theory are unaffected. One can then run these parameters down to the electroweak scale, in exactly the same way as in mSUGRA models, to find the spectrum of sparticle masses.

(The solved-for parameter μ is then indirectly affected by D -terms, through the requirement of correct electroweak symmetry breaking.) The only subtlety involved is an apparent ambiguity in choosing the charges X_i , since any linear combination of $U(1)_X$ and $U(1)_Y$ charges might be used. These charges should be picked to correspond to the basis in which there is no mixing in the kinetic terms of the $U(1)$ gauge bosons. In particular models where $U(1)_X$ and/or $U(1)_Y$ are embedded in non-abelian groups, this linear combination is uniquely determined; otherwise it can be arbitrary.

A test case which seems particularly worthy of study is that of an additional gauged $B - L$ symmetry. In this case the $U(1)_X$ charges for each MSSM scalar field are a linear combination of $B - L$ and Y . If this model is embedded in $SO(10)$ (or certain of its subgroups), then the unmixed linear combination of $U(1)$'s appropriate for (3) is $X = -\frac{5}{3}(B - L) + \frac{4}{3}Y$. The X charges for the MSSM squarks and sleptons are $-1/3$ for Q_L, u_R, e_R and $+1$ for L_L and d_R . The MSSM Higgs fields have charges $+2/3$ for H_u and $-2/3$ for H_d . Here we consider the modifications to a mSUGRA model defined by the parameters $(m_0, m_{1/2}, A_0) = (200, 100, 0)$ GeV, $\mu < 0$, and $\tan\beta = 2$, assuming $m_t = 175$ GeV.

The effects of D -term contributions to the scalar mass spectrum is illustrated in Fig. 1, which shows the masses of \tilde{e}_L, \tilde{e}_R , the lightest Higgs boson h , and the lightest bottom squark \tilde{b}_1 as a function of D_X . The unmodified mSUGRA prediction is found at $D_X = 0$. A particularly dramatic possibility is that D -terms could invert the usual hierarchy of slepton masses, so that $m_{\tilde{e}_L}, m_{\tilde{\nu}} < m_{\tilde{e}_R}$. In the test model, this occurs for negative D_X ; the negative endpoint of D_X is set by the experimental lower bound on $m_{\tilde{\nu}}$. The relative change of the squark masses is smaller, while the change to the lightest Higgs boson mass is

almost negligible except near the positive D_X endpoint where it reaches the experimental lower bound. The complicated mass spectrum perhaps can be probed most directly at the NLC with precision measurements of squark and slepton masses. Since the usual MSSM renormalization group contributions to scalar masses are much larger for squarks than for sleptons, it is likely that the effects of D -term contributions are relatively larger for sleptons.

At the Tevatron and LHC, it has been suggested in these proceedings that SUSY parameter determinations can be obtained by making global fits of the mSUGRA parameter space to various observed signals. In this regard it should be noted that significant D -term contributions could invalidate such strategies unless they are generalized. This is because adding D -terms (3) to a given template mSUGRA model can dramatically change certain branching fractions by altering the kinematics of decays involving squarks and especially sleptons. This is demonstrated for the test model in Fig. 2. Thus we find for example that the product $BR(\tilde{\chi}_1^+ \rightarrow \ell^+ X) \times BR(\tilde{\chi}_2^0 \rightarrow \ell^+ \ell^- X)$ can change up to an order of magnitude or more as one varies D -terms (with all other parameters held fixed). Note that the branching ratios of Fig. 2 include the leptons from two-body and three-body decays, e.g. $\tilde{\chi}_1^+ \rightarrow \ell^+ \nu \tilde{\chi}_1^0$ and $\tilde{\chi}_1^+ \rightarrow \tilde{\ell}^+ \nu \rightarrow \ell^+ \tilde{\chi}_j^0 \nu$. On the other hand, the $BR(\tilde{g} \rightarrow bX)$ is fairly insensitive to D -terms over most, but not all, of parameter space.

Since the squark masses are generally much less affected by the D -terms, and the gluino mass only indirectly, the production cross sections for squarks and gluinos should be fairly stable. Therefore, the variation of $BR(\tilde{g} \rightarrow bX)$ is an accurate gauge of the variation of observables such as the b multiplicity of SUSY events. Likewise, the $\tilde{\chi}_1^+ \tilde{\chi}_2^0$ production cross section does not change much as the D -terms are varied, so the expected trilepton signal can vary like the product of branching ratios – by orders of magnitude. While the results presented are for a specific, and particularly simple, test model, similar variations can be observed in other explicit models. The possible presence of D -terms should be considered when interpreting a SUSY signal at future colliders. An experimental analysis which proves or disproves their existence would be a unique insight into physics at very high energy scales.

To facilitate event generation, approximate expressions for the modified mass spectra are implemented in the *Spythia* Monte Carlo, assuming the D -terms are added in at the unification scale. Sparticle spectra from models with extra D -terms can be incorporated into *ISAJET* simply via the *MSSMi* keywords, although the user must supply a program to generate the relevant spectra via RGE's or analytic formulae.

IV. Non-Universal GUT-Scale Soft SUSY-Breaking Parameters

A. Introduction

We considered models in which the gaugino masses and/or the scalar masses are not universal at the GUT scale, M_{GUT} . We study the extent to which non-universal boundary conditions can influence experimental signatures and detector re-

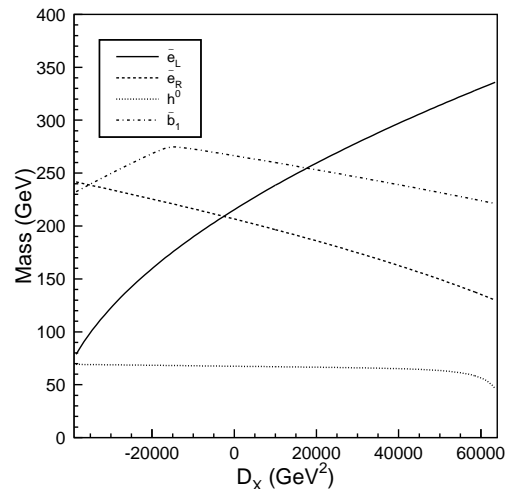


Figure 1: Mass spectrum as a function of D_X .

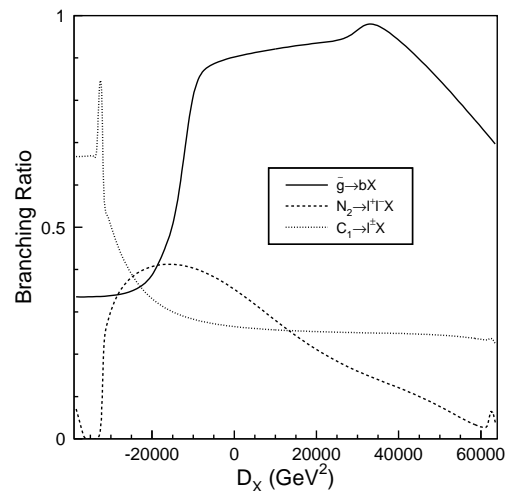


Figure 2: Branching ratios as a function of D_X .

quirements, and the degree to which experimental data can distinguish between different models for the GUT-scale boundary conditions.

1. Non-Universal Gaugino Masses at M_{GUT}

We focus on two well-motivated types of models:

- Superstring-motivated models in which SUSY breaking is moduli dominated. We consider the particularly attractive O-II model of Ref. [11]. The boundary conditions at M_{GUT} are:

$$\begin{aligned} M_a^0 &\sim \sqrt{3}m_{3/2}[-(b_a + \delta_{GS})K\eta] \\ m_0^2 &= m_{3/2}^2[-\delta_{GS}K'] \\ A_0 &= 0 \end{aligned} \quad (4)$$

where b_a are SM beta function coefficients, δ_{GS} is a mixing parameter, which would be a negative integer in the O-II model,

F_Φ	M_{GUT}			m_Z		
	M_3	M_2	M_1	M_3	M_2	M_1
1	1	1	1	~ 6	~ 2	~ 1
24	2	-3	-1	~ 12	~ -6	~ -1
75	1	3	-5	~ 6	~ 6	~ -5
200	1	2	10	~ 6	~ 4	~ 10
$O-II$ $\delta_{GS} = -4$	1	5	$\frac{53}{5}$	~ 6	~ 10	$\sim \frac{53}{5}$

Table II: Relative gaugino masses at M_{GUT} and m_Z in the four possible F_Φ irreducible representations, and in the O-II model with $\delta_{GS} \sim -4$.

and $\eta = \pm 1$. From the estimates of Ref. [11], $K \simeq 4.6 \times 10^{-4}$ and $K' \simeq 10^{-3}$, we expect that slepton and squark masses would be very much larger than gaugino masses.

• Models in which SUSY breaking occurs via an F -term that is not an $SU(5)$ singlet. In this class of models, gaugino masses are generated by a chiral superfield Φ that appears linearly in the gauge kinetic function, and whose auxiliary F component acquires an intermediate-scale vev:

$$\mathcal{L} \sim \int d^2\theta W^a W^b \frac{\Phi_{ab}}{M_{\text{Planck}}} + h.c. \sim \frac{\langle F_\Phi \rangle_{ab}}{M_{\text{Planck}}} \lambda^a \lambda^b + \dots, \quad (5)$$

where the $\lambda^{a,b}$ are the gaugino fields. F_Φ belongs to an $SU(5)$ irreducible representation which appears in the symmetric product of two adjoints:

$$(\mathbf{24} \times \mathbf{24})_{\text{symmetric}} = \mathbf{1} \oplus \mathbf{24} \oplus \mathbf{75} \oplus \mathbf{200}, \quad (6)$$

where only **1** yields universal masses. Only the component of F_Φ that is “neutral” with respect to the SM gauge group should acquire a vev, $\langle F_\Phi \rangle_{ab} = c_a \delta_{ab}$, with c_a then determining the relative magnitude of the gauginos masses at M_{GUT} : see Table II.

Physical masses of the gauginos are influenced by $\tan\beta$ -dependent off-diagonal terms in the mass matrices and by corrections which boost $m_{\tilde{g}}(\text{pole})$ relative to $m_{\tilde{g}}(m_{\tilde{g}})$. If μ is large, the lightest neutralino (which is the LSP) will have mass $m_{\tilde{\chi}_1^0} \sim \min(M_1, M_2)$ while the lightest chargino will have $m_{\tilde{\chi}_1^\pm} \sim M_2$. Thus, in the **200** and O-II scenarios with $M_2 \lesssim M_1$, $m_{\tilde{\chi}_1^\pm} \simeq m_{\tilde{\chi}_1^0}$ and the $\tilde{\chi}_1^\pm$ and $\tilde{\chi}_1^0$ are both Wino-like. The $\tan\beta$ dependence of the masses at m_Z for the universal, **24**, **75**, and **200** choices appears in Fig. 3. The $m_{\tilde{g}}-m_{\tilde{\chi}_1^0}$ mass splitting becomes increasingly smaller in the sequence **24**, **1**, **200** **75**, O-II, as could be anticipated from Table II. It is interesting to note that at high $\tan\beta$, μ decreases to a level comparable to M_1 and M_2 , and there is substantial degeneracy among the $\tilde{\chi}_1^\pm$, $\tilde{\chi}_2^0$ and $\tilde{\chi}_1^0$.

1. Non-Universal Scalar Masses at M_{GUT}

We consider models in which the SUSY-breaking scalar masses at M_{GUT} are influenced by the Yukawa couplings of the corresponding quarks/leptons. This idea is exemplified in the model of Ref. [12] based on perturbing about the $[U(3)]^5$

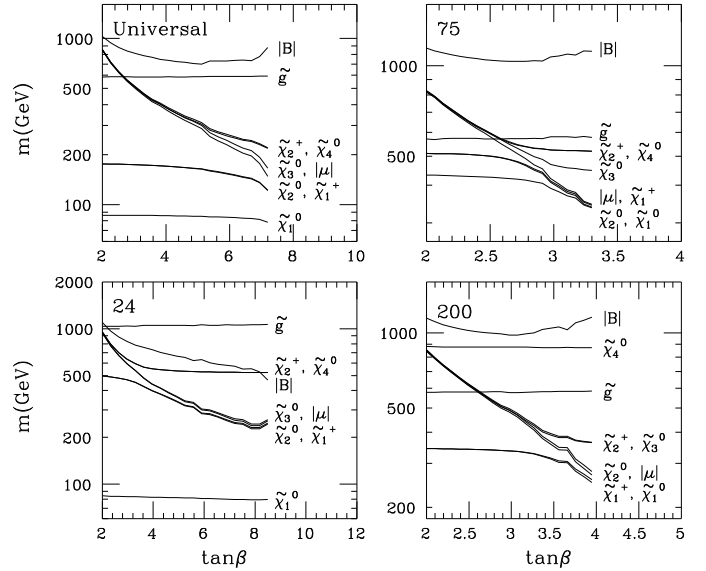


Figure 3: Physical (pole) gaugino masses as a function of $\tan\beta$ for the **1** (universal), **24**, **75**, and **200** F representation choices. Also plotted are $|B|$ and $|\mu|$. We have taken $m_0 = 1$ TeV and $M_3 = 200, 400, 200, 200$ GeV, respectively.

symmetry that is present in the absence of Yukawa couplings. One finds, for example:

$$m_Q^2 = m_0^2 (I + c_Q \lambda_u^\dagger \lambda_u + c'_Q \lambda_d^\dagger \lambda_d + \dots) \quad (7)$$

where Q represents the squark partners of the left-handed quark doublets. The Yukawas λ_u and λ_d are 3×3 matrices in generation space. The \dots represent terms of order λ^4 that we will neglect. A priori, c_Q , c'_Q , should all be similar in size, in which case the large top-quark Yukawa coupling implies that the primary deviations from universality will occur in $m_{\tilde{t}_L}^2$, $m_{\tilde{b}_L}^2$ (equally and in the same direction).¹ It is the fact that $m_{\tilde{t}_L}^2$ and $m_{\tilde{b}_L}^2$ are shifted equally that will distinguish m^2 non-universality from the effects of a large A_0 parameter at M_{GUT} ; the latter would primarily introduce $\tilde{t}_L - \tilde{t}_R$ mixing and yield a low $m_{\tilde{t}_1}$ compared to $m_{\tilde{b}_1}$.

B. Phenomenology

2. Non-Universal Gaugino Masses

We examined the phenomenological implications for the standard Snowmass comparison point (e.g. NLC point #3) specified by $m_t = 175$ GeV, $\alpha_s = 0.12$, $m_0 = 200$ GeV, $M_3^0 = 100$ GeV, $\tan\beta = 2$, $A_0 = 0$ and $\mu < 0$. In treating the O-II model we take $m_0 = 600$ GeV, a value that yields a (pole) value of $m_{\tilde{g}}$ not unlike that for the other scenarios. The masses of the supersymmetric particles for each scenario are given in Table III.

The phenomenology of these scenarios for e^+e^- collisions is not absolutely straightforward.

¹In this discussion we neglect an analogous, but independent, shift in $m_{\tilde{t}_R}^2$.

	1	24	75	200	$O-II$ $\delta_{GS} = -4.7$
$m_{\tilde{g}}$	285	285	287	288	313
$m_{\tilde{u}_R}$	302	301	326	394	-
$m_{\tilde{t}_1}$	255	257	235	292	-
$m_{\tilde{t}_2}$	315	321	351	325	-
$m_{\tilde{b}_L}$	266	276	307	264	-
$m_{\tilde{b}_R}$	303	303	309	328	-
$m_{\tilde{\tau}_R}$	207	204	280	437	-
$m_{\tilde{\tau}_L}$	216	229	305	313	-
$m_{\tilde{\chi}_1^0}$	44.5	12.2	189	174.17	303.09
$m_{\tilde{\chi}_2^0}$	97.0	93.6	235	298	337
$m_{\tilde{\chi}_1^\pm}$	96.4	90.0	240	174.57	303.33
$m_{\tilde{\chi}_2^\pm}$	275	283	291	311	-
m_{h^0}	67	67	68	70	82

Table III: Sparticle masses for the Snowmass comparison point in the different gaugino mass scenarios. Blank entries for the O-II model indicate very large masses.

- In the **75** model, $\tilde{\chi}_1^+ \tilde{\chi}_1^-$ and $\tilde{\chi}_2^0 \tilde{\chi}_2^0$ pair production at $\sqrt{s} = 500$ GeV are barely allowed kinematically; the phase space for $\tilde{\chi}_1^0 \tilde{\chi}_2^0$ is only somewhat better. All the signals would be rather weak, but could probably be extracted with sufficient integrated luminosity.
- In the **200** model, $e^+ e^- \rightarrow \tilde{\chi}_1^+ \tilde{\chi}_1^-$ production would be kinematically allowed at a $\sqrt{s} = 500$ GeV NLC, but not easily observed due to the fact that the (invisible) $\tilde{\chi}_1^0$ would take essentially all of the energy in the $\tilde{\chi}_1^\pm$ decays. However, according to the results of Ref. [13], $e^+ e^- \rightarrow \gamma \tilde{\chi}_1^+ \tilde{\chi}_1^-$ would be observable at $\sqrt{s} = 500$ GeV.
- The O-II model with δ_{GS} near -4 predicts that $m_{\tilde{\chi}_1^\pm}$ and $m_{\tilde{\chi}_1^0}$ are both rather close to $m_{\tilde{g}}$, so that $e^+ e^- \rightarrow \tilde{\chi}_1^+ \tilde{\chi}_1^-$, $\tilde{\chi}_1^0 \tilde{\chi}_1^0$ would *not* be kinematically allowed at $\sqrt{s} = 500$ GeV. The only SUSY “signal” would be the presence of a very SM-like light Higgs boson.

At the LHC, the strongest signal for SUSY would arise from $\tilde{g}\tilde{g}$ production. The different models lead to very distinct signatures for such events. To see this, it is sufficient to list the primary easily identifiable decay chains of the gluino for each of the five scenarios. (In what follows, q denotes any quark other than a b .)

$$\begin{aligned}
\mathbf{1} : \quad & \tilde{g} \xrightarrow{90\%} \tilde{b}_L \bar{b} \xrightarrow{99\%} \tilde{\chi}_2^0 b \bar{b} \xrightarrow{33\%} \tilde{\chi}_1^0 (e^+ e^- \text{ or } \mu^+ \mu^-) b \bar{b} \\
& \xrightarrow{8\%} \tilde{\chi}_1^0 \nu \bar{\nu} b \bar{b} \\
& \xrightarrow{38\%} \tilde{\chi}_1^0 q \bar{q} b \bar{b} \\
& \xrightarrow{8\%} \tilde{\chi}_1^0 b \bar{b} b \bar{b} \\
\mathbf{24} : \quad & \tilde{g} \xrightarrow{85\%} \tilde{b}_L \bar{b} \xrightarrow{70\%} \tilde{\chi}_2^0 b \bar{b} \xrightarrow{99\%} h^0 \tilde{\chi}_1^0 b \bar{b} \xrightarrow{28\%} \tilde{\chi}_1^0 b \bar{b} b \bar{b} \\
& \xrightarrow{69\%} \tilde{\chi}_1^0 \tilde{\chi}_1^0 \tilde{\chi}_1^0 b \bar{b} \\
\mathbf{75} : \quad & \tilde{g} \xrightarrow{43\%} \tilde{\chi}_1^0 g \text{ or } \tilde{\chi}_1^0 q \bar{q}
\end{aligned}$$

$$\begin{aligned}
& \xrightarrow{10\%} \tilde{\chi}_1^0 b \bar{b} \\
& \xrightarrow{20\%} \tilde{\chi}_2^0 g \text{ or } \tilde{\chi}_2^0 q \bar{q} \\
& \xrightarrow{10\%} \tilde{\chi}_2^0 b \bar{b} \\
& \xrightarrow{17\%} \tilde{\chi}_1^\pm q \bar{q} \\
\mathbf{200} : \quad & \tilde{g} \xrightarrow{99\%} \tilde{b}_L \bar{b} \xrightarrow{100\%} \tilde{\chi}_1^0 b \bar{b} \\
\mathbf{O-II} : \quad & \tilde{g} \xrightarrow{51\%} \tilde{\chi}_1^\pm q \bar{q} \\
& \xrightarrow{17\%} \tilde{\chi}_1^0 g \\
& \xrightarrow{26\%} \tilde{\chi}_1^0 q \bar{q} \\
& \xrightarrow{6\%} \tilde{\chi}_1^0 b \bar{b}
\end{aligned}$$

Gluino pair production will then lead to the following strikingly different signals.

- In the **1** scenario we expect a very large number of final states with missing energy, four b -jets and two lepton-antilepton pairs.
- For **24**, an even larger number of events will have missing energy and eight b -jets, four of which reconstruct to two pairs with mass equal to (the known) m_{h^0} .
- The signal for $\tilde{g}\tilde{g}$ production in the case of **75** is much more traditional; the primary decays yield multiple jets (some of which are b -jets) plus $\tilde{\chi}_1^0$, $\tilde{\chi}_2^0$ or $\tilde{\chi}_1^\pm$. Additional jets, leptons and/or neutrinos arise when $\tilde{\chi}_2^0 \rightarrow \tilde{\chi}_1^0 +$ two jets, two leptons or two neutrinos or $\tilde{\chi}_1^\pm \rightarrow \tilde{\chi}_1^0 +$ two jets or lepton+neutrino.
- In the **200** scenario, we find missing energy plus four b -jets; only b -jets appear in the primary decay – any other jets present would have to come from initial- or final-state radiation, and would be expected to be softer on average. This is almost as distinctive a signal as the $8b$ final state found in the **24** scenario.
- In the final O-II scenario, $\tilde{\chi}_1^\pm \rightarrow \tilde{\chi}_1^0 +$ very soft spectator jets or leptons that would not be easily detected. Even the $q\bar{q}$ or g from the primary decay would not be very energetic given the small mass splitting between $m_{\tilde{g}}$ and $m_{\tilde{\chi}_1^\pm} \sim m_{\tilde{\chi}_1^0}$. Soft jet cuts would have to be used to dig out this signal, but it should be possible given the very high $\tilde{g}\tilde{g}$ production rate expected for this low $m_{\tilde{g}}$ value; see Ref. [13].

Thus, for the Snowmass comparison point, distinguishing between the different boundary condition scenarios at the LHC will be easy. Further, the event rate for a gluino mass this low is such that the end-points of the various lepton, jet or h^0 spectra will allow relatively good determinations of the mass differences between the sparticles appearing at various points in the final state decay chain. We are optimistic that this will prove to be a general result so long as event rates are large.

2. Non-Universal Scalar Masses

Once again we focus on the Snowmass overlap point. We maintain gaugino mass universality at M_{GUT} , but allow for non-universality for the squark masses. Of the many possibilities, we focus on the case where only $c_Q \neq 0$ with $A_0 = 0$ (as assumed for the Snowmass overlap point). The phenomenology for this case is compared to that which would emerge if we take $A_0 \neq 0$ with all the $c_i = 0$.

Consider the \tilde{g} branching ratios as a function of $m_{\tilde{t}_L} = m_{\tilde{b}_L}$ as c_Q is varied from negative to positive values. As the common mass crosses the threshold above which the $\tilde{g} \rightarrow \tilde{b}_1 b$ decay becomes kinematically disallowed, we revert to a more standard SUSY scenario in which \tilde{g} decays are dominated by modes such as $\tilde{\chi}_1^\pm q\bar{q}$, $\tilde{\chi}_1^0 q\bar{q}$, $\tilde{\chi}_2^0 q\bar{q}$ and $\tilde{\chi}_2^0 b\bar{b}$. For low enough $m_{\tilde{t}_L}$, the $\tilde{g} \rightarrow \tilde{t}_1 t$ mode opens up, but must compete with the $\tilde{g} \rightarrow \tilde{b}_1 b$ mode that has even larger phase space.

In contrast, if A_t is varied, the \tilde{g} branching ratios remain essentially constant until $m_{\tilde{t}_1}$ is small enough that $\tilde{g} \rightarrow \tilde{t}_1 t$ is kinematically allowed. Below this point, this latter mode quickly dominates the $\tilde{b}_1 b$ mode which continues to have very small phase space given that the \tilde{b}_1 mass remains essentially constant as A_t is varied.

C. Event Generation

A thorough search and determination of the rates (or lack thereof) for the full panoply of possible channels is required to distinguish the many possible GUT-scale boundary conditions from one another. In the program `ISAJET`, independent weak-scale gaugino masses may be input using the `MSSM4` keyword. Independent third generation squark masses may be input via the `MSSM2` keyword. The user must supply a program to generate the relevant weak-scale parameter values from the specific GUT-scale assumptions. Relevant weak-scale MSSM parameters can also be input to `Spythia`; as with `ISAJET`, the user must provide a program for the specific model.

V. MSSM Scenarios Motivated by Data

An alternative procedure for gleaning information about the SUSY soft terms is to use the full (~ 100 parameters) parameter space freedom of the MSSM and match to data, assuming one has a supersymmetry signal. This approach has been used in the following two examples.

A. The CDF $e^+e^- \gamma\gamma + \cancel{E}_T$ Event

Recently a candidate for sparticle production has been reported [14] by the CDF collaboration. This has been interpreted in several ways [15], [16], [17], [18] and later with additional variations [19], [20], [21]. The main two paths are whether the LSP is the lightest neutralino [15], [22], or a nearly massless gravitino [16, 17, 18, 19, 20] or axino [21]. In the gravitino or axino case the LSP is not a candidate for cold dark matter, SUSY can have no effect on R_b or α_s^Z or $BR(b \rightarrow s\gamma)$, and stops and gluinos are not being observed at FNAL. In the case where the lightest neutralino is the LSP, the opposite holds for

$e^+e^- \gamma\gamma + \cancel{E}_T$ constraints on supersymmetric parameters	
\tilde{e}_L	\tilde{e}_R
$100 \lesssim m_{\tilde{e}_L} \lesssim 130$ GeV	$100 \lesssim m_{\tilde{e}_R} \lesssim 112$ GeV
$50 \lesssim M_1 \lesssim 92$ GeV	$60 \lesssim M_1 \lesssim 85$ GeV
$50 \lesssim M_2 \lesssim 105$ GeV	$40 \lesssim M_2 \lesssim 85$ GeV
$0.75 \lesssim M_2/M_1 \lesssim 1.6$	$0.6 \lesssim M_2/M_1 \lesssim 1.15$
$-65 \lesssim \mu \lesssim -35$ GeV	$-60 \lesssim \mu \lesssim -35$ GeV
$0.5 \lesssim \mu /M_1 \lesssim 0.95$	$0.5 \lesssim \mu /M_1 \lesssim 0.8$
$1 \lesssim \tan \beta \lesssim 3$	$1 \lesssim \tan \beta \lesssim 2.2$

Table IV: Constraints on the MSSM parameters and masses in the neutralino LSP scenario.

all of these observables, and we will pursue this case in detail here.

The SUSY Lagrangian depends on a number of parameters, all of which have the dimension of mass. That should not be viewed as a weakness because at present we have no theory of the origin of mass parameters. Probably getting such a theory will depend on understanding how SUSY is broken. When there is no data on sparticle masses and couplings, it is appropriate to make simplifying assumptions, based on theoretical prejudice, to reduce the number of parameters. However, once there may be data, it is important to constrain the most general set of parameters and see what patterns emerge. We proceed by making no assumptions about soft breaking parameters. In practice, even though the full theory has over a hundred such parameters, that is seldom a problem since any given observable depends on at most a few.

The CDF event [14] has a 36 GeV e^- , a 59 GeV e^+ , photons of 38 and 30 GeV, and $\cancel{E}_T = 53$ GeV. A SUSY interpretation is $q\bar{q} \rightarrow \gamma^*, Z^* \rightarrow \tilde{e}^+ \tilde{e}^-$, followed by each $\tilde{e}^\pm \rightarrow e^\pm \tilde{\chi}_2^0$, $\tilde{\chi}_2^0 \rightarrow \gamma \tilde{\chi}_1^0$. The second lightest neutralino, $\tilde{\chi}_2^0$, must be photino-like since it couples strongly to $\tilde{e}e$. Then the LSP = $\tilde{\chi}_1^0$ must be Higgsino-like [23, 24, 25] to have a large $BR(\tilde{\chi}_2^0 \rightarrow \tilde{\chi}_1^0 \gamma)$. The range of parameter choices for this scenario are given in Table IV.

If light superpartners indeed exist, FNAL and LEP will produce thousands of them, and measure their properties very well. The first thing to check at FNAL is whether the produced selectron is \tilde{e}_L or \tilde{e}_R . If \tilde{e}_L , then the charged current channel $u\bar{d} \rightarrow W^+ \rightarrow \tilde{e}_L \tilde{\nu}$ has 5-10 times the rate of $\tilde{e}_L^+ \tilde{e}_L^-$. We expect $\tilde{e}_L \rightarrow e \tilde{\chi}_2^0 (\rightarrow \gamma \tilde{\chi}_1^0)$. Most likely [22] $\tilde{\nu} \rightarrow e \tilde{\chi}_1^\pm$, where $\tilde{\chi}_1^\pm$ is the lightest chargino. If the stop mass $m_{\tilde{t}} < m_{\tilde{\chi}_1^\pm}$, then $\tilde{\chi}_1^\pm \rightarrow \tilde{t} (\rightarrow c \tilde{\chi}_1^0) b$ so $\tilde{\nu} \rightarrow ebc \tilde{\chi}_1^0$; if $m_{\tilde{t}} > m_{\tilde{\chi}_1^\pm}$ then $\tilde{\chi}_1^\pm \rightarrow W^* (\rightarrow jj) \tilde{\chi}_1^0$ so $\tilde{\nu} \rightarrow ejj \tilde{\chi}_1^0$, where $j = u, d, s, c$. Either way, dominantly $\tilde{e}_L \tilde{\nu} \rightarrow ee\gamma jj \cancel{E}_T$ where j may be light or heavy quarks. If no such signal is found, probably the produced selectron was \tilde{e}_R . Also, $\sigma(\tilde{\nu}\tilde{\nu}) \cong \sigma(\tilde{e}_L \tilde{e}_L)$. Cross sections for many channels are given in Ref. [22].

The most interesting channel (in our opinion) at FNAL is $u\bar{d} \rightarrow W^+ \rightarrow \tilde{\chi}_i^+ \tilde{\chi}_2^0$. This gives a signature $\gamma jj \cancel{E}_T$, for which there is only small parton-level SM backgrounds. If $m_{\tilde{t}} < m_{\tilde{\chi}_i^\pm}$, one of j is a b . If $t \rightarrow \tilde{t} \tilde{\chi}_2^0$ (expected about 10% of

the time) and, if \tilde{q} are produced at FNAL, there are additional sources of such events (see below).

If charginos, neutralinos and sleptons are light, then gluinos and squarks may not be too heavy. If stops are light ($m_{\tilde{t}_1} \simeq M_W$), then $BR(t \rightarrow \tilde{t}\tilde{\chi}_i^0) \simeq 1/2$ [26]. In this case, an extra source of tops must exist beyond SM production, because $\sigma \times BR(t \rightarrow Wb)^2$ is near or above its SM value with $BR(t \rightarrow Wb) = 1$. With these motivations, the authors of [27] have suggested that one assume $m_{\tilde{g}} \geq m_t + m_{\tilde{t}}$ and $m_{\tilde{q}} \geq m_{\tilde{g}}$, with $m_{\tilde{q}} \simeq 250 - 300$ GeV. Then there are several pb of top production via channels $\tilde{q}\tilde{g}, \tilde{g}\tilde{g}, \tilde{q}\tilde{q}$ with $\tilde{q} \rightarrow q\tilde{g}$, and $\tilde{g} \rightarrow t\tilde{t}$ since $t\tilde{t}$ is the gluino's only two-body decay mode. This analysis points out that $P_T(t\tilde{t})$ should peak at smaller P_T for the SM than for the SUSY scenario, since the system is recoiling against extra jets in the SUSY case. The SUSY case suggests that if m_t or $\sigma_{t\tilde{t}}$ are measured in different channels one will obtain different values, which may be consistent with reported data. This analysis also argues that the present data is consistent with $BR(t \rightarrow \tilde{t}\tilde{\chi}_i^0) = 1/2$.

At present [28] R_b and $BR(b \rightarrow s\gamma)$ differ from their SM predictions by $1.5-2\sigma$, and α_s measured by the Z width differs by about $1.5-2\sigma$ from its value measured in DIS and other ways. If these effects are real they can be explained by $\tilde{\chi}_i^\pm - \tilde{t}$ loops, using the same SUSY parameters deduced from the $ee\gamma\gamma$ event (+ a light, mainly right-handed, stop). Although $\tan\beta, \mu$, and M_2 a priori could be anything, they come out the same from the analysis of these loops as from $ee\gamma\gamma$ ($\tan\beta \leq 1.5, \mu \sim -m_Z/2, M_2 \sim 60 - 80$ GeV).

The LSP= $\tilde{\chi}_1^0$ apparently escapes the CDF detector in the $ee\gamma\gamma$ event, suggesting it is stable (though only proving it lives longer than $\sim 10^{-8}$ sec). If so it is a candidate for CDM. The properties of $\tilde{\chi}_1^0$ are deduced from the analysis [22] so the calculation of the relic density [29] is highly constrained. The analysis shows that the s-channel annihilation of $\tilde{\chi}_1^0\tilde{\chi}_1^0$ through the Z dominates, so the needed parameters are $\tan\beta, m_{\tilde{\chi}_1^0}$ and the Higgsino fraction for $\tilde{\chi}_1^0$, which is large. The results are encouraging, giving $0.1 \leq \Omega h^2 \leq 1$, with a central value $\Omega h^2 \simeq 1/4$.

The parameter choices of Table IV can be input to event generators such as Spythia or ISAJET (via *MSSMi* keywords) to check that the event rate and kinematics of the $ee\gamma\gamma$ event are satisfied and then to determine other related signatures. Spythia includes the $\tilde{\chi}_2^0 \rightarrow \tilde{\chi}_1^0\gamma$ branching ratio for low $\tan\beta$ values; for ISAJET, the $\tilde{\chi}_2^0 \rightarrow \tilde{\chi}_1^0\gamma$ branching must be input using the *FORCE* command, or must be explicitly added into the decay table.

B. CDF/D0 Dilepton Plus Jets Events

Recently, CDF and D0 have reported various dilepton plus multi-jet events which are presumably top-quark candidate events. For several of these events, however, the event kinematics do not match well with those expected from a top quark with mass $m_t \sim 175$ GeV. The authors of Ref. [30] have shown that the match to event kinematics can be improved by hypothesizing a supersymmetry source for the recalcitrant events. The supersymmetry source is best matched by considering $\tilde{q}\tilde{q}$ production, where each $\tilde{q} \rightarrow q\tilde{\chi}, \tilde{\chi} \rightarrow \nu\tilde{\ell}, \tilde{\ell} \rightarrow \ell\tilde{\chi}_1^0$. A recom-

mended set of parameters is as follows [30]: $m_{\tilde{g}} \simeq 330$ GeV, $m_{\tilde{q}} \simeq 310$ GeV, $m_{\tilde{t}_L} \simeq 220$ GeV, $m_{\tilde{t}} \simeq 220$ GeV, $m_{\tilde{t}_R} \simeq 130$ GeV, $\mu \simeq -400$ GeV, $M_1 \simeq 50$ GeV and $M_2 \simeq 260$ GeV. Note that this parameter set discards the common hypothesis of gaugino mass unification. These parameters can be input into Spythia or ISAJET (via *MSSMi* keywords), taking care to use the non-unified gaugino masses as inputs.

VI. R Parity Violation

R parity (R) is a quantum number which is +1 for any ordinary particle, and -1 for any sparticle. R -violating (R) interactions occur naturally in supersymmetric theories, unless they are explicitly forbidden. Each R coupling also violates either lepton number L , or baryon number B . Together, these couplings violate both L and B , and lead to tree-level diagrams which would make the proton decay at a rate in gross violation of the observed bound. To forbid such rapid decay, such R couplings are normally set to zero. However, what if such couplings are actually present?

In supersymmetry with minimal field content, the allowable R part of the superpotential is

$$W_{\mathcal{R}} = \lambda_{ijk} L_i L_j \bar{E}_k + \lambda'_{ijk} L_i Q_j \bar{D}_k + \lambda''_{ijk} \bar{U}_i \bar{D}_j \bar{D}_k. \quad (8)$$

Here, L, Q, \bar{E}, \bar{U} , and \bar{D} are superfields containing, respectively, lepton and quark doublets, and charged lepton, up quark, and down quark singlets. The indices i, j, k , over which summation is implied, are generational indices. The first term in $W_{\mathcal{R}}$ leads to L -violating (L) transitions such as $e + \nu_\mu \rightarrow \tilde{e}$. The second one leads to L transitions such as $u + \bar{d} \rightarrow \tilde{e}$. The third one produces B transitions such as $\bar{u} + \bar{d} \rightarrow \tilde{d}$. To forbid rapid proton decay, it is often assumed that if R transitions are indeed present, then only the L -violating λ and λ' terms occur, or only the B -violating λ'' term occurs, but not both. While the flavor components of $\lambda'\lambda''$ involving u, d, s are experimentally constrained to be $< 10^{-24}$ from proton decay limits, the other components of $\lambda'\lambda''$ and $\lambda\lambda''$ are significantly less tightly constrained.

Upper bounds on the R couplings λ, λ' , and λ'' have been inferred from a variety of low-energy processes, but most of these bounds are not very stringent. An exception is the bound on λ'_{111} , which comes from the impressive lower limit of 9.6×10^{24} yr [31] on the half-life for the neutrinoless double beta decay ${}^{76}\text{Ge} \rightarrow {}^{76}\text{Se} + 2e^-$. At the quark level, this decay is the process $2d \rightarrow 2u + 2e^-$. If $\lambda'_{111} \neq 0$, this process can be engendered by a diagram in which two d quarks each undergo the R transition $d \rightarrow \tilde{u} + e^-$, and then the two produced \tilde{u} squarks exchange a \tilde{g} to become two u quarks. It can also be engendered by a diagram in which $2d \rightarrow 2\tilde{d}$ by \tilde{g} exchange, and then each of the \tilde{d} squarks undergoes the R transition $\tilde{d} \rightarrow u + e^-$. Both of these diagrams are proportional to $\lambda'_{111}{}^2$. If we assume that the squark masses occurring in the two diagrams are equal, $m_{\tilde{u}_L} \simeq m_{\tilde{d}_R} \equiv m_{\tilde{q}}$, the previously quoted limit on the half-life implies that [32]

$$|\lambda'_{111}| < 3.4 \times 10^{-4} \left(\frac{m_{\tilde{q}}}{100 \text{ GeV}} \right)^2 \left(\frac{m_{\tilde{g}}}{100 \text{ GeV}} \right)^{1/2}. \quad (9)$$

It is interesting to recall that if the amplitude for neutrinoless double beta decay is, for whatever reason, nonzero, then the electron neutrino has a nonzero mass [33]. Thus, if $\lambda'_{1jj} \neq 0$, SUSY interactions lead to nonzero neutrino mass [34].

The way [35] in which low-energy processes constrain many of the L couplings λ and λ' is illustrated by consideration of nuclear β^- decay and μ^- decay. In the Standard Model (SM), both of these decays result from W exchange alone, and the comparison of their rates tells us about the CKM quark mixing matrix. However, in the presence of R couplings, nuclear β^- decay can receive a contribution from \tilde{d} , \tilde{s} , or \tilde{b} exchange, and μ^- decay from \tilde{e} , $\tilde{\mu}$, or $\tilde{\tau}$ exchange. The information on the CKM elements which has been inferred assuming that only W exchange is present bounds these new contributions, and it is found, for example, that [35]

$$|\lambda_{12k}| < 0.04 \left(\frac{m_{\tilde{e}_R^k}}{100 \text{ GeV}} \right), \quad (10)$$

for each value of the generation index k . In a similar fashion, a number of low-energy processes together imply [35] that for many of the L couplings λ_{ijk} and λ'_{ijk} ,

$$|\lambda_{ijk}^{(\prime)}| < (0.03 \rightarrow 0.26) \left(\frac{m_{\tilde{f}}}{100 \text{ GeV}} \right). \quad (11)$$

Here, $m_{\tilde{f}}$ is the mass of the sfermion relevant to the bound on the particular $\lambda_{ijk}^{(\prime)}$.

Bounds of order 0.1 have also been placed on the L couplings λ'_{1jk} by searches for squarks formed through the action of these couplings in e^+p collisions at HERA [36].

Constraints on the B couplings λ'' come from nonleptonic weak processes which are suppressed in the SM, such as rare B decays and $K - \bar{K}$ and $D - \bar{D}$ mixing [37]. For example, the decay $B^+ \rightarrow \bar{K}^0 K^+$ is a penguin (loop) process in the SM, but in the presence of R couplings could arise from a tree-level diagram involving \tilde{u}_R^k ($k = 1, 2, \text{ or } 3$) exchange. The present upper bound on the branching ratio for this decay [38] implies that [37]

$$|\lambda''_{k12} \lambda''_{k23}|^{1/2} < 0.09 \left(\frac{m_{\tilde{u}_R^k}}{100 \text{ GeV}} \right); \quad k = 1, 2, 3. \quad (12)$$

Recently, bounds $\lambda'_{12k} < 0.29$ and $\lambda'_{22k} < 0.18$ for $m_{\tilde{q}} = 100$ GeV have been obtained from data on D meson decays [34]. For a recent review of constraints on R -violating interactions, see Ref. [39].

We see that if sfermion masses are assumed to be of order 100 GeV or somewhat larger, then for many of the R couplings λ_{ijk} , λ'_{ijk} and λ''_{ijk} , the existing upper bound is ~ 0.1 for a sfermion mass of 100 GeV. We note that this upper bound is comparable to the values of some of the SM gauge couplings. Thus, R interactions could still prove to play a significant role in high-energy collisions.

What effects of R might we see, and how would R interactions affect future searches for SUSY? Let us assume that R couplings are small enough that sparticle production and decay are still dominated by gauge interactions, as in the absence of R . The main effect of R is then that the LSP is no longer stable, but decays into ordinary particles, quite possibly within the detector in which it is produced. Thus, the LSP no longer carries away transverse energy, and the missing transverse energy (E_T) signal, which is the mainstay of searches for SUSY when R is assumed to be conserved, is greatly degraded. (Production of SUSY particles may still involve missing E_T , carried away by neutrinos.)

At future e^+e^- colliders, sparticle production may include the processes $e^+e^- \rightarrow \tilde{\chi}_i^+ \tilde{\chi}_j^-, \tilde{\chi}_i^0 \tilde{\chi}_j^0, \tilde{e}_L^+ \tilde{e}_L^-, \tilde{e}_L^+ \tilde{e}_R^-, \tilde{e}_R^+ \tilde{e}_L^-, \tilde{e}_R^+ \tilde{e}_R^-, \tilde{\mu}_L^+ \tilde{\mu}_L^-, \tilde{\mu}_R^+ \tilde{\mu}_R^-, \tilde{\tau}_L^+ \tilde{\tau}_L^-, \tilde{\tau}_R^+ \tilde{\tau}_R^-, \tilde{\nu}_L \tilde{\nu}_L$. Here, the $\tilde{\chi}_i^\pm$ are charginos, and the $\tilde{\chi}_i^0$ are neutralinos. Decay of the produced sparticles will often yield high- E_T charged leptons, which can be sought in seeking evidence of SUSY. Now, suppose the LSP is the lightest neutralino, $\tilde{\chi}_1^0$. If the L , R couplings λ are nonzero, the $\tilde{\chi}_1^0$ can have the decays $\tilde{\chi}_1^0 \rightarrow \mu \bar{\nu}_\mu, e \bar{\nu}_e$.

These yield high-energy leptons, so the strategy of looking for the latter to seek evidence of SUSY will still work. However, if the B , R couplings λ'' are nonzero, the $\tilde{\chi}_1^0$ can have the decays $\tilde{\chi}_1^0 \rightarrow c \bar{d} s, \bar{c} d \bar{s}$. When followed by these decays, the production process $e^+e^- \rightarrow \tilde{\chi}_1^0 \tilde{\chi}_1^0$ yields six jets which form a pair of three-jet systems. The invariant mass of each system is $m_{\tilde{\chi}_1^0}$, and there is no missing energy. This is quite an interesting signature.

Nonvanishing L and R couplings λ would also make possible resonant sneutrino production in e^+e^- collisions. [35] For example, we could have $e^+e^- \rightarrow \tilde{\nu}_\mu \rightarrow \tilde{\chi}_1^\pm \mu^\mp, \tilde{\chi}_1^0 \nu_\mu$. At the resonance peak, the cross section times branching ratio could be large [35].

In future experiments at hadron colliders, one can seek evidence of gluino pair production by looking for the multilepton signal that may result from cascade decays of the gluinos. This signal will be affected by the presence of R interactions. The worst case is where the LSP decays via B , R couplings to yield hadrons. The presence of these hadrons can cause leptons in SUSY events to fail the lepton isolation criteria, degrading the multilepton signal [40]. This reduces considerably the reach in $m_{\tilde{g}}$ of the Tevatron. At the Tevatron with an integrated luminosity of 0.1 fb^{-1} , there is no reach in $m_{\tilde{g}}$, while for 1 fb^{-1} it is approximately 200 GeV [40], if $m_{\tilde{q}} = 2m_{\tilde{g}}$. At the LHC with an integrated luminosity of 10 fb^{-1} , the reach extends beyond $m_{\tilde{g}} = 1 \text{ TeV}$, even in the presence of B and R interactions [41].

If R couplings are large, then conventional SUSY event generators will need many production and decay mechanisms to be re-computed. The results would be very model dependent, owing to the large parameter space in the R sector. If R couplings are assumed small, so that gauge and Yukawa interactions still dominate production and decay mechanisms, then event generators can be used by simply adding in the appropriate expected decays of the LSP (see the approach in Ref. [40, 41]). For ISAJET, the relevant LSP decays must be explicitly added (by hand) to the ISAJET decay table.

VII. Gauge-Mediated Low-Energy Supersymmetry Breaking

A. Introduction

Supersymmetry breaking must be transmitted from the supersymmetry-breaking sector to the visible sector through some messenger sector. Most phenomenological studies of supersymmetry implicitly assume that messenger-sector interactions are of gravitational strength. It is possible, however, that the messenger scale for transmitting supersymmetry breaking is anywhere between the Planck and just above the electroweak scale.

The possibility of supersymmetry breaking at a low scale has two important consequences. First, it is likely that the standard-model gauge interactions play some role in the messenger sector. This is because standard-model gauginos couple at the renormalizable level only through gauge interactions. If Higgs bosons received mass predominantly from non-gauge interactions, the standard-model gauginos would be unacceptably lighter than the electroweak scale. Second, the gravitino is naturally the lightest supersymmetric particle (LSP). The lightest standard-model superpartner is the next to lightest supersymmetric particle (NLSP). Decays of the NLSP to its partner plus the Goldstino component of the gravitino within a detector lead to very distinctive signatures. In the following subsections the minimal model of gauge-mediated supersymmetry breaking, and the experimental signatures of decay to the Goldstino, are presented.

B. The Minimal Model of Gauge-Mediated Supersymmetry Breaking

The standard-model gauge interactions act as messengers of supersymmetry breaking if fields within the supersymmetry-breaking sector transform under the standard-model gauge group. Integrating out these messenger-sector fields gives rise to standard-model gaugino masses at one-loop, and scalar masses squared at two loops. Below the messenger scale the particle content is just that of the MSSM plus the essentially massless Goldstino discussed in the next subsection. The minimal model of gauge-mediated supersymmetry breaking (which preserves the successful predictions of perturbative unification) consists of messenger fields which transform as a single flavor of $\mathbf{5} + \bar{\mathbf{5}}$ of $SU(5)$, i.e. there are triplets, q and \bar{q} , and doublets, ℓ and $\bar{\ell}$. These fields couple to a single gauge singlet field, S , through the superpotential

$$W = \lambda_3 S q \bar{q} + \lambda_2 S \ell \bar{\ell}. \quad (13)$$

A non-zero expectation value for the scalar component of S defines the messenger scale, $M = \lambda S$, while a non-zero expectation value for the auxiliary component, F , defines the supersymmetry-breaking scale within the messenger sector. For $F \ll \lambda S^2$, the one-loop visible-sector gaugino masses at the messenger scale are given by [42]

$$m_{\lambda_i} = c_i \frac{\alpha_i}{4\pi} \Lambda \quad (14)$$

where $c_1 = c_2 = c_3 = 1$ (we define $g_1 = \sqrt{\frac{5}{3}}g'$), and $\Lambda = F/S$. The two-loop squark and slepton masses squared at the messenger scale are [42]

$$\tilde{m}^2 = 2\Lambda^2 \left[C_3 \left(\frac{\alpha_3}{4\pi} \right)^2 + C_2 \left(\frac{\alpha_2}{4\pi} \right)^2 + \frac{3}{5} \left(\frac{Y}{2} \right)^2 \left(\frac{\alpha_1}{4\pi} \right)^2 \right] \quad (15)$$

where $C_3 = \frac{4}{3}$ for color triplets and zero for singlets, $C_2 = \frac{3}{4}$ for weak doublets and zero for singlets, and Y is the ordinary hypercharge normalized as $Q = T_3 + \frac{1}{2}Y$. The gaugino and scalar masses go roughly as their gauge couplings squared. The Bino and right-handed sleptons gain masses only through $U(1)_Y$ interactions, and are therefore lightest. The Winos and left-handed sleptons, transforming under $SU(2)_L$, are somewhat heavier. The strongly interacting squarks and gluino are significantly heavier than the electroweak states. Note that the parameter $\Lambda = F/S$ sets the scale for the soft masses (independent of the λ_i for $F \ll \lambda S^2$). The messenger scale M_i , may be anywhere between roughly 100 TeV and the GUT scale.

The dimensionful parameters within the Higgs sector, $W = \mu H_u H_d$ and $V = m_{12}^2 H_u H_d + h.c.$, do not follow from the ansatz of gauge-mediated supersymmetry breaking, and require additional interactions. At present there is no good model which gives rise to these Higgs-sector masses without tuning parameters. The parameters μ and m_{12}^2 are therefore taken as free parameters in the minimal model, and can be eliminated as usual in favor of $\tan \beta$ and m_Z .

Electroweak symmetry breaking results (just as for high-scale breaking) from the negative one-loop correction to $m_{H_u}^2$ from stop-top loops due to the large top quark Yukawa coupling. Although this effect is formally three loops, it is larger in magnitude than the electroweak contribution to $m_{H_u}^2$ due to the large squark masses. Upon imposing electroweak symmetry breaking, μ is typically found to be in the range $\mu \sim (1-2)m_{\tilde{\ell}_L}$ (depending on $\tan \beta$ and the messenger scale). This leads to a lightest neutralino, $\tilde{\chi}_1^0$, which is mostly Bino, and a lightest chargino, $\tilde{\chi}_1^\pm$, which is mostly Wino. With electroweak symmetry breaking imposed, the parameters of the minimal model may be taken to be

$$(\tan \beta, \Lambda = F/S, \text{sign } \mu, \ln M) \quad (16)$$

The most important parameter is Λ which sets the overall scale for the superpartner spectrum. It may be traded for a physical mass, such as $m_{\tilde{\chi}_1^0}$ or $m_{\tilde{\ell}_L}$. The low energy spectrum is only weakly sensitive to $\ln M_i$, and the splitting between $\ln M_3$ and $\ln M_2$ may be neglected for most applications.

C. The Goldstino

In the presence of supersymmetry breaking the gravitino gains a mass by the super-Higgs mechanism

$$m_G = \frac{F}{\sqrt{3}M_p} \simeq 2.4 \left(\frac{F}{(100 \text{ TeV})^2} \right) \text{eV} \quad (17)$$

where $M_p \simeq 2.4 \times 10^{18}$ GeV is the reduced Planck mass. With low-scale supersymmetry breaking the gravitino is naturally the

lightest supersymmetric particle. The lowest-order couplings of the spin- $\frac{1}{2}$ longitudinal Goldstino component of the gravitino, G_α , are fixed by the supersymmetric Goldberger-Treiman low energy theorem to be given by [43]

$$L = -\frac{1}{F} j^{\alpha\mu} \partial_\mu G_\alpha + h.c. \quad (18)$$

where $j^{\alpha\mu}$ is the supercurrent. Since the Goldstino couplings (18) are suppressed compared to electroweak and strong interactions, decay to the Goldstino is only relevant for the lightest standard-model superpartner (NLSP).

With gauge-mediated supersymmetry breaking it is natural that the NLSP is either a neutralino (as occurs in the minimal model) or a right-handed slepton (as occurs for a messenger sector with two flavors of $\mathbf{5} + \bar{\mathbf{5}}$). A neutralino NLSP can decay by $\tilde{\chi}_1^0 \rightarrow (\gamma, Z^0, h^0, H^0, A^0) + G$, while a slepton NLSP decays by $\tilde{\ell} \rightarrow \ell + G$. Such decays of a superpartner to its partner plus the Goldstino take place over a macroscopic distance, and for \sqrt{F} below a few 1000 TeV, can take place within a detector. The decay rates into the above final states can be found in Ref. [16, 17, 18, 19].

D. Experimental Signatures of Low-Scale Supersymmetry Breaking

The decay of the lightest standard-model superpartner to its partner plus the Goldstino within a detector leads to very distinctive signatures for low-scale supersymmetry breaking. If such signatures were established experimentally, one of the most important challenges would be to measure the distribution of finite path lengths for the NLSP, thereby giving a direct measure of the supersymmetry-breaking scale.

4. Neutralino NLSP

In the minimal model of gauge-mediated supersymmetry breaking, $\tilde{\chi}_1^0$ is the NLSP. It is mostly gaugino and decays predominantly by $\tilde{\chi}_1^0 \rightarrow \gamma + G$. Assuming R parity conservation, and decay within the detector, the signature for supersymmetry at a collider is then $\gamma\gamma X + \cancel{E}_T$, where X arises from cascade decays to $\tilde{\chi}_1^0$. In the minimal model the strongly interacting states are much too heavy to be relevant to discovery, and it is the electroweak states which are produced. At e^+e^- colliders $\tilde{\chi}_1^0$ can be probed directly by t -channel \tilde{e} exchange, yielding the signature $e^+e^- \rightarrow \tilde{\chi}_1^0 \tilde{\chi}_1^0 \rightarrow \gamma\gamma + \cancel{E}_T$. At a hadron collider the most promising signals include $qq' \rightarrow \tilde{\chi}_2^0 \tilde{\chi}_1^\pm, \tilde{\chi}_1^+ \tilde{\chi}_1^- \rightarrow \gamma\gamma X + \cancel{E}_T$, where $X = WZ, WW, W\ell^+\ell^-, \dots$. Another clean signature is $qq' \rightarrow \tilde{\ell}_R^+ \tilde{\ell}_R^- \rightarrow \ell^+\ell^-\gamma\gamma + \cancel{E}_T$. One event of this type has in fact been reported by the CDF collaboration [14]. In all these signatures both the missing energy and photon energy are typically greater than $m_{\tilde{\chi}_1^0}/2$. The photons are also generally isolated. The background from initial- and final-state radiation typically has non-isolated photons with a much softer spectrum.

In non-minimal models it is possible for $\tilde{\chi}_1^0$ to have large Higgsino components, in which case $\tilde{\chi}_1^0 \rightarrow h^0 + G$ can dominate. In this case the signature $bbbbX + \cancel{E}_T$ arises with the b -jets reconstructing m_{h^0} in pairs. This final state topology may be difficult

to reconstruct at the LHC – a systematic study has not yet been attempted.

Detecting the finite path length associated with $\tilde{\chi}_1^0$ decay represents a major experimental challenge. For the case $\tilde{\chi}_1^0 \rightarrow \gamma + G$, tracking within the electromagnetic calorimeter (EMC) is available. A displaced photon vertex can be detected as a non-zero impact parameter with the interaction region. For example, with a photon angular resolution of $40 \text{ mrad}/\sqrt{E}$ expected in the CMS detector with a preshower array covering $|\eta| < 1$ [44], a sensitivity to displaced photon vertices of about 12 mm at the 3σ level results. Decays well within the EMC or hadron calorimeter (HC) would give a particularly distinctive signature. In the case of decays to charged particles, such as from $\tilde{\chi}_1^0 \rightarrow (h^0, Z^0) + G$ or $\tilde{\chi}_1^0 \rightarrow \gamma^* + G$ with $\gamma^* \rightarrow f\bar{f}$, tracking within a silicon vertex detector (SVX) is available. In this case displaced vertices down to the 100 μm level should be accessible. In addition, decays outside the SVX, but inside the EMC, would give spectacular signatures.

4. Slepton NLSP

It is possible within non-minimal models that a right-handed slepton is the NLSP, which decays by $\tilde{\ell}_R \rightarrow \ell + G$. In this case the signature for supersymmetry is $\ell^+\ell^- X + \cancel{E}_T$. At e^+e^- colliders such signatures are fairly clean. At hadron colliders some of these signatures have backgrounds from WW and $t\bar{t}$ production. However, $\tilde{\ell}_L \tilde{\ell}_L$ production can give $X = 4\ell$, which has significantly reduced backgrounds. In the case of $\tilde{\ell}_R \tilde{\ell}_R$ production the signature is nearly identical to slepton pair production with $\tilde{\ell} \rightarrow \ell + \tilde{\chi}_1^0$ with $\tilde{\chi}_1^0$ stable. The main difference here is that the missing energy is carried by the massless Goldstino.

The decay $\tilde{\ell} \rightarrow \ell + G$ over a macroscopic distance would give rise to the spectacular signature of a greater than minimum ionizing track with a kink to a minimum ionizing track. Note that if the decay takes place well outside the detector, the signature for supersymmetry is heavy charged particles rather than the traditional missing energy.

E. Event Generation

For event generation by ISAJET, the user must provide a program to generate the appropriate spectra for a given point in the above parameter space. The corresponding $MSSMi$ parameters can be entered into ISAJET to generate the decay table, except for the NLSP decays to the Goldstino. If $NLSP \rightarrow G + \gamma$ at 100%, the *FORCE* command can be used. Since the G particle is not currently defined in ISAJET, the same effect can be obtained by forcing the NLSP to decay to a neutrino plus a photon. If several decays of the NLSP are relevant, then each decay along with its branching fraction must be explicitly added to the ISAJET decay table. Decay vertex information is not saved in ISAJET, so that the user must provide such information. In *Spythia*, the G particle is defined, and decay vertex information is stored.

VIII. Conclusions

In this report we have looked beyond the discovery of supersymmetry, to the even more exciting prospect of probing the new physics (of as yet unknown type) which we know must be associated with supersymmetry and supersymmetry breaking. The collider experiments which disentangle one weak-scale SUSY scenario from another will also be testing hypotheses about new physics at very high energies: the SUSY-breaking scale, intermediate symmetry-breaking scales, the GUT scale, and the Planck scale.

We have briefly surveyed the variety of ways that weak-scale supersymmetry may manifest itself at colliding beam experiments. We have indicated for each SUSY scenario how Monte Carlo simulations can be performed using existing event generators or soon-to-appear upgrades. In most cases very little simulation work has yet been undertaken. Even in the case of minimal supergravity the simulation studies to date have mostly focused on discovery reach, rather than the broader questions of parameter fitting and testing key theoretical assumptions such as universality. Clearly more studies are needed.

We have seen that alternatives to the minimal supergravity scenario often provide distinct experimental signatures. Many of these signatures involve displaced vertices: the various NLSP decays, LSP decays from R parity violation, chargino decays in the 200 and O-II models, and enhanced b multiplicity in the 24 model. This observation emphasizes the crucial importance of accurate and robust tracking capabilities in future collider experiments.

The phenomenology of some scenarios is less dramatic and thus harder to distinguish from the bulk of the mSUGRA parameter space. In any event, precision measurements will be needed in the maximum possible number of channels. In the absence of a “smoking gun” signature like those mentioned above, the most straightforward way to identify variant SUSY scenarios will be to perform an overconstrained fit to the mSUGRA parameters. Any clear inconsistencies in the fit should point to appropriate alternative scenarios. More study is needed of how to implement this procedure in future experiments with real-world detectors and data.

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