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Effects of $SO(10)$ $D-$Terms on SUSY Signals at the Tevatron

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Abstract

We study signals for the production of superparticles at the Tevatron in supergravity scenarios based on the Grand Unified group $SO(10)$. The breaking of this group introduces extra contributions to the masses of all scalars, described by a single new parameter. We find that varying this parameter can considerably change the size of various expected signals studied in the literature, with different numbers of jets and/or charged leptons in the final state. The ratios of these signal can thus serve as a diagnostic to detect or constrain deviations from the much–studied scenario where all scalar masses are universal at the GUT scale. Moreover, under favorable circumstances some of these signals, and/or new signals involving hard $b-$jets, should be observable at the next run of the Tevatron collider even if the average scalar mass lies well above the gluino mass.

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1 Introduction

The search for Supersymmetry (SUSY) [1], the most attractive candidate for physics beyond the Standard Model (SM), has top priority in present day high energy experiments. From the non–observation of sparticles at various colliders only lower limits on their masses have been obtained so far.

From the model building point of view the limit $M_{\widetilde{\chi}_1^\pm} > 91$ GeV [2] on the mass of the lighter chargino $\widetilde{\chi}_1^\pm$ from LEP2 is of special importance. This limit, in conjunction with the popular assumption of gaugino mass unification [1] with a common gaugino mass $m_{1/2}$ at the GUT scale (M_G), implies that the gluino ($\tilde{g}$) must be rather heavy with $m_{\tilde{g}} \gtrsim 300$ GeV.

The limit on $M_{\widetilde{\chi}_1^\pm}$ becomes much weaker if the $\widetilde{\chi}_1^\pm$ happens to be approximately degenerate with either the lightest supersymmetric particle (LSP) assumed to be the lightest neutralino ($\widetilde{\chi}_1^0$), or with a sneutrino ($\tilde{\nu}$). The first possibility is strongly disfavored in models with gaugino mass unification and radiative electroweak symmetry breaking (see below). The second possibility occurs, e.g., in mSUGRA models [1] for small values of the scalar soft breaking mass m_0. However, a recent result from the DØ collaboration [3] ruled out $m_{\tilde{g}} <$ 300 GeV for small m_0 from direct squark and gluino searches in the multi–jet plus missing E_T channel.

Like most other limits on squark and gluino production at hadron colliders, this bound holds if the squarks are degenerate with each other at M_G with a common mass m_0. They are then also approximately degenerate at the weak scale with average mass $m_{\tilde{q}}$. Under the same assumption, for much heavier squarks ($m_{\tilde{q}} \gg m_{\tilde{g}}$) only $m_{\tilde{g}} \lesssim 200$ GeV can be excluded from present Tevatron data. In this case, even gluino searches after the main injector upgrade of the Tevatron (referred to hereafter as Tevatron Run II), with an estimated integrated luminosity of 2 fb^{-1} at $\sqrt{s} = 2$ TeV, are not likely to improve the indirect limit on the gluino mass from LEP2 searches, if gaugino mass unification holds. A further luminosity upgrade, amounting to an integrated luminosity ~ 25 fb^{-1} at $\sqrt{s} = 2$ TeV (hereafter referred to as TeV33), could reach somewhat higher, but not dramatically different, masses through direct searches for gluinos and squarks [4].

As already mentioned, these estimated search limits have been derived under certain assumptions, like gaugino mass unification and degeneracy of squarks at M_G. Such assumptions about physics at high scales, although attractive, may ultimately prove wrong. It is therefore important to reexamine the search prospects at the Tevatron and its various upgrades if some or all of these assumptions are relaxed. If accompanied by suitable theoretical guidelines, this can be accomplished while still avoiding an unmanageable proliferation in the number of unknown parameters. Moreover, such an analysis indicates the feasibility of obtaining glimpses of the SUSY breaking pattern at a high scale using the Tevatron data. Finally, and most importantly, such analyses are needed to assess the robustness of the SUSY search strategies currently employed by Tevatron experiments.

One possible avenue would be to relax the gaugino mass unification condition. This has been attempted by several groups [5]. However, at least in the framework of SUSY GUT's, the assumption of gaugino mass unification appears quite natural, since it simply follows if the GUT symmetry is respected by the SUSY breaking mechanism at a high scale, irrespective of the specific choice of the GUT group. In this paper we, therefore, assume that gaugino masses do indeed unify.

The assumption that all soft breaking masses for the scalars have a universal value m_0 at scale M_G (referred to hereafter as the conventional scenario), which in turn predicts approximately degenerate squarks at the weak scale, is more model dependent. In the framework of the popular $N = 1$ SUGRA GUT models [6], where SUSY breaking in the hidden sector is transmitted to the observable sector by gravitational–strength interactions alone, two assumptions are required to obtain a degenerate scalar spectrum at scale M_G. First, the Kähler metric has to be the same for all chiral superfields. This ansatz drastically reduces the large number of parameters that may otherwise be present in the Minimal Supersymmetric Standard Model (MSSM), making the model much more predictive. It also avoids problems with flavor changing neutral currents caused by sparticle loops. This assumption guarantees degenerate sfermions at the scale Λ at which SUSY breaking is transmitted to the visible sector. By taking sfermions to be degenerate at scale M_G one therefore implicitly makes the second assumption that $\Lambda \simeq M_G \simeq 2 \cdot 10^{16}$ GeV. However, from the supergravity point of view, it is more natural to choose Λ to be near the Planck scale $M_P \simeq 2.4 \cdot 10^{18}$ GeV. Non-universality in scalar masses at scale M_G can then arise due to several reasons:
i) The running of scalar masses between M_P and M_G [7] may lead to non-universality at M_G. Within the framework of an $SO(10)$ SUSY GUT, however, the resulting non-universality is expected to be negligible for the first two generations of squarks and sleptons with small Yukawa couplings. In contrast, the soft breaking masses of the $SU(2)$ doublet Higgs bosons responsible for electroweak symmetry breaking can be changed significantly, since they have to couple to super-heavy GUT fields in order to generate very large masses for their GUT partners (doublet–triplet splitting); moreover, these Higgs bosons are usually assumed to reside in a different representation of $SO(10)$ than the sfermions do, and hence are renormalized differently by gaugino loop diagrams. The situation for third generation sfermions is intermediate between these two extremes: they reside in the same representation as first and second generation sfermions do, but at least some of them have significant Yukawa couplings. Thus only the first two generations of sleptons and squarks can safely be considered to have a common mass at M_G, which we define to be m_0. The soft breaking masses of third generation sfermions are likely to be somewhat, but not very much, smaller, while those for the Higgs bosons can be either larger or smaller than m_0. Fortunately, the physics of the Higgs sector is relatively unimportant for direct squark and gluino searches at Tevatron energies. In our study we will neglect the (probably not very large) effect of the running between Λ and M_G on third generation soft breaking terms.

ii) It has been known for some time that if the rank of a GUT group (or some symmetry group broken at an intermediate scale) is reduced by spontaneous symmetry breaking, one may obtain $D-$term contributions to scalar masses [8, 9]. These will in general differ for different members of the same GUT multiplet, leading to non-universal squark and slepton masses at the symmetry breaking scale. The size of these new contributions can be comparable to m_0. Note that these nonuniversal terms are generation–independent, so that no additional problems due to flavor changing neutral currents arise.

It has already been emphasized in [10] that right–handed down–type squarks ($\tilde{d}_R$, $\tilde{s}_R$, $\tilde{b}_R$, which are degenerate in mass) may be considerably lighter than other species of squarks ($\tilde{u}_L, \tilde{d}_L$ or $\tilde{u}_R$) due to mechanism (ii). A specific model in which this may happen is an $SO(10)$ SUSY GUT, breaking down to the SM either directly or via the Pati–Salam group

[11] at an intermediate scale [9]. In addition, the L–type (left–handed) or R–type (right–handed) sleptons may also turn out to be considerably lighter, which may yield many novel experimental signatures [10, 12]. The phenomenological consequences of $D-$terms at high scales have also been studied in the first paper of [5].

In the present analysis we shall work within the framework of a $SO(10)$ SUSY GUT breaking down to the SM directly. In this model nonuniversality in the squark and slepton masses at a high scale can be parametrized by only one extra parameter D (compared to the conventional scenario) [9]. $SO(10)$ has several attractive features: it allows to fit a complete generation of (s)quarks and (s)leptons into a single representation of the group. Moreover, it includes right–handed neutrinos, and thus naturally allows to implement the see–saw mechanism [13] to explain the small but nonvanishing neutrino masses that seem to be needed to explain recent data on neutrino oscillations. Direct breaking to $SU(3) \times SU(2) \times U(1)$ is perhaps more natural, in view of the apparent unification of all three gauge couplings in the MSSM in the absence of any intermediate scale. Similar mass patterns are, however, expected even if this GUT group first breaks to the Pati–Salam group.

Our main goal is to compare and contrast the SUSY signals predicted by the nonuniversal scenario with those of the conventional one ($D = 0$) which has 10 roughly degenerate squarks of L and R types excluding the stop ($\tilde{t}$). We shall restrict our analysis to SUSY searches at various phases of the Tevatron. Special attention will be paid to the following issues: i) Is it possible to get a signal at Run II in some region of the parameter space (not already excluded by LEP searches) of the nonuniversal model with $m_0 \gg m_{1/2}$, which corresponds to $m_{\tilde{q}} \gg m_{\tilde{g}}$ in the conventional scenario? ii) If signals are seen at the Tevatron, can they always be accommodated in both models, or might a distinction between the two be possible?

In the next section we shall present the strategy of this analysis, the choice of the parameters and the observables. The results will be presented in section 3. Our conclusions and future outlook will be given in section 4.

2 Strategy, the choice of parameters and observables

In order to analyze signals for the production of squarks and gluinos in a systematic way the large $m_{\tilde{g}} - m_{\tilde{q}}$ parameter space is usually divided into the following three broad regions: (**A**) $m_{\tilde{q}} \gg m_{\tilde{g}}$, (**B**) $m_{\tilde{q}} \approx m_{\tilde{g}}$ and (**C**) $m_{\tilde{q}} \ll m_{\tilde{g}}$.

Out of these, region (C) is not accessible to first and second generation squarks in any grand desert scenario: even if their masses vanish at the GUT scale, large contributions from gluino–quark loops imply that at the weak scale $m_{\tilde{q}} \gtrsim 0.85 m_{\tilde{g}}$. However, some third generation squarks might be significantly lighter than the gluino.

Our analysis will be based on the following strategy. We shall choose a set of GUT scale parameters *viz.*, m_0 and $m_{1/2}$, which evolves into weak scale squark and gluino masses in regions (A) and (B) in the conventional scenario. Moreover, $m_{1/2}$ is chosen such that the resulting chargino mass is always consistent with the LEP bounds. As mentioned in the Introduction, we assume that the masses of $\tilde{b}$ and $\tilde{\tau}$ at scale M_G are the same as those of the corresponding first generation sfermions, i.e. we ignore the effect of RG scaling between Λ and M_G. Since we will concentrate on a moderate value of tanβ, this should be a good approximation for both scalar τ states as well as for $\tilde{b}_R$. This assumption will probably over–estimate the masses of third generation $SU(2)$ doublet squarks slightly, but we expect this effect to be smaller than the new $D-$term contributions which are the main focus of our study.*

We have chosen tan$\beta = 3$ and the trilinear coupling $A_0 = 0$. The latter choice of course has no direct bearing on the phenomenology of the first two generations, which is the focal point of our attention. Now for $D = 0$ the magnitude of the Higgsino mass parameter μ is fixed by the radiative electroweak symmetry breaking condition [15]. We shall work throughout with negative μ, since otherwise LEP searches require $m_{\tilde{g}}$ to be too large for gluino searches at the Tevatron to be viable.

In the presence of nonuniversal Higgs masses the magnitude of μ fixed from the EW symmetry breaking condition can in principle be quite different from that in the conventional scenario [9]. However, in order to compare our results with the predictions of the conventional scenario, we shall work with the same μ as in the $D = 0$ case. This does not affect the cross sections for the production of the sparticles we are interested in, with the sole exception of the light scalar top eigenstate $\tilde{t}_1$. Its mass depends on μ through the off–diagonal entries of the stop mass matrix. However, as will be discussed in more detail below, even a rather light $\tilde{t}_1$ would not contribute very much to the signals we are interested in. We have therefore ignored its contribution completely; however, we have checked that the value of $m_{\tilde{t}_1}$ that follows from our (in this case not very reliable) assumptions satisfies the relevant search limits.

The value of μ also affects the leptonic branching ratio of the second neutralino $\widetilde{\chi}_2^0$ if it decays into 3–body modes. The reason is that the $Z\widetilde{\chi}_1^0\widetilde{\chi}_2^0$ coupling only proceeds through the higgsino components of both neutralinos, and therefore scales roughly $\propto 1/\mu^2$. As a result, slepton exchange contributions to $\widetilde{\chi}_2^0$ decay can compete with Z exchange contributions as long as slepton masses do not greatly exceed $|\mu|$. However, we will see that our model already allows considerable variation of this branching ratio even for fixed $m_{1/2}$ and m_0; allowing μ to vary away from the value predicted in the universal scenario would therefore not lead to new (combinations of) signatures, although it might shift the regions in parameter space where final states containing leptons are important.

Having chosen the parameters, it is now straightforward to compute cross sections at the parton level in both scenarios. We shall consider the cross section σ_n corresponding to events with n leptons + jets + missing transverse energy ($\not{E}_T$) coming from all types of squark and gluino (excluding $\tilde{t}$) production channels. In particular σ_0 and σ_2 will turn out to be crucial observables in our analysis. We also include the cross section of hadronically quiet trilepton events σ_{3l} in our set of observables. These arise from $\widetilde{\chi}_1^\pm - \widetilde{\chi}_2^0$ production and depend strongly on the slepton masses through the leptonic branching ratio of $\widetilde{\chi}_2^0$. In this exploratory study we do not attempt any event simulation, nor do we apply the cuts required to extract these signals. We should clarify that we assume the experimental definition of the hadronic observables σ_n to include fairly hard cuts on the energies of the jets and/or on the missing transverse energy; this is why we do not include contributions from the production of electroweak gauginos here, even though they may give the biggest contribution to the

*Note that we do not assume the Yukawa couplings to unify. In the simplest $SO(10)$ model this assumption would require a large value of tanβ, leading to reduced masses for all third generation sfermions. Signals requiring the presence of energetic electrons or muons will then be greatly depleted even in the conventional scenario [14].

total SUSY cross section at the Tevatron. The same hard cuts would also greatly reduce the contribution of a light $\tilde{t}_1$ to these final states. Moreover, we only count "primary" electrons and muons as "leptons"; we ignore leptons that may result from the decay of $c-$ or $b-$ quarks or $\tau-$leptons.

In the absence of an event simulation which gives the efficiencies of the kinematical cuts to reduce the SM background, a full appreciation of the observability of the signals we have considered is not possible. Fortunately published results of such Monte Carlo simulations allow us to at least estimate the minimal cross sections (before cuts) that lead to detectable signals at future runs of the Tevatron. The present bounds, from data taken during Run I, can be gleaned from experimental searches. Both these searches and theoretical studies of the reach of Tevatron upgrades have generally been performed in the conventional scenario. However, we believe that the efficiencies should not change by more than a factor of 2 when allowing non–vanishing values of D, as long as we compare scenarios with similar values of $m_{1/2}$.

Our estimates for the minimal detectable values of σ_0, σ_2 and σ_{3l} are collected in Table 1. The limits on σ_0 and σ_{3l} from previous Tevatron runs have been extracted from experimental searches [3, 16]. The future sensitivity in the jets plus missing E_T channels (σ_0, σ_2) as well as in the tri–lepton channel (σ_{3l}) has been estimated using results from refs.[4] and [17], respectively. We also used ref.[4] to estimate the sensitivity of Run I to σ_2, since the only published experimental analysis of the jets plus di–lepton channel [18] is based on just 19 pb^{-1} of data.[†]

Signal	Run I	Run II	TeV33
σ_0	5	0.9	0.10
σ_2	0.6	0.14	0.03
σ_{3l}	0.25	0.03	0.009

Table 1: The minimal detectable values of the cross sections for jets plus missing E_T final states with zero (σ_0) or two (σ_2) electrons or muons, as well as for the relatively hadron–free three lepton final state (σ_{3l}). The second, third and fourth column show the sensitivity from present data, from Run II (2 fb^{-1} at $\sqrt{s} = 2$ TeV) and from TeV33 (25 fb^{-1} at $\sqrt{s} = 2$ TeV), respectively. All cross sections are in pb.

3 The spectra and the signals

The sfermion masses at the GUT scale can be parametrized by [9]:

$$m^2_{\tilde{u}_L} = m^2_{\tilde{u}_R} = m^2_{\tilde{e}_R} = m^2_0 + 0.5 D m^2_0; \tag{1a}$$

$$m^2_{\tilde{d}_R} = m^2_{\tilde{e}_L} = m^2_0 - 1.5 D m^2_0, \tag{1b}$$

[†]Both CDF and DØ have announced [19] preliminary limits from searches for final states containing two leptons. However, the CDF search only covers the signal from like–sign di–leptons, which according to ref.[4] has a somewhat smaller reach than the inclusive di–lepton sample, while the DØ search did not include events containing muons.

where the unknown parameter D can be of either sign. As stated earlier we use the same m_0 for the third generation, which amounts to neglecting the running of their masses above M_G.

Requiring positivity of the sfermion mass squares at the GUT scale imposes the restriction $-2 \leq D \leq 0.66$. Slightly stronger constraints on D arise from experimental searches. In particular, LEP searches [20] for $SU(2)$ doublet and singlet sleptons translate into upper bounds on $|D|$ for positive and negative D, respectively. For large m_0, the LEP bound [21] on the mass of the lighter sbottom eigenstate $\tilde{b}_1$ gives the most stringent bound on $|D|$ for $D < 0$.

D	0	0.4	0.6	-0.75	-1.25
$m_{\tilde{u}_L}$	646.7	702.4	728.6	526.6	428.3
$m_{\tilde{d}_L}$	650.7	706.0	732.2	531.5	434.2
$m_{\tilde{u}_R}$	644.3	700.2	726.5	523.7	424.6
$m_{\tilde{d}_R}$	645.1	436.7	279.3	915.8	1058.5
$m_{\tilde{b}_L}$	525.6	592.7	623.6	367.9	203.5
$m_{\tilde{e}_L}$	606.1	376.7	170.9	888.8	1035.2
$m_{\tilde{e}_R}$	602.6	662.1	689.8	471.5	358.3
$m_{\tilde{\nu}}$	601.8	369.9	155.2	885.9	1032.8

Table 2: Squark and slepton masses in GeV at the weak scale for different values of D with $m_0 = 600$ GeV and $m_{1/2} = 105$ GeV.

We first study the signals for the choice $m_0 \approx 6m_{1/2}$, which corresponds to region (A) introduced in the last section. This is the for our purposes most interesting region, since here a signal at the Tevatron or its upgrades is not expected in the conventional scenario. Examples for sfermion masses in this region of parameter space are given in Table 2, for $m_0 = 600$ GeV and $m_{1/2} = 105$ GeV and several values of D. The gaugino spectrum is practically independent of D: $m_{\tilde{g}} = 332.1$ GeV, $M_{\tilde{\chi}_1^\pm} = 95.3$ GeV, $m_{\tilde{\chi}_2^0} = 95.5$ GeV, and $m_{\tilde{\chi}_1^0} = 46.5$ GeV.

Two points are worth noting :

- The masses of $SU(2)$ doublet sleptons $\tilde{l}_L$ and $\tilde{d}_R-$type squarks can be reduced significantly for $D > 0$; for $D \gtrsim 0.55$, $\tilde{g} \to \tilde{d}_R d$ 2–body decays become possible.

- Due to the smaller coefficient of D in eq.(1a) as compared to (1b), the reduction of the masses of $\tilde{u}_L, \tilde{d}_L$ and $\tilde{u}_R$ squarks as well as $\tilde{l}_R$ sleptons for $D < 0$ is at first less significant. However, for large negative D ($D \leq -0.8$) the $\tilde{b}_L$ becomes lighter than the gluino; in fact, as mentioned in the previous section, the lower bound on D comes from the experimental bound on $m_{\tilde{b}_1}$ in this scenario.*

*A recent study [22] of the sparticle spectrum predicted by the minimal $SO(10)$ SUSY GUT also emphasized the reduction of $m_{\tilde{b}_1}$ relative to the other squark masses. However, as mentioned in Sec. 2, minimal $SO(10)$ requires large values of tanβ, leading to a large bottom Yukawa coupling. Successful electroweak symmetry breaking then requires $D > 0$. As a result, the light $\tilde{b}_1$ found in ref.[22] is mostly an $SU(2)$ singlet, while our more general scenario allows $D < 0$ and a light $\tilde{b}_L$.

These changes of the scalar spectrum lead to intriguing features of gaugino branching ratios. In the electroweak gaugino sector, BR($\widetilde{\chi}_2^0 \to \widetilde{\chi}_1^0 l\bar{l}$) increases significantly for $D > 0$, due to the reduction of $m_{\tilde{l}_L}$. For example, we have BR($\widetilde{\chi}_2^0 \to \widetilde{\chi}_1^0 l\bar{l}$) = 0.02, 0.07 and 0.17 (for each lepton flavor) for $D = 0$, 0.4 and 0.6, respectively. As shown in Fig. 1, for $D \lesssim 0.5$ this results in an appreciable enhancement of the dilepton signal (σ_2) compared to the conventional scenario. The clean trilepton signal (σ_{3l}) grows monotonically with increasing D.

Reducing D below zero at first leads to little change of BR($\widetilde{\chi}_2^0 \to \widetilde{\chi}_1^0 l\bar{l}$). Note that in the conventional scenario an increase of m_0 beyond 600 GeV results in an increase of this branching ratio, due to reduced destructive interference between Z and slepton exchange contributions. In our case reducing D below zero leads to lighter $SU(2)$ singlet sleptons, but heavier doublet sleptons; the two effects appear to largely cancel in the leptonic branching ratio of $\widetilde{\chi}_2^0$. However, for $D \lesssim -1.2$ the reduced $\tilde{b}_L$ mass leads to a rapid increase of $BR(\widetilde{\chi}_2^0 \to \widetilde{\chi}_1^0 b\bar{b})$, and a corresponding reduction of the leptonic branching ratio, as well as of σ_{3l}.

The chargino decays do not exhibit any remarkable feature over the entire region of D considered above, since they are dominated by virtual W exchange contributions.

The gluino branching ratios for $D = 0$ and 0.4 are:[†]
$BR(\tilde{g} \to \widetilde{\chi}_1^\pm q\bar{q'}) = 0.47,\ 0.37\ \ (qq' = ud,\ cs)$,
$BR(\tilde{g} \to \widetilde{\chi}_1^0 q\bar{q}) = 0.14,\ 0.35\ \ (q = u,\ d,\ c,\ s,\ b)$,
$BR(\tilde{g} \to \widetilde{\chi}_2^0 q\bar{q}) = 0.39,\ 0.28\ \ (q = u,\ d,\ c,\ s,\ b)$.

Note that the branching ratio for $\tilde{g} \to \widetilde{\chi}_1^0 q\bar{q}$ more than doubles when going from $D = 0$ to $D = 0.4$. This is a consequence of the reduced $\tilde{d}_R$ masses; recall that $SU(2)$ singlet squarks couple only very weakly to $\widetilde{\chi}_2^0$ and $\widetilde{\chi}_1^\pm$, which are predominantly $SU(2)$ gauginos. This leads to an increase of σ_0 (see Fig. 1), a hardening of the missing E_T spectrum, and an increase of the fraction of gluino pair events containing $b-$jets.

These changes become even more dramatic for $D \geq 0.55$. Now $\tilde{d}_R-$type squarks become lighter than the gluino, so that all gluinos decay to $\tilde{d}_R$s with $BR = 1$. Almost all the $\tilde{d}_R$'s in turn decay directly into the LSP. This suppresses σ_2 severely compared to the conventional case. In addition, $\tilde{g}\tilde{d}_R$ production as well as the pair production of $\tilde{d}_R-$type squarks lead to a further increase of σ_0, by up to a factor of 2. Recall that the clean trilepton signal (σ_{3l}) is also enhanced relative to the universal case. This combination of significant σ_0 and σ_{3l} but nearly vanishing σ_2 cannot be accommodated in the conventional scenario for any choice of the parameters.

As mentioned above, taking $D < 0$ decreases the masses of $SU(2)$ doublet as well as $\tilde{u}_R-$type squarks. This leads to increased destructive interference between $s-$ and $t-$channel contributions to $\tilde{g}$ pair production; as a result, the corresponding cross section drops by about a factor 1.6 when D is decreased from 0 to -0.9. We also saw that the decrease in squark mass is especially pronounced for $\tilde{b}_L$, whose mass is also reduced by RG–running from the GUT to the weak scale. Reducing D below 0 therefore also increases the fraction of $\tilde{g}$ decays that produce $b-$jets; we saw a similar effect for $D > 0$ due to the reduction of $m_{\tilde{b}_R}$.

[†]We ignore a small contribution from $\tilde{g} \to \widetilde{\chi}_1^\pm tb$. For our choice of gluino mass the corresponding partial width is suppressed by phase space. Moreover, its exact value depends sensitively on the masses and mixing angle of the stop squarks, which in turn depend on details of the GUT model, as discussed earlier.

However, while this fraction saturates at 1/3 for $D > 0$, where all three $\tilde{d}_R-$type squarks become light, it reaches 100% for $D < -0.9$, where $\tilde{g} \to \tilde{b}_L b$ is the gluino's only 2–body decay mode. Nearly all $\tilde{b}_L$ in turn decay into $\tilde{\chi}_2^0 b$. Since leptonic $\tilde{\chi}_2^0$ decays always lead to lepton pairs with opposite charges, the fraction of like–sign di–lepton pairs in σ_2 becomes very small here.

For $D \lesssim -1$, $\tilde{b}_L$ pair production begins to contribute significantly to σ_0 and σ_2. For example, for $D = -1.3$ (-1.4) we have $m_{\tilde{b}_L} = 174$ (115) GeV, giving $\sigma(\tilde{b}_L \tilde{b}_L^*) \simeq 1.0$ (8.5) pb. Due to the presence of (at least) two hard $b-$jets in the event, a dedicated sbottom pair search is expected to be sensitive to much smaller cross sections than more generic SUSY searches. A recent analysis [23] concludes that data from from Run II (TeV33) should allow to search for $\tilde{b}_L$ with mass up to at least 185 (205) GeV even if all $\tilde{\chi}_2^0$ decay into light quarks; in the case at hand this would cover the range $D \lesssim -1.25$. Scenarios with $D \lesssim -1.4$ can perhaps even be probed with existing data from Run I.[‡] Recall also that for large negative D, $BR(\tilde{\chi}_2^0 \to \tilde{\chi}_1^0 b\bar{b})$ is enhanced, leading to events with up to 6 $b-$jets! In the conventional scenario one would expect σ_0 and σ_2 to be dominated by sbottom pair production only at very large values of tanβ; and even there the $\tilde{b}_1 - \tilde{g}$ mass splitting is less than for large negative D in the non–universal case. Moreover, in the conventional scenario one also expects final states to be rich in $\tau-$leptons [14]; this need not be the case in the non–universal scenario.

It is clear from Table 1 and Figure 1 that if $m_{\tilde{q}} \gg m_{\tilde{g}}$, the latter being tightly constrained by LEP bound on $M_{\tilde{\chi}_1^\pm}$, the conventional scenario offers very little hope for a discovery in any of the three channels up to the TeV33 phase. In contrast, the non–universal scenario allows for cross sections in the striking range of Run II experiments, from sbottom pair production for large negative D, and in the tri–lepton channel at large positive D. Unfortunately observation of a signal in the tri–lepton channel by itself would not be sufficient to discriminate between the two scenarios. In the case at hand this distinction should be possible using TeV33 data, since there σ_0 and σ_{3l} become measurable over the entire range of D shown in Fig. 1. As noted earlier, the combination of large σ_0 and σ_{3l} with small σ_2 would exclude the conventional scenario. Additional clues can be obtained by correlating kinematical distributions, which are sensitive to sparticle masses, with measurements of the total cross sections.

The nonuniversal model allows for an observable trilepton signal even for higher $m_{\tilde{g}}$, albeit for much lighter sleptons ($D \gtrsim 0.6$). For example, with $D = 0.6$, $m_{1/2} = 130$ GeV ($m_{\tilde{g}}$ = 420 GeV), $m_{\tilde{l}_L} = 190$ GeV, $\sigma_{3l} = 0.09$ pb should be easily detectable already at Run II. The corresponding σ_0 is, however, unobservable even at TeV33.

We next consider region (B) of the parameter space. We shall work with $m_{1/2} = 105$ GeV and $m_0 = 300$ or 200 GeV. This is because even the phenomenology of the conventional scenario depends sensitively on m_0. Additional comments will be made on higher $m_{\tilde{g}}$ as and when appropriate. For $m_0 = 300$ GeV, the following mass patterns are obtained:

[‡] A search by the DØ collaboration [24] was able to exclude $m_{\tilde{b}_1} \leq 115$ GeV if $BR(\tilde{b}_1 \to \tilde{\chi}_1^0 b) \simeq 100\%$ and $m_{\tilde{\chi}_1^0} \lesssim 20$ GeV; neither of these conditions is satisfied in our case. However, during Run I CDF had significantly better $b-$tagging capabilities than the DØ detector.

D	0	0.5	-1.00	-1.75
$m_{\tilde{u}_L}$	392.6	421.5	327.4	268.3
$m_{\tilde{d}_R}$	389.9	285.6	541.2	631.3
$m_{\tilde{b}_L}$	339.9	372.3	261.1	181.6
$m_{\tilde{e}_L}$	312.2	164.3	488.2	586.5
$m_{\tilde{\nu}}$	303.9	147.9	482.9	582.1
$m_{\tilde{e}_R}$	305.3	341.6	215.0	105.0

Table 3: Squark and slepton masses (in GeV) at the weak scale for different values of D with $m_0 = 300$ GeV and $m_{1/2} = 105$ GeV.

The spectrum for the gaugino sector is: $m_{\tilde{g}} = 320$ GeV (the difference from the previous case arises due to radiative corrections to the gluino pole mass from squark–quark loops), $M_{\widetilde{\chi}_1^\pm} = 95.16$ GeV, $m_{\widetilde{\chi}_2^0} = 95.06$ GeV, and $m_{\widetilde{\chi}_1^0} = 46.71$ GeV.

Points to be noted are :

- In the conventional scenario the squarks are somewhat heavier than the $\tilde{g}$.
- Increasing D from zero again reduces $m_{\tilde{d}_R}$ and $m_{\tilde{l}_L}$; in particular, for $D \geq 0.5$, $m_{\tilde{d}_R} < m_{\tilde{g}}$ for all three generation of $\tilde{d}_R$−type squarks; and
- For large negative D ($D \leq -1$) the $\tilde{b}_L$ again becomes lighter than the gluino, but the theoretical lower bound on its mass is higher than in the previous case, because of the LEP bounds on $m_{\tilde{e}_R}$.

In contrast to the previous case, all possible squark and gluino production channels can contribute to σ_0 and σ_2. As shown in Fig. 2, the size of these cross sections will be typically larger than the conventional case even for moderate magnitudes of D irrespective of its sign. σ_0 still stays below the current bound of about 5 pb over the entire range of D; however, the region with $D \lesssim -1$ should be accessible in this channel at Run II, and the entire range of D can be probed at TeV33.

Unlike in the scenario with large m_0 we now observe a significant increase of the leptonic branching ratio of $\widetilde{\chi}_2^0$ for large negative D. Note that now slepton exchange contributions dominate the leptonic partial width of $\widetilde{\chi}_2^0$ even in the conventional scenario. Decreasing D away from zero therefore at first actually reduces the leptonic branching ratio of $\widetilde{\chi}_2^0$, due to the increase of $m_{\tilde{e}_L}$; note that for roughly equal mass the exchange of $SU(2)$ doublet sleptons contributes more, due to their stronger coupling to $\widetilde{\chi}_2^0$. However, for $D \lesssim -1$ the reduction of $m_{\tilde{e}_R}$ over–compensates the increase of $m_{\tilde{e}_L}$, leading to an increase of σ_{3l}; in fact, the region around $D = -1.75$ might already be excluded by existing searches [16] for tri–lepton events. For even more negative values of D, the decay $\widetilde{\chi}_1^\pm \to \tilde{\tau}_1 + \nu_\tau$ opens up. Since we do not include secondary leptons from τ decays when computing σ_{3l}, this cross section essentially vanishes once that decay mode dominates. In reality a measurable tri–lepton signal may survive, if softer cuts on the leptons are employed [17]; in addition, one may look for events with two leptons and a τ−jet [14]. Note that unlike $\widetilde{\chi}_1^\pm$, $\widetilde{\chi}_2^0$ also couples to $\tilde{e}_R$ and $\tilde{\mu}_R$; the leptonic branching ratio of $\widetilde{\chi}_2^0$ therefore remains large all the way out to the LEP–imposed lower bound on D.

Together with the reduction of $m_{\tilde{b}_L}$ this can lead to promising signals from events with at least one hard $b-$jet and two opposite–charge, same–flavor leptons; this might extend the reach of Run II to smaller values of $|D|$ than those accessible to sbottom searches based on final states without leptons [23]. Even these kinds of searches should be able to probe $D \lesssim -1.75$ (-1.6) at Run II (TeV33). Moreover, even a moderately negative D leads to a significant increase of σ_2. Recall that in this region all $SU(2)$ doublet squarks, as well as $\tilde{u}_R-$type squarks, are somewhat lighter than in the conventional scenario (for the same value of $m_{1/2}$ and m_0), leading to a significant increase of the total cross section for the production of strongly interacting sparticles. This over–compensates the modest decrease of the leptonic branching ratio of $\widetilde{\chi}_2^0$ for moderately negative D. Nevertheless this region of parameter space can be probed in the inclusive di–lepton channel only at TeV33.

Increasing D from zero leads to quickly increasing leptonic branching ratio of $\widetilde{\chi}_2^0$, and a corresponding rise of σ_{3l}; as for negative D, the region just below the largest value of $|D|$ allowed by slepton searches at LEP is again at best marginally consistent with current search limits for tri–lepton final states. For the largest LEP–allowed value of D, the decays $\widetilde{\chi}_2^0 \to \tilde{\nu}\nu$ open up. The sneutrino then decays into the invisible channel $\tilde{\nu} \to \widetilde{\chi}_1^0\nu$ (the so–called VLSP scenario [25]), leading to a vanishing σ_{3l}. Apart from the most extreme regions, the entire range of D can easily be covered in the tri–lepton mode at TeV33, and much of it should already be accessible to Run II experiments. Finally, sizably positive D again suppresses σ_2, first due to the dominance of $\tilde{g} \to \tilde{d}_R d$ decays (for $D \geq 0.35$), and, for the largest allowed value of D, due to $\widetilde{\chi}_2^0 \to \tilde{\nu}\nu$ decays.

For $m_0 = 200$ we obtain the following sfermion spectra:

D	0	0.4	0.6	-0.75	-1.75
$m_{\tilde{u}_L}$	326.6	339.1	345.2	301.7	264.4
$m_{\tilde{d}_R}$	323.4	282.0	258.8	389.3	462.7
$m_{\tilde{b}_L}$	293.6	307.5	314.2	265.6	222.9
$m_{\tilde{e}_L}$	218.0	149.9	99.7	307.4	396.3
$m_{\tilde{e}_R}$	207.8	227.0	236.1	166.0	81.7
$m_{\tilde{\nu}}$	205.9	131.7	69.5	298.9	389.8

Table 4: Squark and slepton masses (in GeV) at the weak scale for different values of D, for $m_0 = 200$ GeV and $m_{1/2} = 105$ GeV.

The gaugino masses are nearly the same as for $m_0 = 300$ GeV. The behavior of our three cross sections, shown in Fig. 3, is qualitatively similar to the previous case. However, due to the reduced importance of the $D-$terms as compared to the universal contribution from gluino–quark loops, even in the non–universal scenario the mass splitting between different first or second generation squarks now amounts to at most 50%. As a consequence the changes of the three cross sections σ_0, σ_2 and σ_{3l} with varying D are somewhat smaller than for larger values of m_0. In particular, σ_0 now increases by less than a factor of 3 as D is decreased from zero to its most negative allowed value. On the other hand, owing to the reduction of the overall squark mass scale, the entire range of D can now be probed in the jets plus missing E_T channel at Run II of the Tevatron.

Somewhat larger variations can occur for the leptonic cross sections. Since sleptons are still significantly heavier than $\tilde{\chi}_2^0$ and $\tilde{\chi}_1^\pm$ in the conventional scenario ($D = 0$), the leptonic branching ratios can change considerably when D is allowed to take values close to the present upper or lower bounds. Due to the reduction of the slepton masses compared to the previous cases, these cross sections are also generally larger than for scenarios with larger m_0. In particular, the tri–lepton signal should be detectable over the entire range of D at Run II. The availability of at least two distinct SUSY signals should make it easier to distinguish between the universal and non–universal scenarios; on the other hand, the expected differences in the squark spectrum are smaller than before, as already noted.

Yet smaller values of m_0 continue the trend of going from 300 to 200 GeV. The maximal possible mass splitting between squarks keeps decreasing even in the non–universal scenario. Significant differences in the slepton spectrum remain possible; as long as sleptons are heavier than $SU(2)$ gauginos in the conventional scenario, these changes of slepton masses affect the leptonic branching ratios of electroweak gauginos. However, for $m_0 \leq 85$ GeV, these gauginos can decay into on–shell sleptons even in the universal scenario; their leptonic branching ratios are then much less sensitive to D. Moreover, LEP searches severely limit the maximal allowed value of D for such small values of m_0.

4 Summary and conclusions

In this paper we have studied the impact of $SO(10)$ $D-$term contributions to scalar masses on the search for supersymmetric particles at the Tevatron collider. These contributions have opposite signs for sfermions residing in the $\mathbf{5}^*$ representation of $SU(5)$ and those residing in the $\mathbf{10}$ representation. Any non–zero value of D will therefore reduce the masses of some sfermions, while making others heavier. This usually results in an increase of the total cross section for sparticle production at hadron colliders. Heavy sparticles simply decouple, so that further increases of their masses have little impact on any observables; however, the cross section for the production of the lighter sfermions keeps increasing as their masses are reduced. The only exception to this rule occurs for large m_0 and moderately negative D, where the reduction of most squark masses actually reduces the gluino pair cross section through interference effects. However, we found that this can reduce the total cross section for the production of strongly interacting sparticles by no more than a factor of 2; in contrast, for extreme (but allowed!) values of D this cross section can be enhanced by more than an order of magnitude compared to models with degenerate squarks.

Similarly, non–degeneracy in the slepton sector tends to increase leptonic branching ratios of electroweak gauginos, in particular those of $\tilde{\chi}_2^0$, leading to an increase of the relatively hadron–free tri–lepton cross section from $\tilde{\chi}_2^0\tilde{\chi}_1^\pm$ production. For large negative D this cross section can be sufficiently large to yield unambiguous signals at the next run of the Tevatron collider even if the *average* sfermion mass lies well above the gluino mass, which in turn lies well above its lower limit derived from chargino searches at LEP (assuming gaugino mass unification). However, if D is very close to its LEP–imposed upper or lower limit, the tri–lepton signal is greatly depleted, since there decays $\widetilde{\chi}_2^0 \to \widetilde{\chi}_1^0 \tilde{\nu}\nu$ and $\widetilde{\chi}_1^\pm \to \tilde{\tau}_1 \nu_\tau$, respectively, become possible.

One interesting feature of the non–universal scenario is that one can find combinations of parameters where the tri–lepton and multi–jet plus missing E_T signals are both large, but the jets plus di–lepton plus missing E_T signal is very small. This last signal results from cascade decays of the gluino and $SU(2)$ doublet squarks, which are suppressed for $D > 0$, where $\tilde{d}_R-$type squarks can be the only light squarks. Such a combination of signals cannot occur if all sfermions are degenerate at the GUT scale.

Models with non–vanishing D also introduce new signals, which do not hold much promise if squarks are approximately degenerate. These signals originate from the copious production of light sbottom squarks. In models with small or moderate tanβ studied here, light $\tilde{b}_L$ squarks can occur for negative D. Since $\tilde{b}_L$ dominantly decays into $\widetilde{\chi}_2^0 b$, $\tilde{b}_L$ pair production (as well as gluino production followed by $\tilde{g} \to \tilde{b}_L b$ decays) gives rise to several signals. Since for large negative D the branching ratio for $\widetilde{\chi}_2^0 \to \widetilde{\chi}_1^0 b\bar{b}$ decays can be enhanced, searches for final states containing 4, or even 6, $b-$jets may offer the best reach. For somewhat smaller $|D|$ or smaller m_0, searches for final states containing two hard $b-$jets and two opposite–charged, same–flavor leptons might hold even more promise. This reinforces the conclusions of ref.[22] regarding the importance of SUSY signals containing hard $b-$jets.

The maximal possible mass splitting in the squark sector becomes smaller for smaller ratios of the GUT scale parameters $m_0/m_{1/2}$, due to the increased importance of flavor–universal contributions to squark masses from gluino–quark loops. For $m_0 < 2m_{1/2}$ the main impact of sfermion non–universality on SUSY signals at the Tevatron comes from changes of the leptonic branching ratios of electroweak gauginos, which remain sensitive to D as long as $m_0 \gtrsim m_{1/2}$. For even smaller values of $m_0/m_{1/2}$ evidence for non–universality in the sfermion sector can probably best be found in studies of slepton pair production. However, these may have to be postponed until a lepton supercollider has been built.

In summary, in $SO(10)$ models with sizable $D-$term contributions to scalar masses, SUSY searches at future collider runs of the Tevatron hold even more promise than in the conventional mSUGRA model. New regions of the $(m_0, m_{1/2})$ plane become accessible to "standard" searches involving jets and/or leptons plus missing transverse momentum, and searches for new final states containing hard $b-$jets may increase the reach even further. Under favorable circumstances the pattern of observed signals will allow to discriminate between universal and non–universal scenarios. This once again underscores the importance of searching for supersymmetry in as many different channels as possible.

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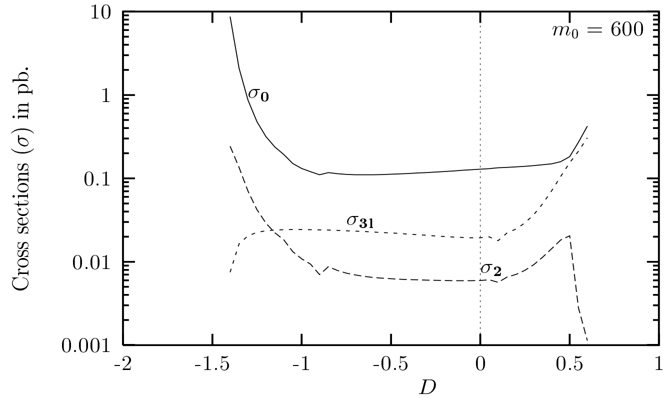

Figure 1: Cross sections (in pb, without any cut) for different final states at the Tevatron collider with $\sqrt{s} = 2$ TeV as a function of D. The other relevant parameters are $m_{1/2} = 105$ GeV, $m_0 = 600$ GeV and tan$\beta = 3$, while μ is fixed by imposing the radiative electroweak symmetry breaking condition for $D = 0$.

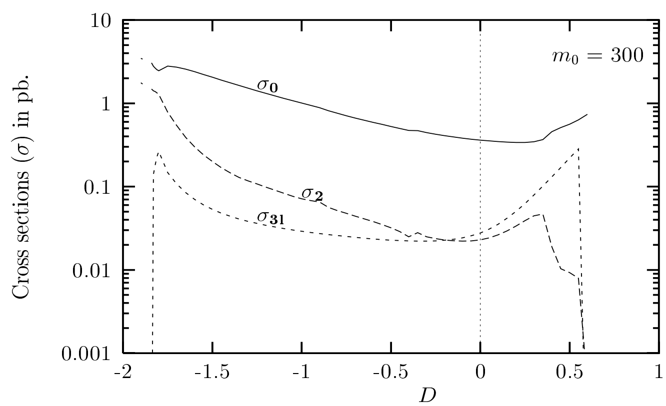

Figure 2: As in Fig. 1, but with $m_0 = 300$ GeV.

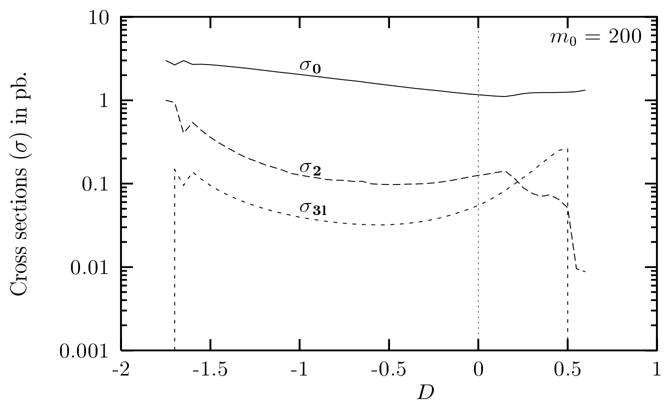

Figure 3: As in Fig. 1, but with $m_0 = 200$ GeV.