Tevatron and LEP-II Probes of Minimal and String-Motivated Supergravity Models

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Abstract

We explore the ability of the Tevatron to probe Minimal Supersymmetry with high energy scale boundary conditions motivated by supersymmetry breaking in the context of supergravity/superstring theory. A number of boundary condition possibilities are considered: dilaton-like string boundary conditions applied at the standard GUT unification scale or alternatively at the string scale; and extreme (“no-scale”) minimal supergravity boundary conditions imposed at the GUT scale or string scale. For numerous specific cases within each scenario the sparticle spectra are computed and then fed into ISAJET 7.07 so that explicit signatures can be examined in detail. We find that, for some of the boundary condition choices, large regions of parameter space can be explored via same-sign dilepton and isolated trilepton signals. For other choices, the mass reach of Tevatron collider experiments is much more limited. We also compare mass reach of Tevatron experiments with the corresponding reach at LEP 200.

1. Introduction

Assessing our ability to experimentally probe supersymmetric extensions of the Standard Model (SM) at existing and future accelerators is a crucial issue for the future of high energy physics. Indeed, $N = 1$ supersymmetric models containing Standard Model matter and gauge fields (and their superpartners) along with exactly two Higgs doublets are remarkable in that the observed values of $\alpha_{\text{QED}}$, $\sin^2 \theta_W$ and $\alpha_s$ at the scale $m_Z$ are highly consistent with unification of the $U(1), SU(2)_L$ and $SU(3)$ gauge coupling constants at a scale of order $M_U \sim 2 \times 10^{16}\text{GeV}$.\textsuperscript{[1]} (Although additional singlet Higgs fields do not affect the unification, we shall focus here on the Minimal Supersymmetric Model (MSSM) in which there are only two Higgs doublet fields, and no extra Higgs singlet field(s).) In a completely general context, the large uncertainty in the soft-supersymmetry-breaking parameters of the MSSM makes it difficult to arrive at definite predictions for the best probes and ultimate experimental accessibility of the superparticles. Even the basic superpartner mass scales (which we generically denote by $M_{\text{SUSY}}$) are rather uncertain, although it is widely accepted that they should lie below about 1 TeV in order to provide an obvious solution to the naturalness problem for the Higgs mass, and, in addition the gauge-coupling unification is only successful if $M_{\text{SUSY}} \lesssim 1\text{TeV}$. However, the success of gauge-coupling unification suggests that we should consider models that also have relatively simple and universal boundary condi-
tions for the soft-supersymmetry-breaking parameters at the unification scale. Implications at low-energies ($\lesssim 1$ TeV) can then be obtained by renormalization group evolution of the high-energy-scale parameters.

Supergravity and superstring theories provide the most attractive context in which gauge unification based on minimal $\tilde{N} = 1$ supersymmetry can be natural. Superstring theory stands out as the only candidate which is known to lead to a consistent theory of quantum gravity. Physics below the Planck scale is determined by the effective non-renormalizable supergravity theory that is believed to arise from the string once the super-heavy $M_P$-scale fields are integrated out. Our goal in this paper will be to assess the extent to which the Tevatron and LEP-II can probe the superparticle spectrum of the MSSM with soft supersymmetry breaking specified by boundary conditions (at the unification scale) as predicted in a limited, but very attractive, set of string and minimal supergravity models. The main focus of the paper will be on assessing a wide range of possible signals at the Tevatron, including the missing-transverse-momentum, same-sign-lepton, tri-lepton, and four-lepton discovery modes. We shall contrast the parameter space range for which the Tevatron can probe the superparticle spectrum with that for which supersymmetry can be observed at LEP-II via chargino-pair, slepton-pair and/or $Z$+Higgs associated production.

The outline of the paper follows. In Sec. 2, we discuss the motivation behind and nature of the string theory “dilaton-like” and “no-scale” minimal supergravity boundary conditions that we employ. In Sec. 3, we delineate the (two-dimensional) parameter spaces that are allowed for each of the eight resulting models, given existing theoretical and experimental constraints and our assumed top-quark mass of $m_t(m_t) = 170$ GeV. The all-important mass spectra and the consequent general phenomenological implications are also detailed in Sec. 3. In Sec. 4, we specify the parameters for the specific scenarios (within each of the eight models) that will be explored at the Tevatron using detailed Monte Carlo simulations. The selected scenarios are particularly chosen to sample the range of parameter space at the edge of Tevatron sensitivity. In Sec. 5, we give details of the simulations and cuts that we employ to analyze and isolate the different types of Tevatron discovery signals. In Sec. 6, we give the numerical results for Tevatron signals and backgrounds for the scenarios specified in Sec. 4. By examining these results as a function of scenario location in parameter space, we determine the portion of the parameter space of each of our eight models for which a supersymmetric signal will be detectable at the Tevatron. Substantial sensitivity to specific boundary condition and unification scale choices emerges. In Sec. 7, we survey the ability of LEP-II to explore the parameter spaces of each of the eight models, and draw comparisons to the Tevatron results. The substantial complementarity of the two machines is discussed. We present final remarks and conclusions in Sec. 8. Earlier analyses of selected Tevatron and/or LEP-II signals for models similar to those considered here appear in Refs. [2-4].

2. String and Minimal Supergravity Models and Boundary Conditions

In string theory the unification scale and the gauge couplings are determined dynamically at tree-level in terms of the vacuum expectation values of the dilaton field, with one-loop corrections coming from moduli field terms. In general, gauge coupling unification is not dependent on a grand unifying group, but instead takes the form: $g_1^2 k_1 = g_2^2 k_2 = g_3^2 k_3$, where the $k_i$ are non-Abelian gauge factors called the Kac-Moody levels. Phenomenologically consistent gauge coupling unification requires $k_3 = k_2 = \frac{3}{5} k_1$. This is, in fact, the prediction of the simplest and most attractive string theories, where one finds Kac-Moody levels $k_3 = k_2 = \frac{4}{5} k_1 = 1$. 

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The most obvious difficulty with the string approach is that the string scale (at which the unification boundary conditions would naively be expected to occur) is determined to be $M_S = 0.7g_U \times 10^{18} \text{ GeV}$, i.e. of order $M_P/\sqrt{8\pi}$ (where $M_P$ is the Planck scale — $M_P = 1.2 \times 10^{19} \text{ GeV}$ — at which quantum effects of gravity must be considered), as compared to the somewhat lower MSSM unification scale of $M_U \sim (2 \times 10^{16}) \text{ GeV}$. However, it is now known that there can be significant threshold corrections associated with the infinite number of Planck-scale states. Calculations in specific cases show that these threshold corrections can effectively cause the couplings at $M_S$ to differ from one another in just such a manner as to yield an effective unification at the lower $M_U$ scale. It is also possible that the unification point, $M_U$, is higher than ‘naively’ predicted (i.e. on the basis of the minimal Standard Model (SM) particles plus associated superpartners plus Higgs doublets) due to the presence of extra vector-like multiplets with masses at an intermediate scale between $m_Z$ and $M_U$. For appropriately chosen masses and representations these can modify the gauge coupling running so as to yield coupling unification at $M_S$. This solution is apparently required in the $4 - D$ free-fermionic string formulation, where the predicted threshold effects only serve to increase the unification scale, and in certain orbifold constructions. For either resolution, the phenomenology of the string models that yield the minimal supersymmetric model as the $N = 1$ low-energy supersymmetric sector is clearly worth examining. By comparing the two approaches, we will gain a first indication of the sensitivity of phenomenology to assumptions about unification-scale physics.

As noted earlier, generically the MSSM has many independent parameters beyond those required in specifying the Standard Model — namely, the magnitudes of soft-supersymmetry-breaking potential terms. The latter are parameterized by scalar masses ($m$), soft Yukawa coupling coefficients ($A$), gaugino masses ($M_a$, where $a$ denotes the group), and the coefficient ($B$) specifying the soft scalar Higgs field mixing term. Fortunately, supergravity and superstring theory both provide insight into the general structure of the soft terms. In supergravity, rather mild assumptions yield universal soft-SUSY breaking, dramatically reducing the number of parameters. In string models there has been much progress in classifying the possible SUSY-breaking sources; associated ‘string-inspired’ forms of soft-SUSY breaking emerge, specified by a relatively small number of $M_U$-scale parameters. Rather universal soft-SUSY breaking can easily emerge. Phenomenologically, FCNC constraints are most easily satisfied if the soft-squark masses are, in fact, generation independent. Of course, in principle a completely self-consistent, calculable string theory should be able to predict all the MSSM parameters. However, we are far from realizing this goal, as numerous difficult and unresolved issues in string theory still exist, including the presumably non-perturbative effects that are important in determining the true vacuum and the details of supersymmetry breaking.

In this paper we shall consider two basic types of boundary conditions for the soft-supersymmetry-breaking parameters at high energy scales. For the first type, we adopt a structure for the soft terms that emerges in a large number of string models — namely $M^0 = -A^0 = \sqrt{3}m^0$, where $m^0$, $A^0$ and $M^0$ are the (common) values for the soft scalar masses, $A$ coefficients, and $M_a$ gaugino masses, respectively. (The superscript 0 denotes values at the high energy scale.) These boundary conditions arise in the universal dilaton-dominated limit of SUSY breaking in all string models, but they are actually much more general, as we outline below. As we shall see, the small value of $m^0$ compared with $M^0$ leads to rather light spectra for the sleptons and sneutrinos of the model. For comparison, we also consider the extreme boundary conditions which yield the lightest possible slepton sector:
\( m^0 = A^0 = 0 \). These latter boundary conditions, supplemented by taking a universal value for the gaugino mass, \( M_0 = M^0 \), are commonly (but perhaps wrongfully) associated with the “no-scale model” label in the literature.

As well as specifying the boundary conditions, the high energy scale at which they are applied must be specified. In this paper we study the consequences of implementing the above two types of boundary conditions at two different high energy scales: \( i \) the standard GUT scale \( M_U \) determined from gauge coupling unification in the absence of any additional matter fields beyond those contained in the minimal supersymmetric model; or, alternatively, \( ii \) the substantially higher string scale. In this latter case, extra matter representations must be introduced — our choices will be detailed below.

### 2.1 String-motivated boundary conditions at \( M_U \)

A systematic analysis of soft-supersymmetry-breaking terms in specific four-dimensional string theories has been presented in Refs. [10,11]. In these references, the role of the dilaton \( D \) and overall \( \text{i.e.} \) associated with the volume or size of the manifold) moduli \( T \) fields is emphasized. Starting with specific compactification choices and the appropriate Kahler potential, explicit forms for the soft-supersymmetry breaking terms can be derived when the dilaton and/or moduli fields acquire non-zero vacuum expectation values as a result of spontaneous supersymmetry breaking. In this approach, the cosmological constant is not automatically zero; setting it to zero further simplifies the soft-breaking boundary conditions. Following the notation of Ref. [11], the soft parameters take particularly simple forms when expressed in terms of the goldstino angle \( \theta \), which specifies the extent to which the source of supersymmetry breaking resides in the dilaton versus moduli sector. If supersymmetry breaking is dominated by the dilaton superfield then \( \theta = \pi/2 \), whereas supersymmetry breaking dominance by the overall moduli superfield occurs in the \( \theta = 0 \) and \( \pi \) limits. For a given string model, the standard soft parameters \( m^0, M^0, A^0, \) and \( B^0 \) can all be expressed in terms of just \( m_{3/2} \) (the gravitino mass) and \( \theta \). For general values of \( \theta \), the precise expressions are model-dependent, although in the dilaton limit of \( \theta = \pi/2 \) (\( \sin \theta = 1 \)), the soft terms take a model-independent universal form (up to small corrections).

We consider a simplified subset of the models explored in Ref. [11]. First, we assume that the Kac-Moody levels of the three gauge groups are related by \( k_3 = k_2 = \frac{3}{5} k_1 = 1 \). If we recall that \( f_a = k_a S \) at tree level, where \( f_a \) are the inverse squared gauge coupling constants at the string scale, we see that this choice will be consistent with the experimentally observed coupling unification if one neglects corrections arising from the difference between \( M_S \) and \( M_U \). Generally, one-loop threshold corrections can alter this relationship somewhat (see Eq. (2.12) of Ref. [11]), perhaps even allowing the apparent unification at \( M_U \) for tree-level \( k_i \) values (related as above) to be consistent with threshold-corrected unification at \( M_S \). Second, we neglect any CP violating phases for the \( A^0 \) and \( B^0 \) parameters. Finally, in the case of orbifold compactification we assume that all fields belong to the untwisted sector \( \text{i.e.} \) we take the modular weights to be \( n_i = -1 \) for all fields \( i \); see Eq. (3.9) of Ref. [11]). This set of choices is certainly the simplest possibility within the context of the four-dimensional superstring models considered in Refs. [10,11].

With these choices, the soft terms for both Calabi-Yau and orbifold compactifications
take the form:

\[ m_i^0 = m_{3/2}^2 [1 - \kappa \cos^2 \theta] \]
\[ A^0 = -\sqrt{3}m_{3/2} \sin \theta \]  
\[ M_a^0 = \sqrt{3}m_{3/2} \sin \theta + m_{3/2} \cos \theta X_a , \]  

where \( a \) labels the gauge group. In the absence of threshold corrections, \( \kappa = 1 \) and \( X_a = 0 \). However, threshold corrections are generally present. For orbifold compactifications (with all \( n_i = -1 \)) one-loop threshold corrections give rise to

\[ \kappa = (1 - \frac{\delta_{GS} \pi^2 Y}{24})^{-1} , \quad X_a \propto b_a - k_a \delta_{GS} \]

where \( Y \) is computable in terms of the dilaton and moduli chiral superfields and numerically is of order 4 or 5, \( b_a \) is the one-loop \( \beta \)-function coefficient for the particular group, and \( \delta_{GS} \) is a model-dependent quantity (often a negative integer). For Calabi-Yau compactifications, much less is known about the threshold corrections, although it is quite probable that \( \kappa \) is not precisely 1 nor \( X_a \) exactly 0.\[^{[12]}\]

Finally, we note that approximate results for the \( B_0^\mu \) parameter were obtained in Ref. [11]. For example, if Higgs superfield mixing appears only in the standard \( \mu \hat{H}_1 \hat{H}_2 \) superpotential term one finds

\[ B_\mu^0 = m_{3/2}(-1 - \sqrt{3} \sin \theta - \cos \theta) . \]  

Another source of \( B_0^0 \) derives from an additional Higgs-mixing term in the Kahler potential that can generally be present in Calabi-Yau compactifications, but is not present for orbifold compactifications. The resulting form of \( B_0^0 \) in the absence of the \( \mu \hat{H}_1 \hat{H}_2 \) superpotential term is

\[ B_0^Z = 2m_{3/2}(1 + \cos \theta) . \]

If both sources of \( B_0^0 \) are simultaneously present, the resulting form of \( B_0^0 \) as a function of \( \theta \) would be more complicated. And the above forms themselves were obtained only after a significant number of approximations and assumptions. Thus, we shall leave \( B_0^0 \) as a free parameter subject only to the requirement that the model be consistent with correct EWSB and phenomenological constraints. However, we shall later describe the values taken on by \( B_0^0 \) within the allowed parameter space regions; we shall see that the models would be very strongly constrained or eliminated altogether for particular choices of \( B_0^0 \). In this regard, it is useful to note that in the context of the above approximate forms, \( B_\mu^0 \leq 0 \), while \( B_0^Z \geq 0 \), with zero values only being reached for \( \theta = -\pi \).

From Eq. (1) we extract the dilaton-dominated model by setting \( \sin \theta = 1 \), which gives identical \( A^0, m_i^0 \) and \( M_a^0 \) results for both Calabi-Yau and orbifold compactification (as noted earlier) and is actually completely independent of the \( n_i \) choices. Indeed, to the extent that threshold corrections can be neglected (in Ref. [11] this is estimated to be true for
sin θ ≳ 0.05), the orbifold (with all \( n_i = -1 \)) and Calabi-Yau model results for \( m_i^0, A^0, \) and \( M_a^0 \) are identical, and are indistinguishable from the strict dilaton-dominated model after the rescaling \( m_{3/2} \rightarrow m_{3/2} \sin \theta \). If one could trust one of the above quoted results for \( B^0 \), then it would provide some discrimination between different \( \theta \) values. However, the uncertainty in \( B^0 \) makes it much more appropriate to allow it to be a free parameter, in which case all models become equivalent for the very large range of \( \theta \) values such that threshold corrections can be neglected. This large class of models, specified by the boundary conditions

\[
M^0 = -A^0 = \sqrt{3} m_i^0,
\]

will be denoted by the symbol \( D \) (for dilaton-equivalent). Leaving \( B^0 \) (as well as \( M^0 \)) free will thus have a large range of validity, and it is only in the sin \( \theta \sim 0 \) case, or by taking some of the \( n_i \) different from \(-1\) in the orbifold models, that SUSY breaking can become model dependent in a manner that goes beyond the boundary conditions that we shall employ.

2.2 Extreme minimal supergravity boundary conditions at \( M_U \)

The boundary conditions obtained in the moduli-dominated limit (corresponding to sin \( \theta = 0 \) in of Eq. (1)) are quite different from the dilaton-equivalent constraints of Eq. (5). First, we note that the scale of \( M_a \) is set by \( X_a \), which therefore cannot be too small on purely phenomenological grounds. More generally, the relative sizes of the scalar masses and the gaugino masses are determined by the relative magnitudes of \( \kappa - 1 \) and \( X_a \). Even in the specific case of orbifold \((n_i = -1)\) moduli-dominated models there are many possibilities. In the simplest case of \( \delta_{GS} = 0, m_i^0 = A^0 = 0 \) and non-universal values for the gaugino masses \( M_a^0 \propto b_a \) are required at \( M_U \). In contrast, for significant (negative, generally negative-integer) non-zero values of \( \delta_{GS}, m_i^0 \) is typically substantially larger than \( M_a^0 \) in the moduli-dominated limit. In the example of \( \delta_{GS} = -5 \) explored in Ref. [11], gaugino masses become smaller than scalar masses for sin \( \theta < \sim 0.05 \). As noted there, this is not unlike the situation obtained in explicit gaugino-condensation models.\[^{[13]}\] The \( \delta_{GS} = 0 \) case, although regarded as atypical (in that \( S - T \) mixing will generally be present in the Kahler potential), represents an interesting extreme case of boundary conditions with specific non-universal gaugino masses at \( M_U \).

The extreme case of \( \delta_{GS} = 0 \) with sin \( \theta = 0 \), yielding the \( M_a^0 \) and \( B^0 \) as the only seeds of supersymmetry breaking, has much in common with the so-called ‘no-scale’ scenario.\[^{[14]}\] The original motivations for no-scale models were i) to guarantee a vanishing cosmological constant at the unification scale, and ii) to yield a flat potential (in a scalar field direction) such that the electroweak symmetry breaking (EWSB) scale is to be generated dynamically. The no-scale scenario requires a Kahler potential of a very specific form (which is in fact realized in certain free-fermionic constructions\[^{[15]}\]). Of course, soft-SUSY breaking introduces a scale, but the underlying motivations and structure of the ‘no-scale’ model are least disturbed for a very simple and specific set of soft-SUSY parameter boundary conditions; namely zero values for the soft scalar masses and \( A \) terms at \( M_U \), but non-zero values for the gaugino
masses. In the simplest models, the chiral superfield density \( f_{\alpha\beta} \) is proportional to \( \delta_{\alpha\beta} \), and gaugino masses take on a universal value at \( M_U \). The resulting boundary conditions are:

\[
m_{i}^{0} = A^{0} = 0; \quad M_{a}^{0} = M^{0}.
\]  

(6)

As in the dilaton-equivalent case, specific choices for the \( B^{0} \) parameter are less motivated, and we will allow any value for \( B^{0} \) consistent with electroweak symmetry breaking in the renormalization group approach. We comment later on the values that \( B^{0} \) actually takes on in the parameter space regions that are consistent with other constraints; \( B^{0} = 0 \) is an interesting special case in that gaugino masses then provide the sole seed for supersymmetry breaking.

The boundary conditions of Eq. (6) are best viewed as specifying an extreme version of the Minimal Supergravity model, for which we adopt the the generic symbol \( MS \). As outlined above, they differ from the boundary conditions that are obtained in the special moduli-dominated \( \delta_{GS} = 0 \) string theory limit where the \( M_{a}^{0} \) are required to be non-universal. To date, no detailed string model has resulted in precisely the minimal-supergravity boundary conditions of Eq. (6) with universal gaugino mass. Nonetheless, the \( MS \) boundary conditions have a long history and are the simplest that can be devised which satisfy the basic constraints outlined below. First, since setting \( M_{a}^{0} = 0 \) would be inconsistent with experimental limits on gaugino masses, Eq. (6) is the simplest choice for which only one of \( A^{0}, m_{i}^{0}, M_{a}^{0} \) is taken to be non-zero. In addition, \( m_{i}^{0} = 0 \) guarantees that the squark masses are sufficiently degenerate after renormalization group evolution to avoid flavor changing neutral current difficulties. (Of course, a universal value for the \( m_{i}^{0} \) as in the \( D \) models, also achieves the same end.) Setting \( A^{0} = 0 \) removes the CP violation that might otherwise be present at a level inconsistent with constraints on the neutron and electron electric dipole moments should \( A^{0} \) have a non-trivial phase with respect to \( M^{0} \). This is because the driving terms in the renormalization group equations (RGE’s) for the various \( A \)’s are proportional to the \( M_{a} \)’s (with real coefficients) so that no phase for any \( A/M_{a} \) can be generated if \( A^{0} = 0 \).

There could still be a non-trivial phase for \( B^{0}/M^{0} \); in our analysis we consider only real values for \( B^{0} \).

2.3 Summary of \( M_U \) boundary conditions

To summarize: in the dilaton-equivalent models we require \( M^{0} = -A^{0} = \sqrt{3}m^{0} \) at \( M_U \), while in the minimal-supergravity models we require \( m^{0} = A^{0} = 0 \) at the scale \( M_U \), and a universal value, \( M_{a}^{0} = M^{0} \) for the gaugino masses — in both models all values of \( B^{0} \) are allowed that yield a correct pattern of electroweak symmetry breaking and that are consistent with existing experimental and phenomenological constraints as detailed below.

So far, we have not discussed the \( \mu \) parameter required to complete the specification of the \( M_U \) boundary conditions. Although the magnitude of \( \mu^{0} \) is determined in the RGE approach in terms of the other \( M_U \)-scale parameters by minimizing the scalar potential, the sign of \( \mu \) is undetermined; we will consider both the positive and negative sign possibilities. This leaves us with two \( D \) and two \( MS \) models that we will explore phenomenologically. We denote these four models by \( D^{+}, D^{-}, MS^{+}, \text{ and } MS^{-} \). We re-emphasize that by leaving the \( B^{0} \) parameter free (subject only to phenomenological and RGE constraints), the
$D$ models actually describe a broad class of superstring-motivated models in which threshold corrections can be neglected, while the $MS$ models are more general than those in which $B^0 = A^0 - m^0^{[14]}$ or $B^0 = 0$ is imposed.

2.4 Boundary conditions at $MS$

As in the case of the gauge couplings, the use of the boundary conditions (5) or (6) at scale $M_U$ as opposed to $M_S$ can be questioned. Certainly, it is simplest to presume that the soft-parameter boundary conditions apply at the same scale at which the coupling constants begin their independent evolutions from a common value. A priori, it cannot be ruled out that the hidden-sector/threshold corrections conspire so that both coupling constant equality and universal values for the soft masses and the $A$ parameters all apply at $M_U$ rather than $M_S$. It is probably too early to even rule out the possibility that the hidden-sector-determined scale at which all evolution begins is significantly below $M_S$.

However, as an alternative, one can presume that the threshold/hidden-sector effects are small or (possibly) serve to push $M_U$ even higher, in which case it is most natural to presume that the coupling constants take a common value at $M_S$. To achieve coupling unification as this higher scale, it is necessary to introduce intermediate-to-GUT-scale (or gap) representations that modify the gauge coupling running at high scales in precisely the correct manner.$^{[16]}$ Of course, this approach implicitly assumes that the nice meeting of the couplings in the MSSM (with minimal field content) is quite accidental. The minimal model along these lines is that suggested in Ref. [17], in which vector-like heavy ‘quark’ field representations of the type

$$Q_L = (3, 2, 1/3), \quad Q'_L = (\bar{3}, 2, -1/3), \quad D^c_R = (3, 1, 2/3), \quad D_R = (\bar{3}, 1, -2/3),$$

are introduced. By inputing $\alpha_s$ and $\sin^2 \theta_W$, and requiring coupling unification at $M_S$, the masses of these gap fields can be determined. The masses required for unification are then $m_{D_R} \simeq 10^7 \text{GeV}$ and $m_{Q_L} \simeq 5 \times 10^{12} \text{GeV}$. Contributions to the $b_i$ are positive for both types of fields, thereby increasing the slopes of the running $\alpha_i$ as a function of the scale, which raises the unification scale as well as $\alpha_U$. Indeed, these extra representations even have a natural home in the flipped $SU(5) \times U(1)$ model: the $Q_L$’s are the fermionic partners of the $X, Y$ gauge bosons and the $(3, 2)$ components of the $10$ Higgs representation; the $D_L$’s are the partners of the Higgs triplet $(3, 1)$ components of the $5$ and $10$ Higgs representations. (Of course, the $Q'_L$ and $D^c_R$ are associated with the conjugate states to those listed above.) The flipped model mass matrices can even yield masses for the $Q$ and $D$ fermions of the required magnitude.$^{[4]}$ We will discuss how the phenomenology changes upon employing this approach and applying the dilaton-equivalent and minimal-supergravity boundary conditions of Eqs. (5) and (6) at $M_S$. The models so generated will be denoted by $SD^\pm$ and $SMS^\pm$ (for string-scale-unified dilaton-equivalent and string-scale-unified minimal-supergravity), where, as before, the models are determined by choosing just two parameters and the sign of $\mu$ (indicated by the superscripts). The masses of the $D$ and $Q$ fields are not independent, being entirely determined by the requirement that unification occur at $M_S$.

Of course, it must be admitted that the choice of unifying at $M_S$ employing intermediate scale representations that are specific to $SU(5) \times U(1)$ is somewhat arbitrary. Indeed, as
noted earlier, there is no need for the super-string to have any group structure beyond the basic $SU(3) \times SU(2)_L \times U(1)$ of the SM. Nonetheless, we hope our results will be representative of those that one would obtain for other reasonably simple choices for the gap fields.

The main effect of raising the unification scale from $M_U$ to $M_S$ is the increased slepton masses (at a given gluino mass) due to the increased amount of evolution lever arm. This effect is clearly greater (on a percentage basis) for the $MS$ models where $m^0 = 0$ (and the entire slepton mass at low energy derives from evolution) as compared to the $D$ models for which $m^0$ is a significant fraction of $M^0$. Important phenomenological differences occur if the mass hierarchy of the sleptons, sneutrinos, charginos and neutralinos is altered. We shall see that mass hierarchies are altered in going from $D^-$ to $SD^-$ and from $MS^\pm$ to $SMS^\pm$. Some amusing patterns will emerge. In particular: the $D^-$ mass hierarchies are converted to the ordering found for $D^+, SD^+$; the $SMS^+$ hierarchies are quite different than for $MS^+$, and are closely related to those for $D^+$; and the $SMS^-$ hierarchies are very similar to those for $D^-$. Thus, both the $SD^+$ and $SD^-$ will have phenomenology similar to that found for $D^+$, while $SMS^-$ results will be very like those for $D^-$. In combination, the eight model scenarios investigated illustrate the vital role played by precise unification boundary condition and scale choices when supersymmetry breaking is dominated by gaugino masses.

2.5 Final introductory remarks

Thus, the models we explore are motivated in two ways. On the theoretical side, supergravity and string theory provide a strong motivation for the types of boundary conditions considered. On the phenomenological side, the boundary conditions employed are distinguished by the rather low masses obtained for the sleptons. This latter provides a very interesting alternative to models in which boundary conditions with large $m^0$ arise, such as gaugino-condensate models and models based on the above-mentioned moduli-dominated string scenarios (with $\delta_{GS} \neq 0$). Light sleptons will turn out to have many crucial phenomenological consequences (such as two-body decays of charginos and neutralinos to slepton-lepton final states) that dramatically alter the phenomenology as compared to models where sleptons have mass of order the gluino mass or higher (so that charginos and neutralinos tend to decay via three-body modes to virtual or real $W,Z$ plus lighter gaugino). We shall see that the predicted slepton masses are generally sufficiently smaller than the gluino mass that slepton-pair production at LEP-II can easily be a much deeper probe of the model parameter space than the Tevatron, whereas the reverse can easily be true if sleptons are heavy. Light sleptons also result in many special situations and much greater variability in the discovery potential at the Tevatron. Signals for supersymmetry will often be more likely to arise from neutralino/chargino-pair and slepton-pair production than from gluino-pair, gluino-squark, or squark-pair production. (Squarks generally turn out have fairly large mass, except possibly the lighter stop squark, as a result of the $\alpha_s$ terms in the RGE’s, which are absent for the sleptons.) Experimental searches at the Tevatron should pay more attention to the types of discovery modes that we shall delineate in the following sections.

3. Model Parameter Space Constraints

The eight different models that we shall explore were delineated in the Introduction. They will be denoted by $D^+$ (dilaton-equivalent, $\mu > 0$), $D^-$ (dilaton-equivalent, $\mu < 0$),
Finally, we note that in evolving the \( \lambda \) threshold corrections are surely larger than those associated with these approximations. In any case, uncertainties associated with influence on the RGE equations for the strongly interacting sparticle masses and Yukawa couplings. For \( \alpha^2 \), for \( \alpha \) RGE’s bring sin2 \( \theta_W \) the lower value of \( \alpha^2 \) preferred value of 0.2324; due to its not-unlikely value for explaining the excess events at the Tevatron. We note that we do not require unification of the bottom and tau Yukawa couplings at \( M_U \). Typically, their ratio at \( M_U, R_{b/\tau} \), is within 15% of unity in the models being considered. Although it would greatly simplify our considerations by reducing the parameter space to just a single dimension, requiring absolutely precise Yukawa unification may well be an artificially strong constraint.

Various choices for the two model-determining parameters can be considered. We have found it convenient to employ \( m_{\tilde{g}} \) and tan \( \beta \), where \( m_{\tilde{g}} \equiv m_{\tilde{g}}(\text{pole}) \) is the mass of the physical gluino state, and tan \( \beta \) is the ratio of the Higgs doublet field vacuum expectation values. Of course, in employing these two low-energy parameters we are using a bottom-up approach to the renormalization group equations, as pioneered in Ref. [20]. The RGE’s are solved by employing one-loop evolution equations, adopting \( \alpha_s(m_Z) = 0.12, \alpha_{QED} = 1/127.9, m_b(m_b) = 4.25 \text{GeV} \) and \( m_t(m_t) = 170 \text{GeV} \). Evolution is performed between \( M_U, \alpha_U \) and \( m_{\tilde{g}} \). The resulting values of \( M_U \) and \( \alpha_U \) are \( M_U = 2.39 \times 10^{16} \text{GeV}, \alpha_U = 0.0413 \) for the \( D, MS \), models and \( M_U = 1 \times 10^{18} \text{GeV}, \alpha_U = 0.0551 \) for the \( SD \) and \( SMS \) models. At one-loop, this procedure predicts sin2 \( \theta_W \) = 0.2305, i.e. outside the errors on the experimentally preferred value of 0.2324, however, it is well-known that the two-loop corrections to the RGE’s bring sin2 \( \theta_W \) into much closer agreement (sin2 \( \theta_W \sim 0.2335 \)) with the experimental result for the chosen value of \( \alpha_s(m_Z) = 0.12 \). At one-loop, we prefer to employ this value for \( \alpha_s(m_Z) \), that is more in the center of its experimentally allowed range, as opposed to the lower value of \( \alpha_s(m_Z) = .11 \) that at one-loop would yield sin2 \( \theta_W \) = 0.2324, due to its influence on the RGE equations for the strongly interacting sparticle masses and Yukawa couplings. In any case, uncertainties associated with \( M_U \) or \( M_S \) boundary conditions and threshold corrections are surely larger than those associated with these approximations. Finally, we note that in evolving the 'L' parameters entering the Higgs potential the various sparticles are decoupled at their respective mass scales; this is accomplished by an iterative method.

The first step in our analysis is to determine the allowed region of the \( m_{\tilde{g}} \)-tan \( \beta \) parameter space for each of the models. The constraints that we shall apply are the following, not all of which turn out to be important.

1. We require a neutral LSP. This requirement determines the upper limit on the mass of the gluino at fixed tan \( \beta \) that arises in the \( MS^+ \) and \( MS^- \) models; if \( m_{\tilde{g}} \) becomes too large, \( m_{\tilde{\chi}_1^0} \) exceeds \( m_{\tilde{\tau}_1} \).

2. We require that the \( h^0 \) and \( A^0 \) of the model not be visible at LEP. This is not constraining, due to the large value of \( m_t(m_t) = 170 \text{GeV} \) that we employ.
3. We require that all sleptons be heavier than $m_Z/2$, since sleptons are not observed in $Z$ decays.\cite{23} Requiring $m_\tilde{\nu} > m_Z/2$ determines the lower $m_\tilde{g}$ boundary for the $SMS^-$ and $D^-$ models at all $\tan \beta$ values, and for the $MS^+$, $SMS^+$ and $SD^-$ models for $\tan \beta \lesssim 2.3$, $4.5$, and $10.5$, respectively. It also determines the upper bound on $\tan \beta$ in the $D$, $SD$ and $SMS$ models — at large $\tan \beta$ splitting between the $\tilde{\tau}$'s becomes sufficiently large that the $\tilde{\tau}_1$ is pushed to a mass below $m_Z/2$.

4. We require that the lightest chargino, the $\tilde{\chi}_1^+$, be heavier than $m_Z/2$, since chargino pairs are not observed in $Z$ decays.\cite{23} This requirement determines the lower $m_\tilde{g}$ bound for the $D^+$ and $SD^+$ models, and for the $MS^+$, $SMS^+$ and $SD^-$ models for $\tan \beta \gtrsim 2.3$, $4.5$, and $10.5$, respectively.

5. We require that $m_{\tilde{t}_1} > m_Z/2$, where $\tilde{t}_1$ is the lighter of the (rather widely split) stop squark mass eigenstates.\cite{23} This requirement is not constraining in the cases studied.

6. We require $m_\tilde{q} > 100$ GeV, for all squarks other than the $\tilde{t}$. This is only a rough lower bound from CDF/D0 data\cite{24,25} but is, in any case, not constraining for the models we study.

7. Similarly, we note that a CDF/D0-like requirement of $m_\tilde{g} > 120$ GeV is not constraining. In fact, the slepton, chargino, and LSP boundary conditions require that the minimum value of $m_\tilde{g}$ is always somewhat above 200 GeV in the models considered.

8. We require that the net contribution from new states to the $Z$ width be smaller than 0.028 GeV, and that any additional contribution to the $Z$'s invisible decay width be $< 0.018$ GeV.\cite{23} Neither requirement is constraining for the models considered.

9. We also demand that the EWSB potential minimum for any acceptable solution to the RGE equations be a true global minimum and that the potential be bounded at $M_U$ or $M_S$.

10. We require that the top quark Yukawa coupling remain perturbative as defined by $h_t \leq 3$. For Yukawa coupling larger than this the two-loop corrections to the one-loop renormalization group equations become large and the perturbative approach begins to break down. This requirement determines the lower boundary, $i.e.$ smallest allowed value of $\tan \beta$, for all the models.

A possible further constraint on the models, that we shall not directly implement, derives from the fact that the $\chi_1^0$ can provide\cite{26} a significant dark-matter density in the early universe. As is well-known,\cite{26} if the $\chi_1^0$ becomes too heavy, and if the cross section for the annihilation of $\chi_1^0$ pairs is not large, then the universe can be overclosed or ‘too young’. However, in the models we consider the sleptons and sneutrinos are generally rather light, which tends to enhance the annihilation cross sections. Thus, we expect much of the parameter space illustrated to remain allowed by even a rather stringent constraint on dark matter relic density. This is illustrated for example in the investigations of Ref. [3]. At most, large $m_\tilde{g}$ values, $i.e.$ beyond those relevant for Tevatron searches, would be eliminated by imposing this constraint.

We also do not implement proton decay constraints. While full gauge group unification
at $M_U$ can lead to difficulties with proton decay, such full unification is not typical of string theories. In string theories, we have noted earlier that coupling constant equality at $M_U$ is instead a result of the simplest and most attractive choices for the Kac-Moody levels, $k_i$. Indeed, many string models with only the minimal SM group structure have been constructed \[28\] In the absence of full gauge-group unification, the $X$ and $Y$ gauge bosons and the especially troublesome Higgs triplets need not be present with the result that there is no definitive constraint on the models coming from proton decay. Of course, the $SD$ and $SMS$ scenarios fit nicely into the flipped $SU(5) \times U(1)$ model, for which there need not be a problem with proton decay (because of the large scale at which unification takes place) despite the fact that true gauge-group unification occurs.

The allowed regions of $m_\tilde{g}$–$\tan \beta$ parameter space for the eight models obtained by imposing these constraints are displayed in Fig. 1. Note that these constraints alone do not serve to determine a right-hand boundary in the $D$, $SD$ or $SMS$ scenarios. Presumably, it would be unreasonable to consider solutions with $m_\tilde{g} > 1$ TeV purely on the aesthetic ground that such a large gluino mass would bring into question the original naturalness motivation for the MSSM. We have confined ourselves to $m_\tilde{g} < 800$ GeV simply because of our focus on the Tevatron in this paper. As noted earlier, we have not imposed unification of the $b$ and $\tau$ Yukawa couplings at $M_U$. However, a choice for $\tan \beta$ (along with $m_t(m_t) = 170$ GeV and $m_\tilde{u}(m_b) = 4.25$) determines the ratio, $R_{b/\tau}(M_U)$, independently of $m_\tilde{g}$. This ratio is given for each scenario along the right hand axes in Fig. 1.

It is also of interest to outline the values taken on by $B^0$ in the allowed regions of Fig. 1 for each of the eight models. For the $D^+$ model, $B^0 < 0$ throughout the allowed region of Fig. 1a. Adopting the approximations of Eqs. (3) and (4), this would exclude a Calabi-Yau model with only $B^0_Z$. However, $B^0/m_3/2 = -2$ and $-(1 + \sqrt{3})$ both fall within the allowed parameter space (at $\tan \beta \sim 4$ and 2.7, respectively — note that $B^0$ becomes less negative as $\tan \beta$ increases). This means that a $B^0_{\mu}$ source would be entirely consistent for all but $\theta$ very near 0 (for which our boundary conditions are not appropriate in any case). The results for the $SD^+$ model are essentially identical. For the $D^-$ and $SD^-$ models, $B^0 > 0$ everywhere in the allowed parameter space depicted in Fig. 1a and b. This requires the presence of a $B^0_Z$ mixing source; in other words, orbifold compactifications are excluded in the context of the approximations of Eqs. (3) and (4). In the $D^-$ and $SD^-$ models $B^0$ increases as $\tan \beta$ decreases; however, the $\theta \sim 0$ limit of $B^0_Z/m_3/2 \sim 2$ falls at tan $\beta$ values below 2, which are excluded by non-perturbative top-quark Yukawa coupling behavior for $m_t = 170$ GeV. (For lower $m_t$ values this limit is reached within the allowed domain.) Turning to the $MS$ and $SMS$ models, the basic features of $B^0$ are easily summarized. For $MS^+$ and $SMS^+$, $B^0 < 0$ throughout the allowed parameter spaces shown in Figs. 1c and d. $B^0$ becomes less negative as $\tan \beta$ increases, but never reaches 0, i.e. the value consistent with $B^0 = A^0 - \mu^0$ for our choice of $A^0 = m^0 = 0$. For the $MS^-$ or $SMS^-$ models, $B^0$ goes from positive to negative values as $\tan \beta$ increases, passing through zero at $\tan \beta \sim 7 - 9$ or $\tan \beta \sim 5 - 9$, respectively (larger $\tan \beta$ for larger $m_\tilde{g}$). Thus, the strict no-scale boundary conditions of $m^0 = A^0 = B^0 = 0$ are only possible for $\mu < 0$. As stated earlier, we shall not place any

\* For example, in $SU(5)$ unification proton decay is frequently too rapid.\[27\]
restriction on $B^0$ in our phenomenological analyses. The above $B^0$ results are presently only of passing theoretical interest, but could become useful should predictions for $B^0$ become more certain at some future date.

The phenomenology of the different models is largely determined by the masses of the super particles. Thus, before turning to the specific scenarios that we shall examine with regard to detection at the Tevatron or at LEP-II, it is useful to illustrate the basic structure for the eigenstate masses that emerge from the four types of models being considered. Since, for a given model, the only mass scale is $m_\tilde{g}$, it is not surprising that all masses when plotted in ratio to $m_\tilde{g}$ exhibit approximate scaling. Only the variation with tan $\beta$ (through the limited range allowed by the parameter space boundaries) yields any scatter. These scaling laws are illustrated in Fig. 2, where the $m_i/m_\tilde{g}$ values exhibit a well-defined band for a given choice of sparticle type, $i$. Note that at large $m_\tilde{g}$ the $m_i/m_\tilde{g}$ ratio for a given sparticle is essentially independent of tan $\beta$. The gaugino masses exhibit the standard relations $^{29}$ $m_{\tilde{\chi}^0_1} \sim M'$ and $m_{\tilde{\chi}^\pm_1} \sim M$ that arise whenever $|\mu| > M', M$, where $M'$ and $M$ are the $U(1)$ and $SU(2)$ low-energy soft masses respectively. ($m_{\tilde{\chi}^0_2}$ is not plotted; as expected in the above limit, $m_{\tilde{\chi}^0_2}$ is always very close to $m_{\tilde{\chi}^+_1}$.) These mass limits are correlated with the fact that the $\tilde{\chi}^0_1$ and $\tilde{\chi}^0_2$ become primarily bino and wino in the large $|\mu|$ limit, a fact which we shall see has some rather important phenomenological consequences.

The crucial role of the mass hierarchies illustrated in Fig. 2 is in determining the production rates and decay chains that are the dominant ingredients in the phenomenological consequences of a particular choice for the model, and the values of tan $\beta$ and $m_\tilde{g}$ within the given model. There are important similarities as well as important differences between the models in this respect. As noted in the introduction, a very important point to note is that since the soft scalar masses $m_i$ are either zero at $M_U$ (for the minimal-supergravity scenarios) or at least smaller than the common gaugino mass by a factor of $\sqrt{3}$ (for the dilaton-like scenarios), the scalar partners of the SM fermions acquire mass largely as the result of evolution from the unification scale $M_U$ down to $\sim m_Z$. Since the evolution of the sleptons is much slower than that of the squarks (there being no strong interaction terms driving them away from zero mass), sleptons will always be very much lighter than squarks in these scenarios. Even the squarks only reach masses as large as the gluino mass for the higher unification scale at $M_S$. In the case of the GUT-scale-unified minimal-supergravity models, the sleptons are quite light, and indeed not much heavier than the $\tilde{\chi}^0_1$. This is why limits on the slepton mass set the upper and lower boundary on $m_\tilde{g}$ in the $MS$ models. In the $SMS$ models, the sleptons and squarks move up in mass as a result of the increased amount of evolution arising from the larger difference between $M_S$ and $m_Z$. For the $D$ models, slepton masses are again larger than in the $MS$ models, but now as a result of the non-zero seed scalar masses ($m_i \neq 0$) at $M_U$. Indeed, the slepton masses are not very different in the $D^\pm$ and $SMS^\pm$ scenarios.

However, there are subtleties in comparing slepton to ino masses that have considerable phenomenological importance. For $m_\tilde{g} \lesssim 500$ GeV, the region of interest in our Tevatron study, the mass hierarchies for $D^+$ and $SMS^+$ are such that $\tilde{\chi}^0_2 \to \nu \tilde{\nu}$ and $\tilde{\chi}^+_1 \to l \tilde{\nu}$ decays are generally forbidden, whereas for $D^-$ and $SMS^-$ these decays are generally allowed. We shall see that this results in significant similarities between the $SMS$ and $D$ phenomenologies for both signs of $\mu$. The $SD^+$ and $SD^-$ hierarchies are such that the lightest chargino is
generally lighter than the lightest slepton, the $\tilde{\nu}$ (except at the lowest allowed $m_{\tilde{g}}$ values). As a result $\tilde{\chi}_0^2 \rightarrow \nu \tilde{\nu}$ (recall that $m_{\tilde{\chi}_0^2} \simeq m_{\tilde{\chi}_1^+}$) and $\tilde{\chi}_1^+ \rightarrow l \tilde{\nu}$ are again forbidden (except for quite low $m_{\tilde{g}}$ and $\tan \beta \lesssim 10.5$ in the $SD^-$ case). This will imply some similarity of $SD^+$ and $SD^-$ phenomenology to that of the $D^+$ model.

In contrast, since the squark masses in all the models considered derive mainly from evolution, the squark masses in the $D$ and $MS$ models are of similar size and somewhat smaller than those predicted in the $SD$ and $SMS$ cases.

In summary, we see that the gluino mass is the largest mass of all sparticles in these scenarios, and hence may not be particularly relevant for phenomenology, contrary to typical expectations. Squark masses are typically $\sim 10 - 50$ GeV below gluino masses. Sleptons are amongst the lightest SUSY particles; hence the charginos and neutralinos often decay via two-body modes to e.g. slepton-lepton, instead of the usually expected three-body decays. We shall see that this can have a substantial influence on the types of collider signatures expected. In addition, we will find that the sleptons are generally sufficiently light that a significant portion of parameter space can be probed via slepton-pair production at LEP-II.

4. Scenarios

The allowed regions in $m_{\tilde{g}} - \tan \beta$ parameter space for the minimal-supergravity and dilaton-equivalent scenarios (for both $\mu > 0$ and $\mu < 0$) are illustrated in Fig. 1. In order to explore the ability of the Tevatron to probe these eight distinct scenarios, we have selected a series of points in each of the eight allowed parameter spaces that comprise a representative sample of cases that might also have some chance of being accessible to the Tevatron. The sampled points are numerically labelled in Fig. 1, and will be referenced by $D_i^+, D_i^-, SD_i^+, SD_i^-, MS_i^+, MS_i^-, SMS_i^+, SMS_i^-$, where $i$ is the numerical index indicating the sampled point within a given scenario. The locations of all these scenarios in the eight parameter spaces are shown in Fig. 1 using the numerical label for a given scenario. In the $SD$ and $SMS$ scenarios, we have focused primarily on points that lie near the limit of experimental sensitivity at the Tevatron.

The phenomenological consequences of a given point within one of the eight scenarios is largely determined by the masses and decay branching ratios of the superpartner sparticles. In Tables 1a–d (for $D$, $SD$, $MS$, $SMS$ models, respectively) we give the scenario label and $m_{\tilde{g}}$, $\tan \beta$ values for a given scenario point, along with the masses of the lightest two neutralinos, $\tilde{\chi}_1^0, \tilde{\chi}_2^0$, the lightest chargino, $\tilde{\chi}_1^+$, the left-handed slepton, $\tilde{l}_L$, the right-handed slepton, $\tilde{l}_R$, the sneutrino, $\tilde{\nu}$, the first and second family squarks, denoted generically by $\tilde{q}$, and the lighter of the two stop eigenstates as obtained after diagonalization, $\tilde{t}_1$. The Higgs masses $m_{h^0}$ and $m_{A^0}$ are also tabulated. Regarding the $h^0$, our procedure is to compute $m_{h^0}$ using the one-loop effective potential corrections (including stop and sbottom mass splitting effects) in the manner of Ref. [30] after having carried out the RGE evolution in the manner described earlier. The $h^0$ masses obtained in this way are somewhat higher (slightly lower) in the $\mu > 0$ ($\mu < 0$) cases than those emerging directly from the RGE’s.

Some important branching ratios for each scenario point appear in Tables 2a–d. In the $D$, $SD$ and $SMS$ scenarios the two-body decays for $\tilde{\chi}_1^+$ and $\tilde{\chi}_2^0$ into $\tilde{l}_L$ are not allowed or are completely negligible, so we do not show these. Also, in the minimal-supergravity scenarios,
the $\tilde{\nu}$ always decays invisibly to $\chi_1^0\nu$; thus, branching ratios for the $\tilde{\nu}$ are not tabulated for the $MS$ case. In the $D^+$ and $SD^\pm$ scenarios, and in one case each for $D^-$ and $SMS^\pm$, there are significant visible decays of the $\tilde{\nu}$, as indicated. Finally, the $D_9^-$ case is rather special in that the two-body decays $\tilde{\chi}_1^+ \rightarrow \tilde{\chi}_1^0W^+$ and $\tilde{\chi}_2^0 \rightarrow \tilde{\chi}_1^0h^0$ (where $h^0$ is the light CP-even Higgs boson) are allowed and dominant (though not listed).

The most crucial distinction amongst models however, derives from the fact that $\tilde{\chi}_1^+$ and $\tilde{\chi}_2^0$ decays tend to be saturated by the modes $\tilde{\chi}_1^+ \rightarrow \tilde{\nu}l$ and $\tilde{\chi}_2^0 \rightarrow \tilde{\nu}\nu$ when kinematically allowed; this is quite characteristic of the $MS^+$, $MS^-$, $D^-$ and $SMS^-$ scenarios. The dominance of the $\tilde{\nu}\nu$ channel in the case of the $\tilde{\chi}_2^0$ occurs despite the fact that the $\tilde{\nu}_R$ channel generally has at least as much kinematic phase space. This is because the $\tilde{\nu}_R$ couples only to the bino component of the $\tilde{\chi}_2^0$, which is quite small whenever $|\mu| > M$, as noted earlier. Indeed, it should be noted that the $\tilde{\chi}_2^0$ approaches a pure wino state much more rapidly for large values of $\mu < 0$ than for $\mu > 0$. Thus, the $\tilde{\nu}_R$ decays of the $\tilde{\chi}_2^0$ are particularly suppressed for the $\mu < 0$ cases. The $\tilde{\nu}\nu$ channel dominance is important phenomenologically since the $\tilde{\nu}$ decays entirely to the invisible $\chi_1^0\nu$ channel whenever $\tilde{\chi}_1^+ \rightarrow \tilde{\nu}l$ and $\tilde{\chi}_2^0 \rightarrow \tilde{\nu}\nu$ decays are allowed. The result is a depletion of the visible event rate from $\tilde{\chi}_2^0$ decays, especially for the $\mu < 0$ scenarios where the $\tilde{\nu}_R$ branching ratio is particularly suppressed.

The net effect of these branching ratios on signatures is difficult to deduce without a complete simulation. For instance, in the search for trilepton events from $pp \rightarrow \tilde{\chi}_1^0\tilde{\chi}_2^0 \rightarrow 3l + E_T$, the chargino and neutralino branching fractions can vary considerably, leading to wide ranges of signal rates. In addition, kinematic effects can be important. For instance, if $\tilde{\chi}_2^0 \rightarrow \tilde{l} + l$, the final state $l$ may be too soft to pass detector requirements, even if the branching ratio is large.

We shall see that the $D^-, SMS^-$ scenarios are distinctly more difficult to detect than the $D^+, SMS^+$ ones. However, the $MS^-$ scenarios will turn out to be as easily probed as the $MS^+$ scenarios. This is because of the ‘inverted’ mass hierarchy, $m_{\tilde{\chi}_1^+} > m_{\tilde{\chi}_1^-}$ for large $m_{\tilde{\chi}_1^-}$ in $MS$ models compared to $D$ and $SMS$ models. The inversion allows for $\tilde{\chi}_2^0 \rightarrow \tilde{l}_Ll$ decays (for which the coupling does not go to zero as the $\tilde{\chi}_2^0$ approaches a pure wino state, although it is small). The resulting production and decay chain $\tilde{\chi}_1^0\tilde{\chi}_2^0 \rightarrow \tilde{l}_Ll + \tilde{l}_Ll$, followed by two $\tilde{l}_L \rightarrow l\tilde{\chi}_1^0$ decays, yields 3l states at a significant rate. Meanwhile, the similarity between the $SD^\pm$ and the $D^+$ mass hierarchies and decays implies that the $SD^\pm$ models will be more or less as easily probed as the $D^+$ case.

Before closing this section, we note that some of the above discussion is peculiar to the $m_{\tilde{\chi}_1^-}$ mass range relevant to Tevatron exploration. In particular, for high enough $m_{\tilde{\chi}_1^-}$ values the $\tilde{\nu}$ becomes heavier than the $\tilde{\chi}_1^\pm$ for the $D^-$ and $SMS^-$ cases, and thus would decay visibly, exactly as for the $D^+$ and $SD^\pm$ model cases explored here. Correspondingly, the $\tilde{\chi}_2^0 \rightarrow \tilde{\nu}\nu$ and $\tilde{\chi}_1^+ \rightarrow \tilde{l}\tilde{l}$ two-body decays become kinematically disallowed. Three-body decay channels would play a more prominent role, as in the $D^+$ and $SD^\pm$ cases. The final leptons would generally be harder as a result (although perhaps less numerous). Overall, there could be a temporary increase in the tri-lepton rate after cuts (as $m_{\tilde{\chi}_1^-}$ is increased). (This is, in fact, the
source of the greater observability that we shall find for the $D^+$ and $SD^\pm$ as compared to the $D^-$ and $SMS^-$ models.) This illustrates how phenomenological considerations could well change significantly in moving to either higher luminosity or higher energy at the Tevatron. And certainly discovery potential at the LHC cannot be extrapolated from the results we shall present here. In general, the phenomenological complexity of the types of models considered is substantial because of the delicate cross-over’s in masses and decay modes.

5. Simulation and Selection Cuts

To simulate signal and background events at the Tevatron collider, we use the event generator program ISAJET 7.07. ISAJET 7.07 has been set up to perform a reasonable simulation of supergravity models, provided $\tan \beta < \sim 10$. For larger $\tan \beta$ values, the approximate degeneracy of sleptons is badly broken, and there are large mixing effects for $\tilde{b}$ and $\tilde{\tau}$ states. For two simulations at large $\tan \beta$ — the $D^+_3$ and $D^-_5$ cases — two-body decays of charginos and neutralino to real staus are dominant: in these cases we force the relevant decays to occur with 100% branching.

Briefly, for a given set of weak scale MSSM parameters, ISAJET calculates branching fractions for all sparticle decay modes. ISAJET then generates all SUSY particle production processes according to their relative cross sections, and decays the various SUSY and SM particles via the calculated or measured branching ratios. Initial and final state QCD radiation is included, as is hadronization of quarks and gluons. Underlying event soft-scattering is included as well.

The relative production cross sections for a variety of superparticle-pair channels are given in Tables 3a–d for all of the numbered scenario cases appearing in Fig. 1 and in Tables 1 and 2.

We see immediately that $\tilde{g}\tilde{g}$, $\tilde{g}\tilde{q}$ and $\tilde{q}\tilde{q}$ production processes often have low rates compared to other SUSY-pair production processes, especially for large values of $m_{\tilde{g}}$. Instead, the dominant production processes are the cumulative $\tilde{\chi}\tilde{\chi}$ subprocesses, especially $\tilde{\chi}^+_1\tilde{\chi}^0_2$ and $\tilde{\chi}^+_1\tilde{\chi}^-_1$. Furthermore, the light slepton and sneutrino masses characteristic of these models results in a substantial rate for $\tilde{\ell}\tilde{\nu}$ production. Finally, there is significant rate for the associated production final states $\tilde{\chi}\tilde{g}$ and $\tilde{\chi}\tilde{q}$, which comprises the remainder of the event rate.

To model collider detector effects, we employ the toy calorimeter simulation ISAPLT. We assume calorimeter extends over the central region out to rapidity of $|\eta| < 4$, with cell size $\Delta \eta \times \Delta \phi = 0.1 \times 0.1$. We take hadronic energy resolution to be $70%/\sqrt{E_T}$, and electromagnetic resolution to be $15%/\sqrt{E_T}$. Jets are coalesced within cones of $R = \sqrt{\Delta \eta^2 + \Delta \phi^2} = 0.7$, using the ISAJET routine GETJET; clusters with $p_T > 15$ GeV are labelled as jets. Muons and electrons are classified as isolated if they have $p_T(l) > 8$ GeV, $|\eta(l)| < 3$, and there is less than $p_T(l)/4$ GeV in a cone of $R = 0.4$ about the lepton direction. For all supersymmetric event topologies, we require at least $E_T > 20$ GeV.

Multi-lepton signals for gluinos and squarks have been examined in Ref. [33]. We follow approximately the cuts given there. The cross sections examined include the following.

1. Multi-jet plus missing energy ($E_T$):
   - $E_T > 50$ GeV,
• no isolated leptons with $p_T(l) > 15$ GeV,
• number of jets $n(jet) \geq 4$,
• at least one central jet ($|\eta(jet)| < 1$) and no jet within $30^\circ$ of $\not{E}_T$.

This is a generic $\not{E}_T$ cut, i.e. it is not optimized for various gluino masses. For instance, for very massive gluinos, a substantially larger $\not{E}_T$ cut may be desirable to improve signal to background rate.

2. Single charged lepton (1l):
• exactly one isolated lepton with $p_T(l) > 15$ GeV,
• veto events with $60 < M_T(l, \not{E}_T) < 100$ GeV, to reduce real $W$ background.

3. A pair of oppositely charged same-flavor leptons (OS):
• two isolated leptons with $p_T(l) > 15$ GeV,
• $30^\circ < \Delta \phi(l^+l^-) < 150^\circ$,
• no jets, and
• veto events with $80 < M(l^+, l^-) < 100$ GeV, to reduce real $Z$ background.

These cuts are designed to extract possible signals from slepton-pair production.\[^{34}\]

4. A pair of same-sign leptons (SS):
• two same-charge isolated leptons with $p_T(l) > 15$ GeV.

This is designed to extract gluino-pair cascade decay events, by exploiting the Majorana nature of the gluino.\[^{35}\]

5. Three leptons (3l):
• three isolated leptons with $p_T(l_1) > 15$ GeV, $p_T(l_2) > 10$ GeV and $p_T(l_3) > 8$ GeV,
• veto events with $80 < M(l^+, l^-) < 100$ GeV, to reduce real $WZ$ background,
• (optional requirement of zero or one jet, or all jets accompanying the event.)

These cuts are designed to extract either gluino and squark cascade decay events \[^{33}\] (with jets), or to extract clean trileptons from $\tilde{\chi}^\pm_1 \tilde{\chi}^0_2$ production \[^{36}\] (with zero or one jet).

6. Four leptons (4l).\[^{33}\]
• four isolated leptons with the first three leptons as in (5.) above, while in addition $p_T(l_4) > 8$ GeV.

In addition, on occasion we picked up events containing five isolated leptons. We do not list these relatively rare event cross sections due to the considerable statistical uncertainty.
6. Numerical Results for Signal and Background

The background cross section levels in the various channels after cuts are given in Table 4:
2720, 1.1 × 10^6, 32, 2, 0.3 (0.7), and ∼ 0 in units of fb for the \( E_T, 1l, OS, SS, 3l \) and 4l channels, respectively. The backgrounds that were included are: \( W + jets, Z + jets, t\bar{t} \) (\( m_t = 170 \text{ GeV} \)), \( W^+W^- \), and \( W^\pm Z \). The quoted rates include the \( \tau \) mode decays of the \( W^\pm \) and \( Z \). The signal rates for each numbered case of Fig. 1 are given in Tables 5a–d, for the \( D, SD, MS \) and \( SMS \) models, respectively.

The \( E_T + jets \) cross section after the above cuts is plotted in Fig. 3 versus \( m_{\tilde{g}} \) for the various scenarios. \( E_T + jets \) events arise from many different sources typically, including \( \tilde{\chi}_1^\pm \tilde{\chi}_1^0, g\tilde{g}, q\tilde{q}, t_1\bar{t}_1, \tilde{g}\tilde{\chi}_1, \) and \( q\tilde{\chi}_1 \) events, where all final state leptons are missed or soft. All signals are below our calculated background of 2720 fb. However, assuming that the background can be normalized by independent measurements and Monte Carlo studies, a 5\( \sigma \) effect for signal over background (given 1 fb^{-1} of integrated luminosity) would allow a search to \( m_{\tilde{g}} \sim 300 \text{ GeV} \). This could be an overestimate given that at \( m_{\tilde{g}} \sim 300 \text{ GeV} \) the signal would only be ∼ 10% of background, implying that the latter would have to be normalized to better than 10%. However, it is also true that optimization of the \( E_T \) cut for these higher \( m_{\tilde{g}} \) values might improve the signal to background ratio somewhat. The value of measuring the \( E_T \) cross section lies in the fact that it roughly scales with \( m_{\tilde{g}} \) in spite of model differences, and different \( \tan \beta \) values. Thus, if a \( E_T + jets \) signal can be found, the size of the cross section will give an indication of the sparticle masses being probed.

The signal cross sections after cuts for the 1l sample are listed in Tables 5a–d. These signals have an enormous background from single \( W \) production, in spite of the transverse mass cut we invoke. Even with optimization of cuts (e.g. looking for events with \( M_T(l, E_T) > 100 \) GeV), detection of such a signal looks dubious, if not hopeless.

The OS dilepton sample of signal events is suited for picking out slepton-pair events. We see from Tables 5a–d that the signal in this channel can range to ∼ 80 fb, although there is a substantial background from \( WW \) production (32 fb). We find signal larger than or of order the background in several cases: \( D_1^+, D_1^-, D_2^+, MS_1^+, MS_2^+ \), and \( SMS_1^- \). These are more optimistic results than those given in Ref. [34], where slepton production was examined for more generic mass spectra. Our larger rates are in part a reflection of the very light slepton masses in the models considered here, and in part due to the fact that numerous SUSY sources other than slepton pairs contribute to the OS signal; these include mainly \( \tilde{\chi}_1^+\tilde{\chi}_1^- \) pairs, but with smaller contributions from \( \tilde{\chi}_1^0\tilde{\chi}_1^0, \tilde{g}\tilde{\chi}_1, \tilde{t}_1\bar{t}_1 \) and \( q\bar{q} \) pairs. Overall, however, observation of the OS dilepton signal looks difficult for most of the cases examined.

The same-sign isolated dilepton signal has been advocated as a means of searching for gluino-pair cascade decays, by exploiting the Majorana nature of the gluino. [35] We see from Tables 5a–d that the signal ranges from a fraction of a fb to several hundred fb, while background is at the 2 fb level, and arises (after cuts) mainly from \( WZ \) production, where one final state lepton is missed. In this case, many background events should be relatively free of jet activity, while the signal may be rich in jets if the SS events originate from strongly produced SUSY particles. We have examined the sources of the SS events for the various \( D, SD, MS \), and \( SMS \) models, and have found that they arise from a variety of SUSY production processes, including \( \tilde{\chi}_1^\pm\tilde{\chi}_1^0 \rightarrow 3l + E_T \) events, where one lepton is missed.
slepton and sneutrino production, $\tilde{g}\tilde{\chi}$ and $\tilde{q}\tilde{\chi}$ associated production events, as well as the expected $\tilde{q}\tilde{g}$, $\tilde{q}\tilde{q}$ and $\tilde{g}\tilde{g}$ events. Hence, we expect the signal events to vary substantially in topology and “jetiness”, depending on the subprocess from which they arise. Given an integrated luminosity of 1 fb$^{-1}$, we take as an estimate at least 5 such events to claim discovery. Examination of Tables 5a–d then shows that scenarios with gluino masses of up to $\sim 300 - 350$ GeV may be probed in this channel. We show in Fig. 4 the SS dilepton signal cross sections as the first entry in the brackets at each of the numbered scenario points appearing in Fig. 1. (Except in cases where overlap forced some slight repositioning, the $m_{\tilde{g}}$–$\tan\beta$ values for a particular point correspond to the location of the lower left-hand corner of the bracket.)

Another promising event topology for the discovery of SUSY is events with three isolated leptons. These events can arise from $\tilde{\chi}_1^{\pm}\tilde{\chi}_2^0 \rightarrow 3l + E_T$ production, in which case they will be relatively free of extra jet activity, or they can arise, for instance, from gluino and squark production processes, where the cascade decays result in leptonically decaying $\tilde{\chi}_1^{\pm}$ and $\tilde{\chi}_2^0$ states, in which case the trileptons will be accompanied by substantial jet activity. We show the cross section for $3l$ events in Tables 5a–d for events containing just 0 or 1 jet, or any number of jets (in parenthesis). These cross sections range from a fraction of a fb, to up to 168 fb, for the cases examined. The background, listed in Table 4, is 0.34 (0.69) fb. Again, assuming that at least five events are needed for discovery (in 1 fb$^{-1}$ of data), we see that cases with gluino mass beyond 500 GeV may be probed in the $MS^+$, $MS^-$, $SMS^+$, $D^+$, $SD^+$ and $SD^-$ models, while in the $D^-$ and $SMS^-$ models discovery reach is restricted to $m_{\tilde{g}} \lesssim 300$ GeV. We list the trilepton rates (for events with all jet multiplicities) in Fig. 4, as the second entry in the bracketed figures.

By combining the discovery potential of both the $SS$ and $3l$ signals, we have estimated the region explorable by Tevatron collider experiments with $L = 1$ fb$^{-1}$ of integrated luminosity. The boundary of this region for each model is drawn on Fig. 4 as the dashed line. It is remarkable that so large a fraction of the model parameter spaces will yield observable rates for these two new-physics signatures. The most difficult scenarios to detect are those associated with the $D^-$ and $SMS^-$ models. This is clearly a result of the suppressed $\tilde{\chi}_1^{\pm}\tilde{\chi}_2^0 \rightarrow 3l, SS$ mode (where the SS events arise when one of the $l$’s is missed), as discussed earlier. In addition, in $\tilde{q}\tilde{g}$ etc. events $\tilde{\chi}_1^{\pm}\tilde{\chi}_1^\pm \rightarrow ll\tilde{\nu}\tilde{\nu}$ yields very soft leptons due to the generally very small $m_{\tilde{\chi}_1^\pm} - m_{\tilde{\nu}}$ mass difference, which is much smaller for these scenarios than any others. As noted earlier, the difficulty of probing the $D^-$ and $SMS^-$ scenarios is to be contrasted with the situation for the $MS^-$ models, for which the inverted mass hierarchy ($m_{\tilde{\chi}_1^\pm} > m_{\tilde{\ell}_L}$ at large $m_{\tilde{g}}$) allows for much larger $3l$ and SS rates. Indeed, the $MS^-$ parameter space can be almost fully explored.

Finally, we list in Tables 5a–d the $4l$ event rates. In a few cases, the rate for $4l + E_T$ events can range up to the 30–60 fb level, giving again a spectacular signature for SUSY. These events usually occur due to events containing a $\tilde{\chi}_2^0$ pair, either from cascade decays, or from direct production, followed by $\tilde{\chi}_2^0 \rightarrow l\tilde{\chi}_1^0$ decay. We were unable to generate any substantial background to this process. We note, however, that it does not occur at a large rate for most of the scenarios considered, and hence would not constitute an optimal discovery channel.
As noted earlier, extrapolation of these results to higher luminosity is somewhat dangerous, but we allow ourselves a few very approximate statements based on the 3l mode which has the lowest background rate. For an integrated luminosity of $L = 30 \text{ fb}^{-1}$ (three years running at 10 fb$^{-1}$ per year as proposed in some versions of a future Tevatron upgrade$^{[37]}$, the 3l background rate (assuming no additional sources of background become important at high luminosity, e.g. from multiple interactions per crossing) is obtained from Table 4; we find a background rate of 21 events for the ‘all-jets’ case. An examination of the signal rates in Tables 5a–d shows that many of the numbered scenarios that are unobservable for $L = 1 \text{ fb}^{-1}$ would then become observable. The increase in parameter space coverage would be especially dramatic for the $D$ and SMS$^-$ models. Adopting a $5\sigma$ criterion (i.e. 23 or more signal events), $m_{\tilde{g}}$ values as high as $\sim 500 \text{ GeV}$ would yield detectable signals for the $D$ and SMS$^-$ models. For the other models, we have not studied scenarios with high enough $m_{\tilde{g}}$ to establish a meaningful estimate of how much higher in $m_{\tilde{g}}$ one can go with $L = 30 \text{ fb}^{-1}$. However, those scenarios we have studied suggest that discovery reach would probably be extended out to $m_{\tilde{g}} \sim 600 - 700 \text{ GeV}$.

7. Comparison of Tevatron and LEP-II

An interesting question is the extent to which LEP-II will be able to explore the $m_{\tilde{g}}$–$\tan \beta$ parameter spaces of the various models considered, and how the discovery reach of LEP-II compares to that of the Tevatron. For the models being considered the discovery reach of LEP-II is determined primarily by the $e^+e^- \rightarrow \tilde{l}_R\tilde{l}_R$, $e^+e^- \rightarrow \tilde{\chi}_1^+\tilde{\chi}_1^-$, and $e^+e^- \rightarrow Zh^0$ production processes, since it is the $\tilde{l}_R$, $\tilde{\chi}_1^\pm$ and $h^0$ that are the lightest observable particles in these models.

Let us first consider an optimistic scenario with $\sqrt{s} = 200 \text{ GeV}$ and a two-year accumulated luminosity of 1 fb$^{-1}$ (corresponding to 500 pb$^{-1}$ per year integrated over two years). Then, roughly speaking, $\tilde{l}_R\tilde{l}_R$, $\tilde{\chi}_1^+\tilde{\chi}_1^-$ production will probably be detectable for $m_{\tilde{l}_R}, m_{\tilde{\chi}_1^\pm} \lesssim 95 \text{ GeV}$, while $Zh^0$ will probably be detectable for $m_{h^0} \lesssim 105 \text{ GeV}$ (recalling that the $h^0$ has ZZ coupling that is close to full strength given the large $m_{A^0}$ values required in our scenarios). The contours for these $\tilde{l}_R$, $\tilde{\chi}_1^\pm$ and $h^0$ mass values are given in Fig. 5, for all eight of the models illustrated in Fig. 1. Specific $m_{\tilde{l}_R}$, $m_{\tilde{\chi}_1^\pm}$ and $m_{h^0}$ values for all the numbered scenarios have been listed in Tables 1a–d. As a further aid to comparing LEP-II and Tevatron results we have also indicated rough LEP-II discovery potential for the $\tilde{l}_R\tilde{l}_R$, $\tilde{\chi}_1^+\tilde{\chi}_1^-$ and $Zh^0$ channels in Fig. 4 as described below.

For $\tilde{l}_R\tilde{l}_R$-pair production, we have used an arrow in each model window to indicate the approximate upper limit in $m_{\tilde{g}}$ for which slepton-pair production will be observable at LEP-200. The width of the arrow characterizes the variation in $m_{\tilde{g}}$ associated with fixed $m_{\tilde{l}_R} = 95 \text{ GeV}$ apparent in Fig. 5. In the $MS$, $D$, $SMS$ and $SD$ models, the values of $m_{\tilde{g}}$ for which $m_{\tilde{l}_R} = 95 \text{ GeV}$ fall in the ranges $[636,665]$, $[351,370]$, $[469,493]$, $[264,279]$ GeV, respectively. $m_{\tilde{l}_R}$ and, hence, these ranges are independent of the sign of $\mu$. If only $m_{\tilde{l}_R} = 92 \text{ GeV}$ could be probed, the corresponding $m_{\tilde{g}}$ ranges become $[611,641]$, $[337,357]$, $[450,475]$ and $[253,268]$ GeV, respectively. Note that $\tilde{l}_R\tilde{l}_R$ detection will be possible for
essentially all of the $MS^-$ allowed parameter space, and for almost none of $SD^+$ parameter space.

A similar procedure is followed in the case of $\tilde{\chi}_1^+ \tilde{\chi}_1^-$-pair production. For $\sqrt{s} = 200$ GeV, we adopt $m_{\tilde{\chi}_1^+} = 95$ GeV as the discovery boundary. For a given model, the $m_{\tilde{g}}$ value corresponding to $m_{\tilde{\chi}_1^+} = 95$ GeV depends upon $\tan \beta$ as indicated in Fig. 5. Note that all the $\mu > 0$ contours are in roughly the same $m_{\tilde{g}}$ mass range, as are all the $\mu < 0$ contours. A rough summary is that the $m_{\tilde{g}}$ ranges corresponding to $m_{\tilde{\chi}_1^+} = 95$ GeV fall within the bands: $[381, 434]$ GeV for $D^+, SD^+, MS^+$ and $SMS^+$; and $[282, 385]$ GeV for $D^-, SD^-, MS^-$ and $SMS^-$. Thus, discovery of a chargino of mass $m_{\tilde{\chi}_1^+} = 95$ GeV on the average probes significantly larger $m_{\tilde{g}}$ values for $\mu > 0$ than for $\mu < 0$. The arrows labelled by $\tilde{\chi}_1^+$ in Fig. 4 reflect the above ranges. This dependence on the sign of $\mu$ is also evident in Fig. 2. There, the $m_{\tilde{\chi}_1^+}/m_{\tilde{g}}$ mass bands for $\mu > 0$ rise towards the large $m_{\tilde{g}}$ asymptotic limit as $m_{\tilde{g}}$ increases, whereas for $\mu < 0$ the bands lie above the large-$m_{\tilde{g}}$ limit, and always somewhat above the $\mu > 0$ band. Note that in the $MS^-$ model, $\tilde{\chi}_1^+ \tilde{\chi}_1^-$ pair production will be detectable for almost none of the allowed parameter space, in sharp contrast to the guaranteed discovery of $lR\tilde{R}$ pairs for this model.

In comparison to the $Zh^0$ discovery limits quoted previously, we see that $\tilde{\chi}_1^+ \tilde{\chi}_1^-$ pair production does not reach to as large $m_{\tilde{g}}$ in the $MS$ and $SMS$ cases, whereas in the $D$ and $SD$ cases $m_{\tilde{g}}$ values probed are larger (comparable) for $\mu > 0$ ($\mu < 0$). These trends become especially apparent by comparing the $m_{\tilde{l}_R}$ and $m_{\tilde{\chi}_1^\pm}$ contours in Fig. 5.

With regard to the $Zh^0$ mode at LEP-II, those numbered scenarios for which $m_{h^0} < 105$ ($> 105$) GeV are surrounded by square (rounded) brackets in Fig. 4. However, the precise boundary of detectability is very sensitive to precise luminosity, detector efficiencies and machine energy. As an indication of this, we note that if only $m_{h^0} \leq 100$ GeV could be probed, then $Zh^0$ production would not be detectable for the numbered scenarios $D_3^+, D_2^+, D_5^+, D_5^-, D_6^-, MS_4^+, MS_5^-, MS_7^-, SMS_5^-, SMS_7^-$ and $SD_2^-$ (i.e. their square brackets would become rounded brackets in Fig. 4). These scenarios, for which $m_{h^0} \gtrsim 100$ GeV, are mostly those with larger values of $\tan \beta$ — recall that $m_{h^0}$ is smallest, even after radiative corrections, for $\tan \beta$ near 1.

The $Zh^0$ detection boundaries are more clearly indicated, however, in Fig. 5. We note that the $m_{h^0} = 105$ GeV contour is absent from the $MS^-$ window of Fig. 5c, since in the $MS^-$ model $m_{h^0} < 105$ GeV for all of allowed parameter space and $Zh^0$ detection will always be possible at LEP-II (for $\sqrt{s} = 200$ GeV).

Of course, detection of the $h^0$ at the Tevatron may also be possible in the $Wh^0$ associated production channel with $h^0 \to b\bar{b}$ decay, provided there is adequate efficiency and purity for $b$ tagging. The upper $m_{h^0}$ limits for which this will be possible are luminosity dependent. For $L = 1, 10$ fb$^{-1}$ the upper limit is $m_{h^0} \lesssim 60, 75$ GeV. As can be seen in Tables 1a–d, rather few of our scenarios have masses below 75 GeV, and only one has mass below 60 GeV. Even $L = 100$ fb$^{-1}$ at the Tevatron would only allow one to probe $m_{h^0}$ masses up to $\sim 95$ GeV.

The bottom line is clear: LEP-II at $\sqrt{s} = 200$ GeV and full luminosity can generally probe much the same parameter space as can the Tevatron. The coverage is comparable for
the $D^+$ and $SMS^+$ models, somewhat greater in the case of the $D^-$, $MS^+$, $MS^-$ and $SMS^-$ models, and somewhat less in the case of the $SD^+$ and $SD^-$ models. However, generally speaking, the Tevatron is sensitive to a much broader set of SUSY particles than is LEP-II, although LEP-II does have sensitivity to both slepton pairs and chargino pairs for $m_{\tilde{g}}$ below a model-dependent value. Obviously, there is substantial complementarity between the two machines.

As already noted, these conclusions are significantly altered if LEP-II only reaches, say, $\sqrt{s} = 176$ GeV. Because the $h^0$ often has mass of the order of 100 GeV, for the chosen $m_t$ and typical $\tilde{t}$ masses predicted in the models being considered, the extent to which $Zh^0$ discovery will probe the allowed parameter spaces of Fig. 1 is extremely sensitive to the exact $\sqrt{s}$ and luminosity that will be achieved at LEP-II. For $\sqrt{s} = 176$ GeV, $Zh^0$ detection will at best only be possible for $m_{h^0} \lesssim 81$ GeV (comparable to the $L = 10$ fb$^{-1}$ Tevatron reach). The $m_{h^0} = 81$ GeV contour for each model is given in Fig. 5. Note how much less of parameter space would allow $Zh^0$ detection. Only for numbered scenarios with low $\tan \beta$ values would $Zh^0$ detection be possible, which amounts to only some 6 or 7 of the scenarios — see the $m_{h^0}$ masses in Tables 1a–d.

At $\sqrt{s} = 176$ GeV, slepton-pair production would be viable only for $m_{l_R} \lesssim 80 - 83$ GeV, which corresponds to $m_{\tilde{g}} \lesssim [504, 568], [277,316], [373,421], [210,238]$ GeV (where the ranges now include both $\tan \beta$ variation and a $80 - 83$ GeV range of possible mass accessibility) for the $MS$, $D$, $SMS$ and $SD$ models, respectively, i.e. some $70 - 100$ GeV below the $m_{\tilde{g}}$ values quoted earlier for $\sqrt{s} = 200$ GeV. Chargino-pair production would be viable for $m_{\tilde{\chi}^\pm_1} \lesssim 83$ GeV, which corresponds to the rough ranges $m_{\tilde{g}} \lesssim [346, 405]$ for $\mu > 0$ cases, and $m_{\tilde{g}} \lesssim [233,344]$ for $\mu < 0$ cases. More precise limits as a function of model and $\tan \beta$ for $\sqrt{s} = 176$ GeV are reflected by the $m_{l_R} = 83$ GeV and $m_{\tilde{\chi}^\pm_1} = 83$ GeV contours given in Fig. 5. The bottom line is clear. The slepton-pair and chargino-pair modes at LEP-II would be generally competitive with the Tevatron for the $D^-$ and $SMS^-$ models, but the Tevatron would provide a signal for SUSY over more of parameter space for the $D^+, SD^+, SD^-, MS^+, MS^-$, and $SMS^+$ models.

8. Conclusion

In this paper, we have explored the phenomenology of the gauge-coupling-unified Minimal Supersymmetric Model employing renormalization group evolution of superstring or supergravity motivated unification-scale boundary conditions (of a rather universal and attractive nature) for the soft-supersymmetry-breaking parameters. The models considered were the minimal-supergravity (or no-scale) and the dilaton-like models. At a theoretical level, the source of supersymmetry breaking and details of the Kahler potential, and so forth, are fairly different for the two model classes, and even the dilaton-like boundary conditions themselves apply for a wide variety of physics as contained in the continuous range of possible values for the goldstino angle characterizing the relative importance of moduli vs. pure dilaton supersymmetry breaking. Despite these theoretical differences, the boundary conditions are sufficiently similar (indeed, identical within the dilaton-like class) that there is a broad similarity of the basic phenomenology of these models, deriving from the presence of relatively light sleptons in all cases. In fact, we have seen that when the uncertainty associated with the question of whether unification should be required at $M_U \sim 2 \times 10^{16}$ GeV or
at $M_S \sim 10^{18}$ GeV is taken into account, the overlap between the mass spectra and resulting phenomenology of the minimal-supergravity and dilaton-like models can be quite substantial. Nonetheless, we have also seen that seemingly small shifts in mass spectra can cause substantial shifts in allowed decay modes and the consequent visibility of crucial detection channels.

Overall, the most remarkable feature of our results is the prediction (summarized in Fig. 4) that these classes of models can be probed by the existing Tevatron (with $L = 1000 \text{ pb}^{-1}$) over such a large portion of the allowed parameter spaces. In terms of the two parameters $m_{\tilde{g}}$ and $\tan \beta$, we find that even the most difficult models, namely the $\mu < 0$ dilaton-equivalent ($D^-$) and superstring-scale-unified minimal-supergravity models ($SMS^-$), yield observable tri-lepton and same-sign-lepton signals for $m_{\tilde{g}} \lesssim 300$ GeV (for all $\tan \beta$). For the $\mu > 0, D^+$ and $SMS^+$ models the $3l$ and $SS$ signals reach observable levels for $m_{\tilde{g}}$ values as high as 600 GeV at low $\tan \beta$ (as preferred if relatively precise Yukawa coupling constant unification is demanded). The $3l$ signal reaches an observable level for $m_{\tilde{g}}$ values up to 450–520 GeV in the $SD^\pm$ models. The GUT-scale-unified minimal-supergravity $\mu > 0$ ($MS^+$) model can be probed for $m_{\tilde{g}} \lesssim 420–480$ GeV. This represents somewhat more than half of the allowed parameter space given that $m_{\tilde{g}}$ has an upper bound (deriving from the requirement of a neutral LSP) in this model of about 700 GeV (or lower at high $\tan \beta$). Meanwhile, the $MS^-$ model can be probed over nearly all of the (rather restricted) parameter space.

Thus, there is cause for optimism that the scheduled main-injector upgrade of the Tevatron will reveal evidence for supersymmetry. However, there is no guarantee. Aside from the regions of parameter space for the models discussed here that lie beyond the reach of the Tevatron, there are also the (still more model dependent) moduli-dominated scenarios in which all sfermion masses are generically expected to be larger than the gluino mass. These, as discussed in Ref. [33] and Ref. [39], are more difficult to probe without a collider of significantly larger energy. Larger energies and luminosities could also improve observability of the $D, SD, MS$ and $SMS$ models. This is under investigation. Here, we note that the predicted $SS$ and $3l$ background rates imply that the discovery reach is not far from being background limited. Thus, simply increasing luminosity may not yield as large an improvement as would otherwise be the case. Increasing the energy may be more advantageous due to increased signal rates, although background rates will also go up and new backgrounds can arise.

Finally, we have noted that for some models LEP-II with $\sqrt{s} = 200$ GeV and integrated luminosity of 500 pb$^{-1} - 1 \text{ fb}^{-1}$ will be able to detect $\tilde{l}_R\tilde{l}_R$ and $\tilde{\chi}_1^+\tilde{\chi}_1^-$ pair production, and $Zh^0$ associated Higgs production, over more of $m_{\tilde{g}}$–$\tan \beta$ parameter space than that for which a SUSY signal will be seen at the Tevatron with $L = 1 \text{ fb}^{-1}$ of integrated luminosity. However, in other models LEP-200 will probe less of parameter space. Further, the relative comparison between the two machines is very dependent upon the precise energy reached by LEP-II and on whether further luminosity upgrades for the Tevatron are implemented. In general, the two machines are quite complementary, with the Tevatron being sensitive to a broader range of SUSY particle types in those regions of parameter space for which SUSY detection is possible.

**Note Added:** As we were completing this manuscript we received a paper[^40] which also
addresses the search for minimal supergravity at the Tevatron and Di-Tevatron.

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Table 1a: A tabulation of supersymmetric particle masses for the $D$ scenarios delineated in Fig. 1a.

| Scenario | $m_\tilde{g}$ | $\tan\beta$ | $m_{h^0}$ | $m_{A^0}$ | $m_{\tilde{\chi}_1^0}$ | $m_{\tilde{\chi}_2^0}$ | $m_{\tilde{\tau}_1^+}$ | $m_{\tilde{l}_L}$ | $m_{\tilde{l}_R}$ | $m_{\tilde{\nu}}$ | $m_{\tilde{q}}$ | $m_{\tilde{t}_1}$ |
|----------|----------------|--------------|-----------|-----------|----------------|----------------|----------------|----------------|----------------|----------------|-------------|--------------|
| $D_1^+$  | 282            | 2.0          | 75.3      | 240       | 23.4          | 59.1          | 47.7          | 96.8          | 75.5          | 74.4          | 252         | 158          |
| $D_2^+$  | 295            | 9.0          | 101       | 180       | 35.5          | 66.3          | 62.4          | 105           | 82.8          | 68.7          | 264         | 178          |
| $D_3^+$  | 310            | 15.0         | 103       | 180       | 39.9          | 72.5          | 70.2          | 109           | 85.9          | 74.4          | 277         | 188          |
| $D_4^+$  | 346            | 3.2          | 93.6      | 250       | 40.4          | 79.2          | 73.5          | 118           | 91.8          | 93.0          | 310         | 195          |
| $D_5^+$  | 431            | 4.5          | 104       | 300       | 58.4          | 109           | 107           | 144           | 111           | 122           | 386         | 250          |
| $D_6^+$  | 435            | 8.0          | 108       | 287       | 60.8          | 113           | 112           | 146           | 112           | 123           | 390         | 258          |
| $D_7^+$  | 503            | 5.0          | 108       | 350       | 71.3          | 134           | 133           | 166           | 127           | 147           | 450         | 297          |
| $D_8^+$  | 609            | 2.0          | 95.6      | 550       | 87.9          | 169           | 168           | 197           | 150           | 187           | 545         | 357          |
| $D_1^-$  | 232            | 2.0          | 58.4      | 190       | 37.1          | 83.5          | 83.3          | 82.3          | 65.0          | 54.1          | 207         | 215          |
| $D_2^-$  | 242            | 3.2          | 76.9      | 162       | 37.8          | 73.8          | 76.0          | 88.1          | 70.2          | 50.0          | 217         | 192          |
| $D_3^-$  | 295            | 9.0          | 98.4      | 180       | 42.6          | 76.5          | 77.4          | 105           | 82.8          | 68.7          | 264         | 195          |
| $D_4^-$  | 301            | 2.2          | 69.0      | 244       | 47.3          | 100           | 100           | 103           | 80.2          | 79.8          | 269         | 242          |
| $D_5^-$  | 310            | 15.0         | 101       | 180       | 43.9          | 78.7          | 79.0          | 109           | 60.3          | 74.4          | 277         | 198          |
| $D_6^-$  | 346            | 3.2          | 86.0      | 250       | 53.4          | 106           | 106           | 118           | 91.8          | 93.0          | 310         | 246          |
| $D_7^-$  | 431            | 4.5          | 98.4      | 300       | 65.6          | 128           | 128           | 144           | 111           | 122           | 386         | 285          |
| $D_8^-$  | 503            | 5.0          | 103       | 350       | 76.9          | 150           | 151           | 166           | 127           | 147           | 450         | 329          |
| $D_9^-$  | 609            | 2.0          | 81.3      | 550       | 95.4          | 193           | 193           | 197           | 150           | 187           | 545         | 436          |

Table 1b: A tabulation of supersymmetric particle masses for the $SD$ scenarios delineated in Fig. 1b.

| Scenario | $m_\tilde{g}$ | $\tan\beta$ | $m_{h^0}$ | $m_{A^0}$ | $m_{\tilde{\chi}_1^0}$ | $m_{\tilde{\chi}_2^0}$ | $m_{\tilde{\tau}_1^+}$ | $m_{\tilde{l}_L}$ | $m_{\tilde{l}_R}$ | $m_{\tilde{\nu}}$ | $m_{\tilde{q}}$ | $m_{\tilde{t}_1}$ |
|----------|----------------|--------------|-----------|-----------|----------------|----------------|----------------|----------------|----------------|----------------|-------------|--------------|
| $SD_1^+$ | 471            | 15.0         | 112       | 357       | 67.4          | 129           | 129           | 193           | 157           | 175           | 464         | 289          |
| $SD_2^+$ | 503            | 5.0          | 110       | 424       | 71.1          | 136           | 135           | 205           | 166           | 190           | 496         | 303          |
| $SD_3^+$ | 510            | 2.0          | 93.7      | 542       | 70.5          | 136           | 135           | 206           | 167           | 196           | 503         | 308          |
| $SD_1^-$ | 471            | 15.0         | 111       | 357       | 69.1          | 134           | 134           | 193           | 157           | 176           | 464         | 301          |
| $SD_2^-$ | 503            | 5.0          | 105       | 424       | 75.4          | 149           | 149           | 205           | 166           | 190           | 496         | 339          |
| $SD_3^-$ | 510            | 2.0          | 78.5      | 542       | 78.0          | 159           | 159           | 206           | 167           | 196           | 503         | 392          |
### Table 1c: A tabulation of supersymmetric particle masses for the $MS$ scenarios delineated in Fig. 1c.

| Scenario | $m_{\tilde{g}}$ | $\tan \beta$ | $m_{h^0}$ | $m_{A^0}$ | $m_{\tilde{\chi}^0_1}$ | $m_{\tilde{\chi}^0_2}$ | $m_{\tilde{\chi}^0_{1,2}}$ | $m_{\tilde{\chi}^+_1}$ | $m_{\tilde{\chi}^0_{1,2}}$ | $m_{\tilde{\chi}^0_{1,2}}$ | $m_{\tilde{\chi}^0_{1,2}}$ |
|----------|----------|----------|----------|----------|-----------------|-----------------|-----------------|-----------------|-----------------|-----------------|-----------------|
| $MS_{1}^{+}$ | 296      | 2.2      | 78.3     | 220      | 26.0            | 62.6            | 51.3            | 82.5            | 52.9            | 51.0            | 257             |
| $MS_{2}^{+}$ | 324      | 9.0      | 102      | 180      | 40.3            | 74.4            | 70.7            | 92.6            | 60.9            | 48.2            | 281             |
| $MS_{3}^{+}$ | 424      | 8.0      | 106      | 252      | 58.4            | 108             | 106             | 115             | 71.0            | 83.6            | 370             |
| $MS_{4}^{+}$ | 471      | 4.5      | 104      | 300      | 65.2            | 122             | 119             | 125             | 75.2            | 99.0            | 410             |
| $MS_{5}^{+}$ | 491      | 2.2      | 92.0     | 387      | 67.0            | 127             | 125             | 127             | 74.3            | 110             | 428             |
| $MS_{6}^{+}$ | 492      | 7.0      | 108      | 300      | 69.9            | 130             | 129             | 130             | 78.1            | 104             | 429             |
| $MS_{7}^{+}$ | 550      | 5.0      | 108      | 350      | 79.2            | 149             | 148             | 143             | 84.3            | 121             | 479             |
| $MS_{8}^{+}$ | 605      | 2.0      | 82.5     | 510      | 87.3            | 167.4           | 166             | 154.1           | 87.4            | 141.1           | 528             |
| $MS_{1}^{-}$ | 296      | 2.2      | 68.0     | 220      | 47.0            | 98.5            | 98.7            | 82.5            | 52.9            | 51.0            | 257             |
| $MS_{2}^{-}$ | 324      | 9.0      | 99.1     | 180      | 47.1            | 84.6            | 85.5            | 92.6            | 60.9            | 48.2            | 281             |
| $MS_{3}^{-}$ | 368      | 2.0      | 69.4     | 300      | 58.0            | 121             | 120             | 98.1            | 59.5            | 76.1            | 321             |
| $MS_{4}^{-}$ | 373      | 4.5      | 94.5     | 230      | 56.7            | 108             | 109             | 103             | 64.7            | 68.7            | 325             |
| $MS_{5}^{-}$ | 400      | 7.0      | 101      | 238      | 59.8            | 112             | 113             | 109             | 68.3            | 76.2            | 349             |
| $MS_{6}^{-}$ | 450      | 3.0      | 89.1     | 311      | 69.8            | 139             | 139             | 119             | 71.4            | 95.1            | 392             |
| $MS_{7}^{-}$ | 477      | 5.0      | 101      | 300      | 72.9            | 141             | 141             | 127             | 76.1            | 100             | 416             |

### Table 1d: A tabulation of supersymmetric particle masses for the $SMS$ scenarios delineated in Fig. 1d.

| Scenario | $m_{\tilde{g}}$ | $\tan \beta$ | $m_{h^0}$ | $m_{A^0}$ | $m_{\tilde{\chi}^0_1}$ | $m_{\tilde{\chi}^0_2}$ | $m_{\tilde{\chi}^0_{1,2}}$ | $m_{\tilde{\chi}^+_1}$ | $m_{\tilde{\chi}^0_{1,2}}$ | $m_{\tilde{\chi}^0_{1,2}}$ | $m_{\tilde{\chi}^0_{1,2}}$ |
|----------|----------|----------|----------|----------|-----------------|-----------------|-----------------|-----------------|-----------------|-----------------|-----------------|
| $SMS_{1}^{+}$ | 390      | 8.0      | 106      | 274      | 52.9            | 98.8            | 97.4            | 124             | 82.1            | 96.2            | 370             |
| $SMS_{2}^{+}$ | 570      | 2.2      | 97.6     | 529      | 81.0            | 156             | 155             | 173             | 108             | 160             | 541             |
| $SMS_{3}^{+}$ | 591      | 4.5      | 110      | 453      | 85.5            | 164             | 164             | 180             | 114             | 163             | 561             |
| $SMS_{1}^{-}$ | 290      | 9.0      | 98.6     | 192      | 41.5            | 76.2            | 77.0            | 97.5            | 67.6            | 57.1            | 275             |
| $SMS_{2}^{-}$ | 300      | 3.0      | 81.7     | 234      | 45.9            | 92.9            | 93.4            | 98.2            | 66.5            | 67.2            | 284             |
| $SMS_{3}^{-}$ | 300      | 5.0      | 93.4     | 212      | 44.6            | 85.0            | 86.0            | 99.6            | 68.3            | 63.4            | 285             |
| $SMS_{4}^{-}$ | 301      | 6.0      | 95.9     | 208      | 44.2            | 83.2            | 84.1            | 100             | 68.7            | 62.9            | 285             |
Table 2a: A tabulation of some important branching ratios for the $D$ scenarios delineated in Fig. 1a. Results are quoted for a single $l$ or $\nu$ type (e.g. $l = e$); for the $\tilde{\chi}_2^0$ particle-antiparticle and antiparticle-particle channels are summed together.

| Scenario | $\tilde{\chi}_1^+ \rightarrow \bar{\nu} l$ l-3-body | $\tilde{\chi}_2^0 \rightarrow \bar{\nu} \bar{l}$ l-3-body | $l_L \rightarrow \tilde{\chi}_1^+ \nu$ | $\tilde{\chi}_1^+ l$ | $\tilde{\chi}_2^0 \nu$ |
|----------|------------------|------------------|------------------|------------------|------------------|
| $D_1^+$  | 0.0 0.21         | 0.0 0.0 0.19     | 0.64 ~ 0.36     | 0.48 0.025       |
| $D_2^+$  | 0.0 0.29         | 0.0 0.0 0.046    | 0.54 0.050      | 0.16 0.005       |
| $D_3^+$  | 0.0 0.30         | 0.0 0.0 0.041    | 0.51 0.094      | 0.085 0.004      |
| $D_4^+$  | 0.0 0.21         | 0.0 0.0 0.15     | 0.59 0.037      | 0.37 0.046       |
| $D_5^+$  | 0.0 0.21         | 0.0 0.0 0.16     | 0.52 0.16       | 0.26 0.063       |
| $D_6^+$  | 0.0 0.23         | 0.004 0.0 0.04   | 0.48 0.22       | 0.18 0.049       |
| $D_7^+$  | 0.0 0.20         | 0.089 0.0 0.044  | 0.46 0.27       | 0.20 0.06        |
| $D_8^+$  | 0.0 0.14         | 0.096 0.0 0.069  | 0.40 0.38       | 0.21 0.077       |
| $D_1^-$  | 0.33 ~ 0         | 0.006 0.32 ~ 0   | 0.0 1.0         | 0.0 0.0          |
| $D_2^-$  | 0.33 0.001       | ~ 0 0.33 0.001   | 0.17 0.70       | 0.13 0.0         |
| $D_3^-$  | 0.33 ~ 0         | 0.0 0.33 0.001   | 0.29 0.31       | 0.40 0.0         |
| $D_4^-$  | 0.32 0.01        | 0.004 0.32 ~ 0   | 0.02 0.97       | 0.01 0.0         |
| $D_5^-$  | 0.33 ~ 0         | 0.11 0.22 0.001  | 0.43 0.24       | 0.33 0.0         |
| $D_6^-$  | 0.32 0.007       | ~ 0 0.32 0.007   | 0.16 0.74       | 0.096 0.0        |
| $D_7^-$  | 0.31 0.014       | ~ 0 0.30 0.026   | 0.22 0.65       | 0.13 0.0         |
| $D_8^-$  | 0.26 0.036       | 0.02 0.23 0.076  | 0.19 0.70       | 0.11 0.0         |
| $D_9^-$  | 0.036 0.022      | ~ 0 0.020 0.015  | 0.011 0.98      | 0.005 0.0        |

Table 2b: A tabulation of some important branching ratios for the $SD$ scenarios delineated in Fig. 1b. Results are quoted for a single $l$ or $\nu$ type (e.g. $l = e$); for the $\tilde{\chi}_2^0$ particle-antiparticle and antiparticle-particle channels are summed together.

| Scenario | $\tilde{\chi}_1^+ \rightarrow \bar{\nu} l$ l-3-body | $\tilde{\chi}_2^0 \rightarrow \bar{\nu} \bar{l}$ l-3-body | $l_L \rightarrow \tilde{\chi}_1^+ \nu$ | $\tilde{\chi}_1^+ l$ | $\tilde{\chi}_2^0 \nu$ |
|----------|------------------|------------------|------------------|------------------|------------------|
| $SD_1^+$ | 0.0 0.16         | 0.0 0.0 0.09     | 0.51 0.18 0.30   | 0.49 0.20        |
| $SD_2^+$ | 0.0 0.14         | 0.0 0.0 0.10     | 0.53 0.15 0.31   | 0.50 0.20        |
| $SD_3^+$ | 0.0 0.14         | 0.0 0.0 0.13     | 0.55 0.13 0.32   | 0.50 0.20        |
| $SD_1^-$ | 0.0 0.19         | 0.0 0.0 0.10     | 0.49 0.23 0.28   | 0.48 0.21        |
| $SD_2^-$ | 0.0 0.17         | 0.0 0.0 0.13     | 0.47 0.28 0.25   | 0.46 0.22        |
| $SD_3^-$ | 0.0 0.11         | 0.0 0.0 0.004    | 0.43 0.37 0.20   | 0.43 0.20        |
Table 2c: A tabulation of some important branching ratios for the $MS$ scenarios delineated in Fig. 1c. Results are quoted for a single $l$ or $\nu$ type (e.g. $l = e$); for the $\tilde{\chi}^0_2$ particle-antiparticle and antiparticle-particle channels are summed together.

| Scenario | $\tilde{l}l$ | $\tilde{\chi}^+_1 \rightarrow \tilde{l}_L \nu$ l-3-body | $\tilde{\nu}_R$ | $\tilde{\nu}_L$ $\nu \tilde{\nu}$ l-3-body | $\tilde{l}_L \rightarrow \tilde{\chi}^+_1 \nu$ $\tilde{\chi}^0_1 \tilde{\chi}^0_2$ |
|----------|-------------|---------------------------------|----------------|---------------------------------|---------------------------------|
| $MS^+_1$ | 0.27        | 0.0                             | 0.18           | 0.0                             | 0.15 $\sim 0$                   | 0.71 $\sim 0$ 0.29             |
| $MS^+_2$ | 0.33        | 0.0                             | $\sim 0$       | 0.049                           | 0.28 $\sim 0$                   | 0.54 0.12 0.34                  |
| $MS^+_3$ | 0.33        | 0.0                             | 0.001          | 0.048                           | 0.28 $\sim 0$                   | 0.25 0.64 0.11                  |
| $MS^+_4$ | 0.32        | 0.0                             | 0.002          | 0.058                           | 0.27 $\sim 0$                   | 0.13 0.84 0.03                  |
| $MS^+_5$ | 0.31        | 0.0                             | 0.006          | 0.080                           | 0.24 $\sim 0$                   | 0.04 0.96 $\sim 0$             |
| $MS^+_6$ | 0.32        | 0.0                             | 0.002          | 0.058                           | 0.27 $\sim 0$                   | 0.13 0.84 0.03                  |
| $MS^+_7$ | 0.29        | 0.024                           | $\sim 0$       | 0.028                           | 0.382 $\sim 0$                  | 0.0 1.0 0.0                     |
| $MS^-_1$ | 0.21        | 0.078                           | 0.014          | 0.022                           | 0.106 $\sim 0$                  | 0.0 1.0 0.0                     |
| $MS^-_2$ | 0.25        | 0.084                           | $\sim 0$       | 0.004                           | 0.056 $\sim 0$                  | 0.0 1.0 0.0                     |
| $MS^-_3$ | 0.33        | 0.0                             | $\sim 0$       | 0.008                           | 0.32 $\sim 0$                   | 0.14 0.73 0.13                  |
| $MS^-_4$ | 0.22        | 0.11                            | $\sim 0$       | 0.004                           | 0.085 $\sim 0$                  | 0.0 1.0 0.0                     |
| $MS^-_5$ | 0.30        | 0.030                           | $\sim 0$       | 0.022                           | 0.31 $\sim 0$                   | 0.0 1.0 0.0                     |
| $MS^-_6$ | 0.32        | 0.016                           | $\sim 0$       | 0.003                           | 0.011 $\sim 0$                  | 0.0 1.0 0.0                     |
| $MS^-_7$ | 0.23        | 0.098                           | 0.001          | 0.003                           | 0.084 $\sim 0$                  | 0.0 1.0 0.0                     |

Table 2d: A tabulation of some important branching ratios for the $SMS$ scenarios delineated in Fig. 1d. Results are quoted for a single $l$ or $\nu$ type (e.g. $l = e$); for the $\tilde{\chi}^0_2$ particle-antiparticle and antiparticle-particle channels are summed together.

| Scenario | $\tilde{\chi}^+_1 \rightarrow \tilde{l}_L \nu$ l-3-body | $\tilde{\chi}^0_2 \rightarrow \tilde{l}_R \nu \tilde{l}_L$ l-3-body | $\tilde{l}_L \rightarrow \tilde{\chi}^+_1 \nu$ $\tilde{\chi}^0_1 l_{\chi}^0_2 l_{\chi}^0_2$ |
|----------|------------------------------------------------------|---------------------------------------------------------------|--------------------------------------------------|
| $SMS^+_1$ | 0.35 0.02                                            | 0.16 0.16 0.006                                              | 0.48 0.23 0.29                                    |
| $SMS^+_2$ | 0.0 0.21                                             | 0.14 0.0 0.08                                                | 0.31 0.53 0.16                                    |
| $SMS^+_3$ | 0.006 0.076                                          | 0.20 0.0 0.10                                                | 0.25 0.62 0.13                                    |
| $SMS^-_1$ | 0.33 $\sim 0$                                        | 0.002 0.33 $\sim 0$                                         | 0.37 0.36 0.27                                    |
| $SMS^-_2$ | 0.33 0.004                                           | 0.002 0.33 0.004                                             | 0.05 0.92 0.03                                    |
| $SMS^-_3$ | 0.33 0.001                                           | $\sim 0$ 0.33 0.001                                          | 0.25 0.59 0.16                                    |
| $SMS^-_4$ | 0.33 $\sim 0$                                        | $\sim 0$ 0.33 0.001                                          | 0.29 0.51 0.20                                    |

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Table 3a: Per cent of SUSY particles produced at Tevatron in $2 \rightarrow 2$ subprocesses for $D$ scenarios. The quantity $\tilde{\chi}\tilde{\chi}$ includes $\tilde{\chi}^\pm \tilde{\chi}_0^0$, while $\tilde{q}\tilde{q}$ doesn’t include $\tilde{t}_1\tilde{t}_1$. The remaining sparticle production fraction is taken up by associated production mechanisms.

| Scenario | $\tilde{g}\tilde{g}$ | $\tilde{q}\tilde{q}$ | $\tilde{q}\tilde{q}$ | $\tilde{t}_1\tilde{t}_1$ | $\tilde{u}\tilde{u}$ $\tilde{c}\tilde{c}$ | $\tilde{\chi}\tilde{\chi}$ $\tilde{\chi}^\pm \tilde{\chi}_0^0$ | $l\bar{l}$ | $l\bar{\nu} + \nu\bar{\nu}$ |
|----------|-----------------|-----------------|-----------------|-----------------|-----------------|-----------------|-----------------|-----------------|
| $D_1^+$  | 0.5             | 2.5             | 6.6             | 5.3             | 79              | 20              | 0.8             | 2.1             |
| $D_2^+$  | 0.7             | 4.4             | 12              | 7.2             | 63              | 30              | 1.3             | 6.4             |
| $D_3^+$  | 0.8             | 4.1             | 13              | 6.8             | 59              | 32              | 3.7             | 7.1             |
| $D_4^+$  | 0.3             | 1.5             | 6.3             | 7.6             | 74              | 35              | 1.5             | 4.6             |
| $D_5^+$  | 0.07            | 0.3             | 2.9             | 5.9             | 78              | 43              | 3.1             | 6.9             |
| $D_6^+$  | 0.04            | 0.23            | 3.6             | 5.5             | 76              | 42              | 3.7             | 7.2             |
| $D_7^+$  | 0.01            | 0.05            | 1.2             | 3.7             | 82              | 46              | 4.4             | 7.2             |
| $D_8^+$  | 0.00            | 0.00            | 0.14            | 2.0             | 85              | 47              | 6.0             | 6.2             |
| $D_9^+$  | 3.9             | 23              | 41              | 1.3             | 10              | 4.6             | 1.9             | 10              |
| $D_1^-$  | 3.1             | 19              | 33              | 2.8             | 17              | 9.9             | 1.6             | 14              |
| $D_2^-$  | 1.0             | 7.5             | 21              | 6.1             | 44              | 26              | 2.3             | 10              |
| $D_3^-$  | 1.7             | 10              | 31              | 2.5             | 29              | 15              | 3.9             | 12              |
| $D_4^- $ | 1.0             | 5.4             | 16              | 7.0             | 49              | 29              | 4.2             | 9.9             |
| $D_5^- $ | 0.7             | 4.3             | 18              | 4.4             | 46              | 26              | 4.3             | 14              |
| $D_6^- $ | 0.1             | 0.6             | 5.6             | 3.8             | 66              | 38              | 6.4             | 12              |
| $D_7^- $ | 0.02            | 0.08            | 1.9             | 2.4             | 73              | 42              | 7.5             | 12              |
| $D_8^- $ | 0.00            | 0.00            | 0.3             | 0.7             | 74              | 38              | 12              | 13              |

Table 3b: Per cent of SUSY particles produced at Tevatron in $2 \rightarrow 2$ subprocesses for $SD$ scenarios. The quantity $\tilde{\chi}\tilde{\chi}$ includes $\tilde{\chi}^\pm \tilde{\chi}_0^0$, while $\tilde{q}\tilde{q}$ doesn’t include $\tilde{t}_1\tilde{t}_1$. The remaining sparticle production fraction is taken up by associated production mechanisms.

| Scenario | $\tilde{g}\tilde{g}$ | $\tilde{q}\tilde{q}$ | $\tilde{q}\tilde{q}$ | $\tilde{t}_1\tilde{t}_1$ | $\tilde{u}\tilde{u}$ $\tilde{c}\tilde{c}$ | $\tilde{\chi}\tilde{\chi}$ $\tilde{\chi}^\pm \tilde{\chi}_0^0$ | $l\bar{l}$ | $l\bar{\nu} + \nu\bar{\nu}$ |
|----------|-----------------|-----------------|-----------------|-----------------|-----------------|-----------------|-----------------|-----------------|
| $SD_1^+$ | 0.03            | 0.1             | 0.8             | 4.7             | 88              | 53              | 1.4             | 2.5             |
| $SD_2^+$ | 0.02            | —               | 0.3             | 3.5             | 91              | 54              | 1.4             | 2.0             |
| $SD_3^+$ | 0.02            | 0.01            | 0.4             | 2.8             | 92              | 55              | 1.3             | 1.5             |
| $SD_1^- $| 0.09            | 0.1             | 0.8             | 3.7             | 87              | 53              | 2.0             | 3.1             |
| $SD_2^- $| 0.02            | 0.05            | 0.5             | 1.8             | 90              | 52              | 2.5             | 3.1             |
| $SD_3^- $| 0.03            | 0.02            | 0.5             | 0.4             | 89              | 50              | 3.1             | 3.9             |
Table 3c: Per cent of SUSY particles produced at Tevatron in $2 \rightarrow 2$ subprocesses for $MS$ scenarios. The quantity $\tilde{\chi}\tilde{\chi}$ includes $\tilde{\chi}_{1}^{+}\tilde{\chi}_{2}^{0}$, while $\tilde{q}\tilde{q}$ doesn’t include $\tilde{t}_{1}\tilde{t}_{1}$. The remaining sparticle production fraction is taken up by associated production mechanisms.

| Scenario | $\tilde{g}\tilde{g}$ | $\tilde{q}\tilde{q}$ | $\tilde{t}_{1}\tilde{t}_{1}$ | $\tilde{\chi}\tilde{\chi}$ | $\tilde{\chi}_{1}^{+}\tilde{\chi}_{2}^{0}$ | $ll$ | $l\tilde{\nu} + \nu\tilde{l}$ |
|----------|----------------------|----------------------|-----------------------------|-----------------------------|--------------------------------|------|--------------------------|
| $MS_{1}^{+}$ | 0.4                  | 2.4                  | 6.8                         | 3.7                         | 72                             | 19   | 2.1                      | 9.6 |
| $MS_{2}^{+}$ | 0.4                  | 2.3                  | 9.5                         | 3.4                         | 47                             | 33   | 14                        | 20  |
| $MS_{3}^{+}$ | 0.1                  | 0.4                  | 3.6                         | 2.6                         | 57                             | 31   | 11                        | 22  |
| $MS_{4}^{+}$ | 0.00                 | 0.1                  | 2.0                         | 2.4                         | 59                             | 33   | 14                        | 20  |
| $MS_{5}^{+}$ | 0.04                 | 0.15                 | 1.3                         | 3.0                         | 60                             | 33   | 18                        | 16  |
| $MS_{6}^{+}$ | 0.02                 | 0.09                 | 1.4                         | 1.6                         | 57                             | 31   | 17                        | 22  |
| $MS_{7}^{+}$ | 0.01                 | 0.02                 | 0.6                         | 1.0                         | 56                             | 31   | 22                        | 20  |
| $MS_{8}^{+}$ | 0.00                 | 0.00                 | 0.2                         | 1.1                         | 53                             | 29   | 29                        | 16  |
| $MS_{1}^{-}$ | 1.2                  | 7.9                  | 24                          | 1.5                         | 16                             | 7.9  | 7.8                      | 34  |
| $MS_{2}^{-}$ | 0.5                  | 3.0                  | 13                          | 2.8                         | 30                             | 18   | 5.3                      | 41  |
| $MS_{3}^{-}$ | 0.5                  | 2.3                  | 14                          | 1.2                         | 27                             | 13   | 18                        | 30  |
| $MS_{4}^{-}$ | 0.3                  | 1.6                  | 10                          | 2.3                         | 35                             | 20   | 12                        | 33  |
| $MS_{5}^{-}$ | 0.1                  | 0.9                  | 7.2                         | 2.3                         | 42                             | 24   | 13                        | 31  |
| $MS_{6}^{-}$ | 0.05                 | 0.2                  | 3.8                         | 1.2                         | 37                             | 20   | 22                        | 32  |
| $MS_{7}^{-}$ | 0.01                 | 0.1                  | 2.4                         | 1.2                         | 42                             | 24   | 22                        | 30  |

Table 3d: Per cent of SUSY particles produced at Tevatron in $2 \rightarrow 2$ subprocesses for $SMS$ scenarios. The quantity $\tilde{\chi}\tilde{\chi}$ includes $\tilde{\chi}_{1}^{+}\tilde{\chi}_{2}^{0}$, while $\tilde{q}\tilde{q}$ doesn’t include $\tilde{t}_{1}\tilde{t}_{1}$. The remaining sparticle production fraction is taken up by associated production mechanisms.

| Scenario | $\tilde{g}\tilde{g}$ | $\tilde{q}\tilde{q}$ | $\tilde{t}_{1}\tilde{t}_{1}$ | $\tilde{\chi}\tilde{\chi}$ | $\tilde{\chi}_{1}^{+}\tilde{\chi}_{2}^{0}$ | $ll$ | $l\tilde{\nu} + \nu\tilde{l}$ |
|----------|----------------------|----------------------|-----------------------------|-----------------------------|--------------------------------|------|--------------------------|
| $SMS_{1}^{+}$ | 0.2                  | 0.6                  | 3.1                         | 3.7                         | 72                             | 42   | 5.7                      | 11  |
| $SMS_{2}^{+}$ | $\sim 0$             | 0.01                 | 0.03                        | 1.2                         | 77                             | 44   | 12                        | 8.6 |
| $SMS_{3}^{+}$ | $\sim 0$             | 0.04                 | 0.54                        | 0.54                        | 75                             | 44   | 14                        | 10  |
| $SMS_{1}^{-}$ | 1.4                  | 6.4                  | 14                          | 3.8                         | 43                             | 27   | 3.5                      | 20  |
| $SMS_{2}^{-}$ | 1.4                  | 8.6                  | 18                          | 2.6                         | 35                             | 19   | 6.0                      | 19  |
| $SMS_{3}^{-}$ | 1.0                  | 6.4                  | 15                          | 3.3                         | 41                             | 25   | 4.8                      | 20  |
| $SMS_{4}^{-}$ | 1.1                  | 5.9                  | 14                          | 3.8                         | 43                             | 26   | 4.8                      | 19  |
Table 4: Cross sections in fb after cuts given in the text, for various background processes. The $W$ and $Z$ backgrounds include decays to $\tau$ leptons. Non-parenthetical (parenthetical) numbers in the 3$l$ case are cross sections for 0 + 1 jets (any number of jets).

| Process     | $E_T$  | $1l$     | OS | SS  | 3$l$ | 4$l$ |
|-------------|--------|----------|----|-----|------|------|
| $W + jets$  | 1450   | $1.1 \times 10^6$ | —  | —   | —    | —    |
| $Z + jets$  | 1065   | 6850     | —  | —   | —    | —    |
| $t\bar{t}(178)$ | 200 | 491      | 0.02 | 0.35 | 0.07 (0.18) | —    |
| $WW$        | 1.2    | 106      | 31.7 | —   | 0.01 (0.05) | —    |
| $WZ$        | 0.1    | 3.9      | 0.17 | 1.8  | 0.3 (0.4)  | —    |
| total       | 2716.3 | $1.1 \times 10^6$ | 31.9 | 2.15 | 0.38 (0.63) | —    |
Table 5a: Cross sections in fb after cuts given in the text, for various signals for the \( D \) scenarios. Non-parenthetical (parenthetical) numbers in the 3l case are cross sections for 0 + 1 jets (any number of jets).

| Scenario \( m_{\tilde{g}} \) \( \tan \beta \) \( E_T \) \( 1l \) \( OS \) \( SS \) 3l 3l/SS 4l 5l |
|-----------------|--------|------|-----|-----|-----|-----|-----|-----|-----|
| \( D_1^+ \)     | 282    | 2.0  | 479 | 2010| 53.9| 53.9| 112 (168)| 2.1 | 32.4| 4.3 |
| \( D_2^+ \)     | 295    | 9.0  | 301 | 1050| 16.2| 18.6| 16.2 (26.8)| 0.9 | —   | —   |
| \( D_3^+ \)     | 310    | 15.0 | 304 | 310 | 9.7 | 7.0 | 1.1 (3.2) | 0.2 | —   | —   |
| \( D_4^+ \)     | 346    | 3.2  | 103 | 491 | 16.1| 20.4| 42.0 (56.2)| 2.1 | 2.0 | 0.8 |
| \( D_5^+ \)     | 431    | 4.5  | 23.5| 128 | 8.8 | 4.8 | 16.0 (22.8)| 3.3 | 0.8 | —   |
| \( D_6^+ \)     | 435    | 8.0  | 20.2| 120 | 5.5 | 3.2 | 4.7 (6.3) | 1.5 | 0.3 | —   |
| \( D_7^+ \)     | 503    | 5.0  | 9.2 | 55.0| 2.8 | 3.0 | 4.2 (5.7) | 1.4 | 0.3 | 0.03|
| \( D_8^+ \)     | 609    | 2.0  | 1.6 | 20.6| 2.2 | 1.1 | 4.1 (5.6) | 3.8 | 0.2 | 0.03|
| \( D_9^+ \)     | 232    | 2.0  | 1340| 2530| 39.3| 78.5| 0.0 (5.8) | —   | —   | —   |
| \( D_{10}^+ \)  | 243    | 3.2  | 1080| 2130| 36.6| 64.4| 1.3 (10.1)| 0.02| 1.3 | —   |
| \( D_{11}^+ \)  | 295    | 9.0  | 297 | 483 | 6.7 | 6.3 | 1.0 (1.4) | 0.2 | —   | —   |
| \( D_{12}^+ \)  | 301    | 2.2  | 196 | 480 | 17.3| 5.9 | 5.4 (8.8) | 0.9 | 0.3 | 0.3 |
| \( D_{13}^- \)  | 310    | 15.0 | 278 | 261 | 8.9 | 2.8 | 0.8 (1.2) | 0.3 | —   | —   |
| \( D_{14}^- \)  | 346    | 3.2  | 72.8| 207 | 6.1 | 1.2 | 1.2 (2.5) | 1.0 | —   | —   |
| \( D_{15}^- \)  | 431    | 4.5  | 12.4| 30.3| 3.4 | 0.1 | 0.9 (1.1) | 6.7 | —   | —   |
| \( D_{16}^- \)  | 503    | 5.0  | 5.0 | 15.1| 3.5 | 0.04| 0.5 (0.8)| 13.0| 0.04| 0.02|
| \( D_{17}^- \)  | 609    | 2.0  | 4.0 | 8.5 | 1.1 | 0.04| 0.07 (0.08)| 1.9 | —   | —   |

Table 5b: Cross sections in fb after cuts given in the text, for various signals for the \( SD \) scenarios. Non-parenthetical (parenthetical) numbers in the 3l case are cross sections for 0 + 1 jets (any number of jets).

| Scenario | \( m_{\tilde{g}} \) \( \tan \beta \) \( E_T \) \( 1l \) \( OS \) \( SS \) 3l 3l/SS 4l 5l |
|----------|--------|------|-----|-----|-----|-----|-----|-----|-----|
| \( SD_1^+ \)     | 471    | 15.0 | 15.4| 51.4| 2.3 | 1.1 | 4.1 (5.6) | 3.6 | 0.21| —   |
| \( SD_2^+ \)     | 503    | 5.0  | 12.1| 40.0| 1.5 | 0.93| 4.0 (5.1) | 4.3 | 0.25| 0.03|
| \( SD_3^+ \)     | 510    | 2.0  | 10.9| 38.9| 1.4 | 1.0 | 4.2 (5.2) | 4.1 | 0.12| 0.03|
| \( SD_1^- \)     | 471    | 15.0 | 11.2| 45.4| 2.3 | 1.5 | 4.9 (6.7) | 3.2 | 0.24| 0.03|
| \( SD_2^- \)     | 503    | 5.0  | 7.7 | 27.4| 1.6 | 0.7 | 3.9 (5.0) | 5.5 | 0.16| 0.05|
| \( SD_3^- \)     | 510    | 2.0  | 3.7 | 18.4| 0.8 | 0.1 | 0.14 (0.28)| 1.0 | 0.01| —   |
### Table 5c: Cross sections in fb after cuts given in the text, for various signals for the MS scenarios. Non-parenthetical (parenthetical) numbers in the 3l case are cross sections for 0 + 1 jets (any number of jets).

| Scenario | $m_{\tilde{g}}$ | $\tan \beta$ | $E_T$ | 1l | OS | SS | 3l | 3l/SS | 4l | 5l |
|----------|-----------------|--------------|-------|----|----|----|----|-------|----|----|
| MS$^+_1$ | 296             | 2.2          | 411   | 393| 75.4| 7.0| 26.3|(43.9)| 3.8| 5.3|—   |
| MS$^+_2$ | 324             | 9.0          | 141   | 983| 51.1| 17.4| 24.6|(47.2)| 1.4| 3.2|—   |
| MS$^+_3$ | 424             | 8.0          | 15.9  | 214| 8.5 | 2.3 | 7.7 |(10.7)| 3.4| 0.34|—   |
| MS$^+_4$ | 471             | 4.5          | 8.3   | 126| 6.1 | 2.9 | 5.4 |(6.5)| 1.9| 0.07|—   |
| MS$^+_5$ | 491             | 2.2          | 7.9   | 96.5| 4.2 | 3.0 | 3.3 |(4.2)| 1.1| 0.06|—   |
| MS$^+_6$ | 492             | 7.0          | 5.1   | 96.9| 5.7 | 1.3 | 1.9 |(2.5)| 1.5| 0.15|—   |
| MS$^+_7$ | 550             | 5.0          | 27.6  | 53.9| 3.1 | 2.6 | 1.2 |(1.5)| 0.45| 0.03|—   |
| MS$^+_8$ | 603             | 2.0          | 1.9   | 31.2| 2.3 | 1.5 | 2.6 |(3.4)| 1.7| 0.05|—   |
| MS$^-_1$ | 296             | 2.2          | 257   | 809| 22.7| 22.2| 14.1|(27.8)| 0.6|—    |0.5 |
| MS$^-_2$ | 324             | 9.0          | 119   | 784| 19.2| 8.4 | 2.8 |(6.5)| 0.3| 0.9 |—   |
| MS$^-_3$ | 368             | 2.0          | 38.3  | 242| 13.7| 5.1 | 10.8|(18.2)| 2.1| 0.4 |—   |
| MS$^-_4$ | 373             | 4.5          | 34.6  | 306| 14.1| 7.4 | 1.9 |(4.4)| 0.2|—    |—   |
| MS$^-_5$ | 400             | 7.0          | 5.1   | 96.9| 5.7 | 1.3 | 1.9 |(2.5)| 1.5| 0.15|—   |
| MS$^-_6$ | 450             | 3.0          | 6.5   | 101| 1.9 | 6.4 | 9.0 |(16.0)| 3.4| 0.16|—   |
| MS$^-_7$ | 477             | 5.0          | 2.7   | 31.2| 2.3 | 1.5 | 2.6 |(3.4)| 1.7| 0.05|—   |

### Table 5d: Cross sections in fb after cuts given in the text, for various signals for the SMS scenarios. Non-parenthetical (parenthetical) numbers in the 3l case are cross sections for 0 + 1 jets (any number of jets).

| Scenario | $m_{\tilde{g}}$ | $\tan \beta$ | $E_T$ | 1l | OS | SS | 3l | 3l/SS | 4l | 5l |
|----------|-----------------|--------------|-------|----|----|----|----|-------|----|----|
| SMS$^+_1$ | 390             | 8.0          | 31.6  | 55.5| 23.8| 0.81| 4.6 |(7.4)| 5.7| 0.94|—   |
| SMS$^+_2$ | 570             | 2.2          | 3.63  | 30.1| 3.2 | 0.96| 6.3 |(8.2)| 6.6| 0.12|—   |
| SMS$^+_3$ | 591             | 4.5          | 5.8   | 16.5| 2.2 | 0.59| 3.1 |(3.9)| 5.3| 0.28|0.03|
| SMS$^-_1$ | 290             | 9.0          | 207   | 864| 30.9| 10.1| 1.5 |(3.0)| 0.15|—    |—   |
| SMS$^-_2$ | 300             | 3.0          | 15.7  | 581| 21.4| 8.1 | 1.8 |(2.7)| 0.22|—    |—   |
| SMS$^-_3$ | 300             | 5.0          | 153   | 635| 23.8| 6.0 | 0.36|(3.2)| 0.06|—    |—   |
| SMS$^-_4$ | 301             | 6.0          | 138   | 627| 23.0| 5.1 | 1.1 |(4.0)| 0.21|—    |—   |
Figure 1a,b: We plot the boundaries of allowed $m_g$–$\tan \beta$ parameter space for the dilaton-equivalent and string-scale-unified dilaton-equivalent scenarios for both $\mu > 0$ and $\mu < 0$. Values of $R_{b/\tau}$ for a given $\tan \beta$ are given on the right-hand axis. The numbers indicate the different cases considered for each scenario (see text). They are positioned so that the actual $m_g$–$\tan \beta$ value associated with a point is at the lower left-hand side of the number.
Figure 1c,d: We plot the boundaries of allowed $m_{\tilde{g}}$--tan $\beta$ parameter space for the minimal-supergravity and string-scale-unified minimal-supergravity scenarios for both $\mu > 0$ and $\mu < 0$. Values of $R_{h}\tau$ for a given tan $\beta$ are given on the right-hand axis. The numbers indicate the different cases considered for each scenario (see text). They are positioned so that the actual $m_{\tilde{g}}$--tan $\beta$ value associated with a point is at the lower left-hand side of the number.
Figure 2a,b: We plot the masses of the various superpartner particles in terms of the ratio $m_i/m_\tilde{g}$ ($i$ denoting a particular sparticle) as a function of $m_\tilde{g}$ for the $D$ and $SD$ models. The scatter in the points indicates the variation as $\tan \beta$ is allowed to vary at fixed $m_\tilde{g}$. The band corresponding to a given particle type $i$ is labelled, according to its relative position at large $m_\tilde{g}$, on the r.h.s. of the graph.
Figure 2c,d: We plot the masses of the various superpartner particles in terms of the ratio $m_i/m_\tilde{g}$ ($i$ denoting a particular sparticle) as a function of $m_\tilde{g}$ for the MS and SMS models. The scatter in the points indicates the variation as tan $\beta$ is allowed to vary at fixed $m_\tilde{g}$. The band corresponding to a given particle type $i$ is labelled, according to its relative position at large $m_\tilde{g}$, on the r.h.s. of the graph.
Figure 3: The $E_T$ cross section (in fb) for $p\bar{p}$ collisions at $\sqrt{s} = 1.8$ TeV is plotted for all the numbered scenarios in Fig. 1 as a function of $m_{\tilde{g}}$. $D$, $SD$, $MS$ and $SMS$ cases are indicated by $d$, $D$, $m$ and $M$, respectively. The $m_{\tilde{g}}$ mass value can be used to identify the scenario within each such case.
Figure 4a,b: We display the same-sign dilepton (SS) and tri-lepton (3l) cross sections (fb) for $\sqrt{s} = 1.8$ TeV $p\bar{p}$ collisions for the D and SD scenarios defined in Fig. 1a,b. No restrictions are placed on the number of associated jets. The format is [SS, 3l] or (SS, 3l) where cases with square (curved) brackets have $m_{\tilde{h}_0} < 105$ GeV ($> 105$ GeV). The dashed line indicates the approximate boundary beyond which detection of a SS or 3l signal would not be possible for an integrated luminosity $L$ of 1 fb$^{-1}$ or smaller. The arrows, labelled by $\tilde{t}_R$ and $\tilde{\chi}_1^+$, indicate the approximate $m_{\tilde{g}}$ values corresponding to $m_{\tilde{t}_R} = 95$ GeV and $m_{\tilde{\chi}_1^+} = 95$ GeV, respectively. The width of each arrow reflects the range of $m_{\tilde{g}}$ values obtained as $\tan \beta$ is varied at fixed $m_{\tilde{t}_R}, m_{\tilde{\chi}_1^+} = 95$ GeV. For $m_{\tilde{t}_R} < 95$ GeV, $m_{\tilde{\chi}_1^+} < 95$ GeV, $m_{h_0} < 105$ GeV detection of $\tilde{t}_R\tilde{t}_R$, $\tilde{\chi}_1^+\tilde{\chi}_1^+$, $Zh^0$ production would be possible at LEP-200 with $L = 500$ pb$^{-1}$ - 1 fb$^{-1}$.
Figure 4c,d: We display the same-sign dilepton (SS) and tri-lepton (3l) cross sections (fb) for $\sqrt{s} = 1.8\text{ TeV}$ $\sqrt{p\bar{p}}$ collisions for the MS and SMS scenarios defined in Fig. 1c,d. No restrictions are placed on the number of associated jets. The format is [SS,3l] or (SS,3l) where cases with square (curved) brackets have $m_{h^0} < 105\text{ GeV} (> 105\text{ GeV})$. The dashed line indicates the approximate boundary beyond which detection of a SS or 3l signal would not be possible for an integrated luminosity $L$ of 1 fb$^{-1}$ or smaller. The arrows, labelled by $\tilde{l}_R$ and $\tilde{\chi}^1_1$, indicate the approximate $m_{\tilde{g}}$ values corresponding to $m_{\tilde{l}_R} = 95\text{ GeV}$ and $m_{\tilde{\chi}^1} = 95\text{ GeV}$, respectively. The width of each arrow reflects the range of $m_{\tilde{g}}$ values obtained as $\tan \beta$ is varied at fixed $m_{\tilde{l}_R}, m_{\tilde{\chi}^1} = 95\text{ GeV}$. For $m_{\tilde{l}_R} < 95\text{ GeV}$, $m_{\tilde{\chi}^1} < 95\text{ GeV}$, $m_{h^0} < 105\text{ GeV}$ detection of $\tilde{l}_R \tilde{l}_R$, $\tilde{\chi}^1_1 \tilde{\chi}^1_1$, $ZH^0$ production would be possible at LEP-200 with $L = 500\text{ pb}^{-1} - 1\text{ fb}^{-1}$.
Figure 5a,b: We plot the contours for $m_{\tilde{l}_R} = 83, 95$ GeV (dots), $m_{\tilde{\chi}^+_1} = 83, 95$ GeV (dashes) and $m_{h^0} = 81, 105$ GeV (dotdash) within the allowed $m_g^- - \tan\beta$ parameter space for the dilaton-equivalent and for the string-scale-unified dilaton-equivalent scenarios, for both $\mu > 0$ and $\mu < 0$. The lower and upper mass values for each particle type represent rough upper limits for the mass reach at LEP-II with $\sqrt{s} = 176, 200$ GeV, respectively. In the $SD^+$ case, $m_{\tilde{l}_R} \geq 83$ GeV for all allowed $m_g^- - \tan\beta$, and only the $m_{\tilde{l}_R} = 95$ GeV contour appears. In the $D^+$ and $SD^+$ windows, the $m_{h^0} = 81$ GeV contour is barely visible in the low-$\tan\beta$, low-$m_g^-$ parameter space corner.
Figure 5c,d: We plot the contours for \( m_{\tilde{l}_R} = 83, 95 \text{ GeV} \) (dots), \( m_{\tilde{\chi}_1^\pm} = 83, 95 \text{ GeV} \) (dashes) and \( m_{h^0} = 81, 105 \text{ GeV} \) (dotdash) within the allowed \( m_{\tilde{g}} - \tan \beta \) parameter space for the minimal-supergravity and for the string-scale-unified minimal-supergravity scenarios, for both \( \mu > 0 \) and \( \mu < 0 \). The lower and upper mass values for each particle type represent rough upper limits for the mass reach at LEP-II with \( \sqrt{s} = 176, 200 \text{ GeV} \), respectively. In the \( MS^- \) case, only the \( m_{\tilde{l}_R} = 83 \text{ GeV} \) contour (barely) appears since \( m_{\tilde{l}_R} \leq 95 \text{ GeV} \) throughout all of the allowed parameter space. In the \( SMS^+ \) window, the \( m_{h^0} = 81 \text{ GeV} \) contour is barely visible in the low-\( \tan \beta \), low-\( m_{\tilde{g}} \) parameter space corner. In the \( MS^- \) model, \( m_{h^0} \leq 105 \text{ GeV} \) for all allowed \( m_{\tilde{g}}, \tan \beta \) values, and only the \( m_{h^0} = 81 \text{ GeV} \) contour appears.