Beyond the Standard Model physics at the Tevatron

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Abstract.

The purpose of this lecture is to provide an overview of the analyses currently pursued by the Tevatron collaborations in the search for processes beyond the Standard Model. The focus has been put on Supersymmetry and – wherever possible – a pedagogical approach has been followed in presenting the analyses. For each considered process, except for combined results, only one analysis – either by CDF or DØ – has been discussed.

The lecture has been divided in three main sections: phenomenology of the Higgs sector of the Minimal Supersymmetric Standard Model and Supersymmetry with conserved and violated $R$-parity. The layouts of the Fermilab accelerator complex and of the two Tevatron detectors are not discussed in the following; appropriate descriptions can be found in the references [1, 2, 3].

1. Introduction

Despite the general agreement observed between the experimental data and theoretical predictions, there are reasons to believe that the current theory describing the building blocks of Nature and their interactions – the Standard Model (SM) – is not final. Indeed, it is a fact that the SM does not include any treatment of gravitational interactions, which makes its description of the fundamental forces clearly incomplete; besides, some of the SM couplings are not asymptotically free, a fact that endangers the consistency of the model as a formal quantum field theory. For these reasons, despite its many successes, the SM should be considered no more than an effective low-energy theory valid up to some cut-off scale $\Lambda$.

The issue that modern high energy physics has to address in the near future is then two-fold: determine the magnitude of the cut-off $\Lambda$ at which the current model breaks down and new physics arises and provide a suitable description of the more fundamental theory extending the SM beyond that cut-off. Among all the possibilities, $\Lambda \sim O(1 \text{ TeV})$ seems to be most appealing from a theoretical point of view, the reason relying mainly on the so-called naturalness problem affecting the SM [4]: this is related to the difficulty of reconciling the value of the Higgs boson mass, $M_H$, with the corresponding quantum corrections, quadratically divergent in $M_H$.

In this scenario, one of the most successful ideas comes from low-energy Supersymmetry (SUSY), which – by extending the symmetry of the SM to the spin sector – is capable of solving the naturalness problem by providing an accurate cancellation (at all perturbative orders) of the quadratic divergences affecting the masses of the spin-0 bosons of the model. The drawback is the introduction of a large number of new particles, a sort of mirror-image of the SM spectrum induced by the spin symmetry. Besides being a viable solution of the naturalness problem, SUSY represents a step forward toward the unification of all interactions: not only does it allow the
unification of the electroweak and strong coupling constants (a fact that is not realized within the SM framework), but also seems to play an important role for the consistency of superstrings – candidate unified theories of all interactions, gravitational included.

Despite all its desirable features, SUSY cannot be realized in Nature as an exact symmetry; in other words, if supersymmetry exists, it must be broken. If this were not the case, then, for each SM particle discovered so far, a SUSY partner of equal mass should exist; conversely, no such particle has ever been observed. Nevertheless, the idea of supersymmetry can be rescued by assuming that a mass splitting is introduced by the SUSY-breaking mechanism in such a way that the masses of the unobserved SUSY particles are pushed beyond the current experimental limits. The possibility of choosing between different SUSY-breaking scenarios complicates the emerging phenomenological picture: this is a consequence of the fact that different SUSY-breaking mechanisms lead to different SUSY structures and therefore phenomenology. In order to parametrize the phenomenology at the electroweak scale without relying on a particular scenario, a general supersymmetric SM extension, named Minimal Supersymmetric Standard Model (MSSM), is considered; in the MSSM Lagrangian, supersymmetry is explicitly violated by the introduction of soft terms, which – despite the symmetry breaking – preserve the good ultraviolet properties of the theory. As a result, however, the MSSM is characterized by even more free parameters than the SM; for this reason, experimental results are often interpreted in terms of models constrained on the basis of specific SUSY-breaking scenarios and Grand Unified Theory (GUT) scale relations or in terms of special benchmarking models.

2. The Minimal Supersymmetric Standard Model Higgs sector

Contrary to the SM, which is characterized by one SU(2) Higgs doublet, in the MSSM two Higgs doublets are required in order to give mass to all quarks and leptons, to cancel gauge anomalies and to avoid massless charged fermions [5]. These give rise to five physical states, namely two CP-even and one CP-odd neutral states (h, H and A respectively) and two charged Higgs bosons (H±). Spectrum and phenomenology of the Higgs sector of the MSSM at tree-level is completely controlled by two parameters, conventionally chosen in $M_A$ (the mass of the pseudo-scalar Higgs boson) and $\tan \beta$ (the ratio of the vacuum expectation values of the up to down-type neutral Higgs field); quantum corrections entering at higher order are dominated by loops involving quarks and scalar quarks (also known as squarks) of the third generation as a consequence of their large Yukawa couplings.

The spectrum of the MSSM Higgs sector is characterized by a striking peculiarity: the mass of the lightest CP-even Higgs state (h) is constrained by an upper bound, $M_h^{max}$; this is in contrast with the SM, where the mass of the Higgs boson is a free parameter of the theory. At tree-level, the prediction for $M_h$ corresponds to its theoretical upper bound, $M_h^{max} = M_Z$; once quantum corrections are included, $M_h^{max}$ is substantially enhanced as a consequence of an incomplete cancellation between top quark and top squark loops. A residual dependence is found on the degree of mixing affecting the top squark\(^1\), with $M_h^{max}$ reaching a maximum for maximal \(\tilde{t}\)-mixing. Another interesting feature of the spectrum of the MSSM Higgs sector concerns the high $\tan \beta$ regime, where either of the CP-even Higgs bosons is predicted to be nearly mass-degenerate with the CP-odd state.

\[ m_{\tilde{t}} (A_{\tilde{t}} - \mu \cot \beta) \]
\[ m_{\tilde{d}} (A_{\tilde{d}} - \mu \tan \beta) \]

where $m_t$ and $A_t$ are the mass of the corresponding SM fermion partner and the trilinear coupling. These terms are dominant for third-generation fermions thanks to the large mass of their corresponding SM partners.

\(^1\) Supersymmetric partners of both $L$ and $R$ chirality states are introduced for each SM fermion; in this way, two different sfermions, $\tilde{f}_L$ and $\tilde{f}_R$, are defined. The off-diagonal elements in their $2 \times 2$ mass matrix are:
Decays of the MSSM neutral Higgs bosons are essentially driven by the Yukawa coupling to down-type fermions (b¯b and τ+τ−): playing a dominant role in most of the parameter space, it is enhanced by large values of tan β. As a result, branching ratios of neutral Higgs bosons (Φ=h,H,A) decaying into b¯b or τ+τ− pairs account to almost 90 and 10% respectively, regardless of the Higgs mass.

2.1. Searching the Standard Model Higgs
In the previous paragraphs, it was noted that there is an upper bound on the value of the lightest CP-even neutral Higgs boson of the MSSM. Based on accurate loop calculations, it can be shown that, assuming an average squark mass parameter $M_S \sim \mathcal{O}(\text{TeV}/c^2)$, $M_{h_{\text{max}}}^2$ ranges between 120 and 135 GeV/c^2 according to the degree of ˜t-mixing and the value of the top mass [6]; higher values of this constraint (up to 150 GeV/c^2) are indeed possible, although being discouraged by theoretical arguments. This suggests that a significant region of the MSSM parameter space can be probed by searching the SM Higgs.

![Figure 1. Comparison between Standard Model predictions and experimental sensitivities obtained by the Tevatron collaborations on the Higgs boson searches exploiting 194–950 pb$^{-1}$ of Run II data. The experimental sensitivities are divided by the product of cross-section $\sigma(p\bar{p} \rightarrow H_{\text{SM}}, VH_{\text{SM}})$ ($V = W, Z$) times branching ratio for different final states as predicted by the SM is plotted with respect to the SM Higgs boson mass.](image)

A simple example can be considered in order to understand how this can be achieved. Consider the MSSM in a particularly simple scenario, the so-called decoupling regime: when $M_A \gg M_Z$, the heavy CP-even, the CP-odd and the charged Higgs decouple from the low energy effective theory and the lightest CP-even state behaves as the SM Higgs boson. If this is the case, the results obtained for the SM (for instance, a 95% confidence level exclusion of a certain Higgs mass range) can be applied directly to the MSSM. In other regions, the comparison between SM and MSSM Higgs bosons is less trivial; nevertheless, in this case the similarities between the SM and the CP-even MSSM Higgs bosons can be exploited. In this way, the results of the searches for the SM Higgs bosons (see figure 1 and [7] for the current status of SM Higgs searches at the Tevatron) can be reinterpreted in the MSSM framework. At the end of the procedure, the limits on the SM Higgs boson mass can be mapped onto the $(M_A, \tan \beta)$ plane [5]; the results of the unfolding process is illustrated in figure 2. The effect of the upper bound on $M_H$ results in wide exclusions of the MSSM parameter space already with an integrated luminosity of 5 fb$^{-1}$; the
extent of the excluded regions, however, depends on the specific benchmarking model used for tuning the remaining MSSM parameters. Note that, while in the high \( \tan \beta \) regime the two CP-even Higgs bosons are nearly mass-degenerate and therefore both contribute to the same signal mass range, in the moderate \( \tan \beta \) region where \( M_A \approx M_{\text{max}}^h \) this is not true any longer, giving rise to a problematic region which is visible in both plots shown in figure 2. Other problematic regions, on the other hand, can be the result of conspiring parameters within a specific scenario and can be explored by means of redundant analyses that rely on different search channels.

![Graphical representation of exclusion contours](image)

**Figure 2.** Expected 95\% confidence level exclusion contours in the MSSM (\( M_A, \tan \beta \)) plane; results are based on fast detector simulations of the Higgsstrahlung process (\( p\bar{p} \rightarrow VH_{\text{SM}}, \text{with } V=W, Z \)) with \( H_{\text{SM}} \rightarrow b\bar{b} \) and are expressed in terms of integrated luminosity needed per experiment in order to achieve the exclusion [5]. Left: maximal stop-mixing scenario. Right: scenario with suppressed \( \Phi b\bar{b} \) coupling (\( \Phi=h, H \)).

### 2.2. Higgs bosons decaying in tau pairs

The gluon-fusion processes are mediated by triangle loops involving the heavy quarks and squarks belonging to the third generation. Gluon-fusion is the dominant MSSM neutral Higgs boson production mechanism at the Tevatron, with a cross-section ranging from 0.03 to 30 pb with increasing \( \tan \beta \) for an assumed Higgs boson mass between 100 and 200 GeV/\( c^2 \). However, as in the case of the SM Higgs boson, the dominant \( b\bar{b} \) final states are overwhelmed by the copious QCD background. This essentially rules out the possibility of an analysis targeting the leading \( p\bar{p} \rightarrow \Phi \rightarrow b\bar{b} \) (with \( \Phi=h/H, A \)); at large \( \tan \beta \), however, the next-to-leading \( \Phi \rightarrow \tau^+\tau^- \) decay modes become promising.

Tau leptons are not as easily identified as electrons or muons. The CDF analysis [8] searches for a signature obtained by the combination of one tau decaying leptonically with the other tau decaying hadronically\(^3\). Tau leptons decaying hadronically are reconstructed as narrow, pencil-like jets with low track and \( \pi^0 \) multiplicity that can be distinguished from the broader jets originating from quark and gluon showers. In order to identify hadronic tau decays, CDF relies on a double-cone algorithm (see figure 3). The algorithm is seeded by a charged track

\(^2\) Such is the case, for instance, of the region visible in the high \( \tan \beta \) end of the suppressed \( V\Phi \rightarrow Vb\bar{b} \) scenario: appropriate coverage in this region can be achieved by considering the \( p\bar{p} \rightarrow \Phi b\bar{b} \rightarrow b\bar{b}b\bar{b} \) search channel (see §2.3 and figure 5).

\(^3\) A similar analysis has been performed by DØ [9] based on an integrated luminosity of 325 pb\(^{-1}\).
with $p_T > 6 \text{ GeV}/c$ pointing to an energy deposit $E_{cl}$ in the calorimeter. Then additional tracks and neutral pions are associated with the seed to form the $\tau$ candidate; the positions of $\pi^0$’s are determined from the signals released in the shower maximum detector\textsuperscript{4} and their energy is measured by the electromagnetic calorimeter. The direction of the seed track defines the axis of the signal cone and an isolation annulus; the semi-aperture of the signal cone ranges between 10 and 30° (corresponding to the outer semi-aperture of the isolation annulus) according to the law $5 \text{ GeV}/E_{cl}$. The sums of the transverse momenta of tracks and of the transverse energies of $\pi^0$’s in the isolation annulus must not exceed $1 \text{ GeV}/c$ and $1 \text{ GeV}$ respectively. The four-momentum of the tau candidate is then computed from all tracks and $\pi^0$’s falling within the signal cone; tau candidates with $p_T \leq 15 \text{ GeV}/c$ or $m \geq 1.8 \text{ GeV}/c^2$ are rejected.

Figure 3. In CDF tau leptons are identified as isolated pencil-like jets.

An obvious variable that can be used for increasing the purity of the tau reconstruction comes from the charged track multiplicity, $N^\text{trk}_{\text{sig}}$, defined as the number of charged tracks with $p_T > 1 \text{ GeV}/c$ associated with the $\tau$ candidate; $N^\text{trk}_{\text{sig}}$ is restricted to one or three, consistently with the dominant $\tau$ decay modes (see figure 4); furthermore, the charges of tracks falling in the $N^\text{trk}_{\text{sig}} = 3$ bin need to sum up to ±1. The algorithm defined in this way achieves an efficiency of 46% with a misidentification rate – estimated from data – ranging from 1.5 to 0.1% per jet for jet transverse energies between 20 and 100 GeV.

The dominant background for this search is given by the irreducible Drell-Yan (DY) production of $Z/\gamma^*$ with subsequent decay in tau pairs; it has been estimated based on Monte Carlo simulations with a normalization corresponding to the measured value of $\sigma(p\bar{p} \rightarrow Z/\gamma^*) \times BR(Z/\gamma^* \rightarrow \ell^+\ell^-)$ in the Z di-lepton mass window\textsuperscript{5}. QCD di-jet and multijet events, $W$ produced in association with jets and photon plus jets events represent the next-to-leading background sources; their contribution is estimated by applying the jet → $\tau$ misidentification rates to jets in events passing all selection criteria except the hadronic tau reconstruction.

The data sample used in this analysis has been collected by means of a highly efficient trigger – especially designed to target tau production – which requires an electron or muon and an isolated track. Events are then filtered by requiring one high-$p_T$ ($p_T > 10 \text{ GeV}/c$) electron or muon and one tau lepton decaying hadronically with opposite charge. Events with $Y_T \equiv |p_T^f| + |p_T^\tau| + |E_T| < 50 \text{ GeV}$ are rejected in order to reduce the contamination due to low-energy multijet events; $W$ production in association with jets, with a jet misidentified as a tau

\textsuperscript{4} The shower maximum detector – a wire/strip chamber or a scintillator tile array – is embedded in the calorimeter at a depth corresponding to the maximum development of the electromagnetic shower.

\textsuperscript{5} $\sigma(p\bar{p} \rightarrow Z/\gamma^*) \times BR(Z/\gamma^* \rightarrow \ell^+\ell^-) = 254.9 \text{ pb}$ with $\ell = e, \mu$, CDF measurement [10].
lepton, are suppressed by a topological cut acting on the relative directions of the reconstructed tau candidate and the missing $E_T$.

No excess of events is observed in the first 310 pb$^{-1}$ of Run II data; the uncertainties in the background estimation – dominated by the systematic effects underlying the determination of the jet → τ misidentification rates and of the luminosity of the data sample – are comparable or larger than the expected signal for a large variety of MSSM scenarios. In order to increase the sensitivity of the search, a binned likelihood fit has been performed in terms of the visible mass:

$$m_{\text{vis}} \equiv \sqrt{(p_\ell^2 + p_\tau^2 + P_T^2) \cdot (p_\ell^2 + p_\tau^2 + P_T^2)}$$

where $P_T = (E_x, E_y, 0, E_T)$. In the hypothesis of absence of signal, limits have been set on $\sigma(pp→Φ) \times BR(Φ→τ^+τ^-)$ ($Φ = h/H, A$) for different values of $\tan \beta$ and $M_A$; the results are summarized in figure 5 in the ($M_A, \tan \beta$) plane and interpreted in terms of two MSSM scenarios.

### 2.3. Higgs bosons in multijet environments

In the MSSM, the radiation of a neutral Higgs boson off bottom quarks becomes important in the high $\tan \beta$ regime, with cross-sections of the order of 1 to 0.1 pb corresponding to a Higgs mass range of 100 to 200 GeV/c^2. In this case, it is possible to consider the production process in conjunction with the decay of the Higgs boson into a bb pair, thanks to the additional b quarks that provide a powerful signature to keep the QCD background under control. Since for large $\tan \beta$ values either of the CP-even Higgs bosons is mass-degenerate with the CP-odd state, and since their width is small if compared to the di-jet experimental resolution, the h/H system cannot be resolved from A: therefore, under this assumption, the total cross-section for signal production is assumed to be twice that predicted for the A alone.

The DØ analysis [11] in the four b-jet channel requires at least three jets$^6$ in each event being originated from b quarks. It is a known fact that B hadrons are massive objects characterized by the following properties: 

There are several algorithm for jet reconstruction: the most used identifies jets with cones of fixed semi-aperture (defined as $ΔR \equiv \sqrt{Δ\eta^2 + Δφ^2}$ – where $φ$ is the azimuthal angle and $η \equiv -\ln\tan(θ/2)$, $θ$ being the polar angle with respect to the proton beam direction, the pseudorapidity). The value chosen for $ΔR$ needs to be chosen according to the working environment: in multijet events a cone size of 0.5 and 0.4 is commonly used by DØ and CDF respectively.

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**Figure 4.** Search for $pp→Φ→τ^+τ^-$ ($Φ = h/H, A$). Left: charged track multiplicity of tau candidates (before the opposite sign and $N_{\text{trk}}^{\text{sig}} = 1, 3$ requirements) compared to main backgrounds. Right: visible mass spectrum for signal and background. Shown plots by CDF with 310 pb$^{-1}$ [8].
by relatively long lifetimes (of the order of 1.5 ps): in the fragmentation of a b quark, when they acquire large momenta, this results in decay lengths of the order of $c\tau \approx 300$–500 $\mu$m, which become detectable by means of vertex detectors. This fact is exploited in the b-tagging technique that relies on the explicit reconstruction of a secondary vertex: a jet is identified as a b-jet if a secondary vertex can be found within a distance $\Delta R < 0.5$ from its axis and if the transverse displacement of the secondary vertex from the primary interaction vertex exceeds five times the uncertainty on the displacement itself. Jets are b-tagged up to $|\eta| < 2.5$, although the algorithm is about twice as efficient in the central region of the detector ($|\eta| < 1.1$) due to tracking coverage. The b-tagging efficiency reaches 55% for central high-$p_T$ jets $p_T > 35$ GeV/$c$, with a light quark (or gluon) tag rate of about 1%; thanks to its very low misidentification rate, which translates in a high background rejection, it is particularly indicated in multijet environments. The uncertainty on the b-tagging efficiency represents the dominant systematic effect in this analysis, followed by the uncertainty on the jet energy scale.

7 Vertices are reconstructed on the basis of the outgoing tracks that are identified by the central and inner tracking devices of a detector. Typical procedures require a minimum of three quality tracks in order to identify a secondary vertex; the number of tracks can be relaxed to two provided their quality and number of associated hits in the inner tracking devices is sufficiently high.

8 An alternative procedure for matching secondary vertices to jets prescribes the tracks to be used in the secondary vertex reconstruction to be embedded in the jet cone.
Figure 6. DØ 260 pb$^{-1}$ search for $p\bar{p} \to \Phi b\bar{b} \to b\bar{b}b\bar{b}$ ($\Phi = h/H/A$). Left: invariant mass spectrum of the two leading-$p_T$ jets in the doubly b-tagged data compared to a sum of backgrounds; QCD contributions from heavy flavour and mis-tagged events are shown separately, other backgrounds – $Z(\to b\bar{b})$+jets, $Zb$, $tt$ – are summed together. Right: same distribution for events with at least three b-tagged jets. Data compared to total background and to the $M_A=120\text{ GeV}/c^2$ signal that can be excluded at 95% of confidence level [11].

The inclusive three-jet data sample, corresponding to an integrated luminosity of 260 pb$^{-1}$, are filtered by requiring at least three jets being b-tagged according to a secondary vertex reconstruction procedure; the leading jet needs to have $E_T > 35\text{ GeV}$, while for the third jet the same request is lowered to 15 GeV. The estimation of the QCD background is performed directly from data: first, the mis-tag functions are determined by measuring the probability to b-tag a jet on the initial data sample (before applying any b-tagging requirement); then, any true heavy flavour contamination is subtracted and the mis-tag functions corrected. Finally, the mis-tag functions are applied to all the untagged jets within the doubly-tagged sample in order to get the background for triply b-tagged events. The overall background normalization scale is determined by a fit of the leading di-jet invariant mass spectrum outside the signal region in the sample with three b-tags. Figure 6 illustrates the level of agreement achieved by this procedure.

Figure 7. DØ 260 pb$^{-1}$ search for $p\bar{p} \to \Phi b\bar{b} \to b\bar{b}b\bar{b}$ ($\Phi = h/H/A$): expected and observed 95% confidence level upper limits on $\sigma(p\bar{p} \to \Phi b\bar{b}) \times BR(\Phi \to b\bar{b})$. The shaded band indicates the $\pm 1\sigma$ range on the expected limit. The dotted line represents the signal cross-section for $\tan \beta = 80$ in the no-mixing scenario; the overlaid band reflects theoretical uncertainty [11].

A modified frequentist method applied to the di-jet invariant mass distribution in triply b-tagged events is used to set limits on the $\sigma(p\bar{p} \to \Phi b\bar{b}) \times BR(\Phi \to b\bar{b})$; the value of $\tan \beta$ is varied until the confidence level for signal is less than 5%. Figure 6 shows the result of this procedure, depicting the signal distribution for a 120 GeV$/c^2$ Higgs boson $A$ at the exclusion limit; signal sensitivity for a fixed $\tan \beta = 80$ is instead shown in figure 7. By varying the $\tan \beta$ and $M_A$ values simultaneously, the limit is most suitably represented on the $(M_A, \tan \beta)$ plane: the result has
already been shown in figure 5.

2.4. Search for charged Higgs bosons

At the Tevatron, direct pair production of charged Higgs bosons is expected to have a cross section of the order of 0.1 pb; the t̄t production, on the other hand, with a SM expected cross-section of about 7 pb, may offer another source of charged Higgs production. In this case – assuming the reaction is kinematically allowed – a top quark could decay to a H+i b pair, which would then compete with the dominant SM decay t→W+i b. This mechanism might then provide a larger production of charged Higgs bosons with a cleaner signature than direct pair production.

Figure 8. CDF search for a charged Higgs boson in the top quark decay products with 193 pb−1 [12]. Left: 95% confidence level upper limit on the branching ratio for the decay t→H+i b. Right: regions of the (M_{H+i}, tan β) plane that have been excluded at 95% confidence level by this search for a particular scenario.

CDF has searched for charged Higgs production in the decay products of t̄t events using 193 pb−1 of Run II data [12]; the search takes in consideration all the accessible decay modes of a light9 charged Higgs bosons in SM particles10 – namely H+i→τντ, H+i→c̄s, H+i→t̄b (→W+i b̄b) and H+i→W+iΦ (→W+i b̄b) – and investigates all the signatures that can eventually emerge from the decay of a t̄t system in the assumption that such a Higgs boson exists. t̄t events are classified in four exclusive categories: di-lepton11, lepton plus jets with one b-tag, lepton plus jets with two or more b-tags and lepton in association with a tau decaying hadronically. A likelihood function of M_{H+i} and tan β is then used to compare the number of observed events with the number of expected events in each category – being clear that, depending on the top and charged Higgs branching ratios (affected in turn by the values assumed by the likelihood variables M_{H+i} and tan β), this number can show discrepancies with respect to the corresponding SM prediction.

A Bayesian approach has been used to set limits in the hypothesis of no signal; the results are summarized in figure 8, where the upper limit on BR(t→H+i b) as a function of M_{H+i} is shown together with the excluded region in the (tan β, M_{H+i}) plane corresponding to one of the benchmark models considered by the analysis.

9 M_{H+i} < M_t – M_b
10 Although the possibility of a charged Higgs boson decaying into SUSY particles, such as a chargino-neutralino pair, is not formally excluded, the area of the MSSM parameter space in which these decays are kinematically allowed is not large and will be neglected in the following.
11 The term lepton here refers to electrons and muons only.
3. Supersymmetry with conserved $R$–parity

It is a known fact that the total baryon number ($B$) and the individual lepton numbers ($L_e, L_\mu, L_\tau$) are accidental symmetries of the SM, meaning that they are not violated by any of the gauge-invariant Lagrangian terms. In other words, in the SM the total number of baryons and the number of leptons of each individual flavour happen to be conserved without imposing any symmetry other than the gauge SU(3)$_C \times$ SU(2)$_L \times$ U(1)$_Y$.

In the construction of the MSSM, however, the most general superpotential that can be written under the assumptions of gauge and SUSY invariance and renormalizability is the following:

$$w = h_{ab}^U Q_a U_b^C H_2 + h_{ab}^D Q_a D_b^C H_1 + h_{ab}^L L_a E_b^C H_1 + \mu H_1 H_2 + \lambda_{abc} Q_a D_b^C L_c + \lambda'_{abc} L_a H_b^C E_c + \mu'_{abc} L_a H_2 + \lambda''_{abc} U_a^C D_b^C D_c^C,$$

where $a, b$ and $c$ are generation indices. It can be noted that, while the first line of equation (2) contains only terms conserving both $B$ and $L$, terms in the second and third lines violate lepton ($\Delta B = 0, |\Delta L| = 1$) and baryon numbers ($|\Delta B| = 1, \Delta L = 0$) respectively. As a result, the simultaneous presence of these terms in the superpotential would lead to potentially catastrophic situations, such as, for instance, a squark-mediated fast proton decay.

A common way out from this situation is to suppress the potentially dangerous terms by means of an additional discrete, multiplicative symmetry: $R$–parity, defined as:

$$R \equiv (-1)^{2S + 3B + L},$$

where $S$ is the spin quantum number; in practice, $R = 1$ for all standard particles (quarks, leptons, gauge and Higgs bosons), while $R = -1$ labels their supersymmetric partners (squarks, sleptons, gauginos and higgsinos).

Two additional features characterize the phenomenology of supersymmetric models with $R$–parity conservation: first, SUSY particles can only be produced in pairs; second, the lightest supersymmetric particle (LSP) needs to be stable. No evidence for such a particle has been found so far, which leads to the hypothesis of a neutral and weakly interacting LSP

$^{12}$For this reason, the LSP is considered an eligible candidate for cold dark matter.

$^{13}$The naming scheme is inherited from the fact that in the mSUGRA – the minimal supergravity model – the parameter $M_0$ is the mass of all scalar particles (squarks and sleptons) at the GUT scale.

$^{14}$Two jets are acoplanar if they are not back-to-back in the plane transverse with respect to the beamline.
with gluino-pair production being dominant, this gives rise to final states characterized by $E_T$ and a four-jet topology. An additional (“three-jet”) scenario is also considered in order to cover the intermediate situations, in which squarks and gluinos have a similar mass.

The analysis performed by DØ [13] makes a preselection, which requires at least two acoplanar central jets\(^\text{15}\) with a transverse momentum greater than 60 and 40 GeV/c; the electromagnetic energy fraction in each jet cannot exceed 95% of the total energy in order to reject purely electromagnetic sources (such as electrons and photons). A comparison between the jet energy as measured in the calorimeter and its counterpart carried by the charged tracks recorded in the tracking devices\(^\text{16}\) reduces the possibilities of selecting false jets and wrong primary vertices (which eventually lead to energy mismeasurements). Finally, the requirement $H_T \equiv |\sum_{\text{jets}} p_T| > 40 \text{ GeV/c}$ is applied.

![Figure 9. DØ search for strongly interacting supersymmetric particles in the missing $E_T$+jets signature [13]. Clockwise, from top left: missing $E_T$ spectra for the preselection stage and marginal $E_T$ distributions for “low $M_0$”, “three-jet” and “high $M_0$” analyses. In all plots, the non-QCD SM background is compared to signal and data; for the “high $M_0$” analysis, the fitted QCD background is also shown (solid curve). All analyses rely on 310 pb\(^{-1}\).](image)

After the preselection, dedicated analyses\(^\text{17}\) are performed to reach the maximum sensitivity in each of the three topologies. In all cases, leptons are vetoed and the $E_T$ vector is required to be well separated by any reconstructed jet in the event; this is aimed at reducing the dominant backgrounds, namely SM Z/W+jets and QCD multijet production.

The final requirements are performed on the missing $E_T$ (marginal distributions for the considered scenarios being shown in figure 9) and on the jet-based quantity $H_T \equiv |\sum_{\text{jets}} p_T|$: $(E_T > 175 \text{ GeV}, H_T > 250 \text{ GeV/c})$, $(E_T > 100 \text{ GeV}, H_T > 325 \text{ GeV/c})$ and

\(^{15}\) Central calorimeters – due to their larger distance from the beamline – suffer less from beam halo contaminations and, being located within the geometrical acceptance of the inner tracking devices, benefit from more reliable calibration techniques combining calorimetry with tracking information. As previously mentioned, they are typically contained in the $|\eta| \lesssim 1$ region.

\(^{16}\) This procedure leads to the definition of the so called jet charged fraction.

\(^{17}\) A similar analysis is performed by CDF based on 371 pb\(^{-1}\) [14].
(\(E_T > 175\ \text{GeV}, H_T > 250\ \text{GeV/c}\)) are required for the “low \(M_0\)”, “three-jet” and “high \(M_0\)” topology respectively. The overall signal efficiency ranges on the few percent level.

Once the analyses are optimized for the three topologies, additional signal points are generated under the same \(\tan \beta, A_0\) and sign(\(\mu\)) conditions\(^{18}\): only \(M_0\) and \(M_{1/2}\) are varied in order to scan the \((\tilde{m}_q, \tilde{m}_g)\) plane; each point is classified as “low \(M_0\)”, “three-jet” or “high \(M_0\)” according to the signature it leads to. Based on the agreement between events observed in data and events predicted by the SM, 95% confidence level cross-section limits are set in the no-signal hypothesis; these limits establish the criteria according to which each generated signal point is tested against exclusion. This is the procedure eventually leading to Figure 10.

\[\begin{align*}
\text{Figure 10.} & \quad \text{CDF and DØ 95\% confidence level excluded regions in the} (M_{\tilde{g}}, M_{\tilde{q}}) \text{ plane for the mSUGRA framework with tan} \beta = 3, \ A_0 = 0, \ \mu < 0 \ [13, 14].
\end{align*}\]

3.2. Searches for scalar top quarks

In \(R\)-parity conserving scenarios, assuming a stop\(^{19}\) lighter than the top quark\(^{20}\), three possible decay modes should be considered: the two-body decay channel \(\tilde{t}_1 \rightarrow b \tilde{\chi}_1^+ (\rightarrow b W^+ \tilde{\chi}_1^-)\) dominates when \(M_t > M_b + M_{\tilde{\chi}_1^+}\) or \(M_t > M_b + M_W + M_{\tilde{\chi}_0}\); if the previous mode is kinematically forbidden and the sneutrino is sufficiently light (\(M_t > M_b + M_{\tilde{\chi}_0}\)), then the three-body decay \(\tilde{t}_1 \rightarrow b \ell \tilde{\nu}\) takes place. If both tree-level channels are kinematically forbidden, then the top decays according to \(\tilde{t}_1 \rightarrow c \tilde{\chi}_1^0\), which is driven by a loop diagram. The last two decay modes have been studied by DØ and CDF respectively.

CDF has searched for stop quark pair production in final states consisting of two charm quarks and two neutralinos \([15]\), leading to two acoplanar heavy-flavour jets and large missing transverse energy. The preselection is aimed at correcting the missing \(E_T\) for instrumental noise, accelerator background and cosmic-ray contamination; variables based on jet electromagnetic energy fractions and charged fractions (which have been defined previously in this document).

\(^{18}\)The mSUGRA fixed parameters are \(\tan \beta = 3, \ A_0 = 0\) and \(\mu < 0\). The three template points correspond to \((M_0, M_{1/2}) = (25, 145)\ \text{GeV/c}^2, (M_0, M_{1/2}) = (500, 80)\ \text{GeV/c}^2\) and \(M_{\tilde{g}} = M_{\tilde{q}} = 325\ \text{GeV/c}^2\) for the “low \(M_0\)”, “high \(M_0\)” and “three-jet” scenario respectively.

\(^{19}\)In presence of mixing, sfermion mass eigenstates are commonly labelled as \(\tilde{t}_1\) and \(\tilde{t}_2\) in order of increasing mass.

\(^{20}\)This assumption is not at all restrictive, since the top squark mass range up to the top quark mass has not been completely excluded yet.
have been used for this purpose. In order to compensate for non-linearities in the calorimeter response, jets and \( E_T \) are corrected by the jet energy corrections, which account also for adjustments in the calorimeter energy scale. In order to avoid energy mismeasurements, a minimum separation is required between the missing \( E_T \) vector and each of the jets; complying with the signal signature, a minimum of two high-\( E_T \) jets is required in the \( |\eta| < 2.5 \) region, with the leading pair (\( E_T > 35 \) and \( 25 \) GeV) being central and acoplanar.

Non-QCD SM backgrounds, dominated by W/Z+jets production, can be reduced by vetoing isolated leptons; the remaining \( W \to \tau \nu + \)jets can be controlled by exploiting the different characteristics of the pencil-like jet arising from hadronic decay of taus\(^{21} \) with respect to hadronic jets: the latter are found to embed a larger number of tracks. Track multiplicity is then required to be greater than three, while an upper limit (at 12) is put in order to reduce the contribution of b-jets from \( t\bar{t} \) and QCD production. The missing \( E_T \) spectrum after requiring \( E_T > 55 \) GeV and before flavour-tagging is shown in figure 11.

Identification techniques relying on secondary vertex reconstruction are much less efficient when targeting c instead of b quarks: the tagging efficiency falls by a factor three or four below the corresponding b-tagging level as a consequence of the shorter lifetime of C hadrons. A viable alternative to secondary vertex reconstruction is offered by the jet probability tagging method: based on the charged tracks embedded in the jet, it computes the probability of the jet being originated from the primary interaction vertex\(^{22} \). According to this definition, jets from heavy-flavour fragmentation exhibit small values of jet probabilities; typical tagging probabilities are 1 and 5\% , corresponding to efficiencies on heavy-flavour identification of 25 and 30\% and mis-tag probabilities\(^{23} \) of the order of 1 and 5\% respectively. At least one of the leading central jets is required to be flagged with a 5\% jet probability tag; the remaining jet – if in the acceptance of the tagger – is required to fulfill a mild 45\% jet probability tag.

At the end of the selection (with a total acceptance of the order of 1\% in the hypothesis of a 40 GeV/c\(^2 \) neutralino), the number of observed events in 163 pb\(^{-1} \) of data is compatible with the SM predictions; based on the assumption of no signal, and with the estimated systematic

\(^{21} \) Note that taus are not removed by the lepton veto, which refer to isolated electrons and muons only.

\(^{22} \) The acceptance of the jet probability procedure is determined by the geometrical coverage of tracking, approximately \( |\eta| < 1.5 \) (forward tracking is not included in the jet-probability algorithm).

\(^{23} \) Defined as the probabilities of mistakenly identifying light quark or gluon jets as heavy-flavour jets.
uncertainties, a 95% confidence level upper limit on $\sigma(pp\to \tilde{t}_1 \tilde{t}_1) \times BR(\tilde{t}_1 \to c\tilde{\chi}^0_1)\times BR(\tilde{\chi}^0_1 \to \mu\nu)$ is computed by means of a Bayesian likelihood method. As shown in figure 11, no lower limit on $M_{\tilde{t}}$ is reached since the experimental sensitivity is still limited by systematic uncertainties, dominated by scale effects on the jet probability algorithm.

Figure 12. DØ search for direct pair production of scalar top quarks in final states with two leptons, two b quarks and missing $E_T$ using 350 pb$^{-1}$ [16, 17, 18]. Left: di-muon invariant mass spectrum (data compared to major background contributions). Right: 95% confidence level excluded range in the $(M_{\tilde{t}}, M_{\tilde{\nu}})$ plane for the combined search.

If an intermediate stop mass is considered, stop decays are dominated by the three-body $\tilde{t}_1 \to b\ell\tilde{\nu}$ channel. DØ has searched for stop pair production in final states characterized by two b-jets, two opposite-signed leptons and $E_T$, probing both the electron-muon and di-muon signatures [16, 17, 18].

Both analyses require a pair of oppositely charged leptons: in one case a pair of muons of $p_{T}\mu > 8$ and 6 GeV/c, in the other a muon of $p_{T}\mu > 8$ GeV/c and an electron with $E_{T}\mu > 12$ GeV. Muons are selected using tracks in the central tracking system in combination with patterns of hits recorded by the muon chambers in correspondence of the track extrapolated position; muons are required to be isolated in both the calorimeter and the tracker. Electrons are reconstructed based on their characteristic energy deposition in the calorimeter, including the transverse and longitudinal shower profiles. In order to distinguish electrons from photons, a track must point to the energy deposition in the calorimeter; besides, its momentum must be consistent with the energy recorded in the calorimeter.

Muons essentially behave as minimum ionizing particles and the energy they release in the electromagnetic and hadronic calorimeters is almost independent of their momentum. For this reason, the presence of a muon in an event is associated to a fake $E_T$ signal of magnitude approximately given by the difference between the muon $p_T(\sim E_T)$ and the transverse energy deposited in the calorimeter, with the direction of the $E_T$ vector being collinear with the muon. Therefore, once muon candidates are identified, a corrected $E_T$ can be obtained from the vector sum:

$$E_T^{corr} = E_T^{raw} - \sum_{N_{\mu\mu\mu}} [p_T(\mu \text{ track}) - E_T(\mu \text{ calorimetry})].$$

That is a top squark with mass such that $M_{\tilde{b}} + M_{\tilde{\nu}} + M_{\tilde{t}} < M_{\tilde{t}} < \min(M_{\tilde{b}} + M_{\tilde{\chi}^+_1}, M_{\tilde{b}} + M_{W} + M_{\tilde{\chi}^0_1}).$

The requirement mimicks the well known $E/p \simeq 1$, which holds for relativistic electrons.

24 That is a top squark with mass such that $M_{\tilde{b}} + M_{\tilde{\nu}} + M_{\tilde{t}} < M_{\tilde{t}} < \min(M_{\tilde{b}} + M_{\tilde{\chi}^+_1}, M_{\tilde{b}} + M_{W} + M_{\tilde{\chi}^0_1}).$

25 The requirement mimicks the well known $E/p \simeq 1$, which holds for relativistic electrons.
After correcting the $E_T$, in order to avoid other sources of energy (and $E_T$) mismeasurement, a minimum opening angle is required between the direction of the $E_T$ vector and the directions of the reconstructed leptons projected on the transverse plane.

Physical background is dominated by $Z$+jets production, with contributions from $Z\rightarrow \mu\mu$ in the $\mu\mu$ and $Z\rightarrow \tau\tau$ in the $e\mu$ channels. One of the main characteristic of the signal is given by jets originating from the hadronization of $b$ quarks; $Z$+jets, on the other hand, owes the presence of jets to initial state radiation gluons which typically give rise to less energetic jets, yielding a lower jet multiplicity. This fact is exploited in the reduction of the $Z\rightarrow \mu\mu$+jets contribution (shown, together with other backgrounds, in figure 12) by requiring at least one central jet being $b$-tagged. The absence of neutrinos in $Z\rightarrow \mu\mu$ is used to keep ˜χ0 extrapolate the QCD contribution at higher values of $\frac{M}{E_T}$; they are completely rejected after at least one jet is required to be $b$-tagged. The same action is not taken against the homologous $Z\rightarrow \tau\tau$+jets, which is already suppressed by the branching ratios for the ˜τ→eντντ and ˜τ→μντμτ decays: instead, the different jet activity observed in signal and background is parametrized in terms of number of non-isolated tracks in the event (NIT) and used as an additional input for the procedure used for setting limits. The absence of neutrinos in $Z\rightarrow \mu\mu$ decays allows further background reduction by vetoing events with $75 \leq M_{\mu\mu} \leq 120$ GeV/$c^2$ and $E_T < 50$ GeV. A final cut on $E_T$ is applied to both samples: $E_T > 15$ GeV is required for $Z\rightarrow \tau\tau$+jets, while for $Z\rightarrow \mu\mu$+jets a two-dimensional cut exploiting the correlation between the missing transverse energy and the opening angle $\Delta \phi (E_T, p_T(leading \mu))$ is found to maximize the signal statistical significance.

After the selection, the number of observed $\mu\mu$ and $e\mu$ events in samples of 339 and 350 pb$^{-1}$ of data are found to be in agreement with the corresponding SM predictions. Limits are then set under the assumption of no signal: a modified frequentist method based on $H_T (\mu\mu$ channel) and on a combination of NIT and $p_T \equiv p_T^\ell+p_T^\nu+\Delta \phi = 20, \mu = 225$ GeV, $M_2 = 500$ GeV/$c^2$ and $M_A = 800$ GeV/$c^2$ are common parameters; slepton parameters, including the trilinear stau mixing, have been set to obtain equal branching ratios of ˜χ± to all lepton flavours (assuming a 100% decay ˜χ± → ℓν). Finally, $A_t$ has been varied to obtain the stop mass, $M_1$ (the bino mass) has been adjusted to keep ˜χ0 as LSP and $M_2$ (the wino mass) has been adjusted to keep ˜χ± virtual.

### 3.3. Search for scalar bottom quarks

The Tevatron analyses dedicated to bottom squark searches assume the decay of the sbottom into LSP, according to $b_1 \rightarrow b\tilde{\chi}_1^0$; both direct and indirect production mechanisms have been investigated.

The D0 search looks for sbottom squark direct pair production in topologies characterized by two acoplanar heavy-flavour jets and missing transverse energy [19]. Events are preselected with two or three high-$p_T$ jets; the two leading ones ($p_T > 40$ and 15 GeV/$c$) are also required to be central and acoplanar. Background from t¯t and diboson production is reduced by vetoing isolated leptons and tracks. Events in which the direction of the missing $E_T$ vector is close in azimuthal angle to any of the reconstructed jets are rejected; this reduces the possibility of energy mismeasurements and fake missing $E_T$ signals. Figure 13 shows the agreement between data and non-QCD SM background at high values of $E_T$; on the other hand, the discrepancy visible at low $E_T$ is interpreted in terms of QCD background. A fit to the $E_T < 60$ GeV region is then performed after subtraction of the other background (see the figure insert): this allows to extrapolate the QCD contribution at higher values of $E_T$. After the requirement $E_T > 60$ GeV only 80 QCD events survive; they are completely rejected after at least one jet is required to be $b$-tagged.

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26 Physical background, opposed to instrumental background, is due to SM processes with genuine leptons; instrumental background, on the other hand, is caused by an incorrect identification of leptons (fake leptons) and uncorrected sources of $E_T$.

27 Despite a signal signature with two jets, the applied cut takes into account the possibility of one jet being below threshold due to small values of $M_\ell - M_\nu$. 

The analysis performed up to this point is aimed at retaining a good sensitivity to signal, even for small sbottom masses – when jets are characterized by a softer spectrum. However, higher sensitivities can be reached by tightening the $E_T$ and jet $p_T$ thresholds; for this reason, two additional sets of final cuts are considered to target medium and large sbottom masses: $E_T > 80$ GeV with $p_T(1^{st} jet) > 40$ GeV/$c$, $p_T(2^{nd} jet) > 15$ GeV and $E_T > 100$ GeV with $p_T(1^{st} jet) > 70$ GeV/$c$, $p_T(2^{nd} jet) > 40$ GeV. No excess is observed in 310 pb$^{-1}$ of data relative to SM expectations in any of the considered regions.

A 95% confidence level exclusion on the $\sigma(pp \to \tilde{b}_1\bar{\tilde{b}}_1) \times BR(\tilde{b}_1 \to b\tilde{\chi}_1^0)$ is set by means of a modified frequentist method in the hypothesis of absence of signal; the result is interpreted within a MSSM framework in which only the sbottom and neutralino masses are varied, leading to an excluded region in the $(M_{\tilde{b}}, M_{\tilde{\chi}_1^0})$ plane, as shown in figure 13.

![Figure 13](image.png)

Figure 13. Search for direct pair production of scalar bottom quarks in 310 pb$^{-1}$ of data collected by the DØ detector [19]. Left: missing $E_T$ for data and non-QCD Standard Model background; detail shows agreement between data after Standard Model background subtraction and fitted QCD contamination. Right: 95% confidence level excluded range in the $(M_{\tilde{b}}, M_{\tilde{\chi}_1^0})$ plane.

The cross-section for gluino pair production is large compared to direct bottom squark pair production of similar mass; if the sbottom is light enough, then the two-body decay $\tilde{g} \to b\tilde{b}$ becomes kinematically allowed. The CDF analysis searches for sbottom production in gluino decays [20]: assuming each of the pair-produced gluinos decays according to $\tilde{g} \to b\tilde{1}\bar{b}$ with subsequent sbottom decay to a b quark and lightest neutralino ($b_1 \to b\tilde{\chi}_1^0$), a striking signature of four b-jets and large $E_T$ is achieved. Signal topology, however, depends crucially on the mass difference between the gluino and the sbottom: for small values of $M_{\tilde{g}} - M_{b_1}$, b quarks from gluino decays are characterized by low transverse energies due to the reduced phase space, which increases the chance of the corresponding b-jets of failing the reconstruction criteria. Conversely, if $M_{\tilde{g}} - M_{b_1}$ is large, neutralinos from sbottom decays receive a considerable boost.

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28 If the gluino is lighter than the average squark mass, three-body decays ($\tilde{g} \to q\tilde{\chi}_1^0; \tilde{g} \to q'\tilde{\chi}_1^\pm$) and the one-loop decay $\tilde{g} \to g\tilde{\chi}_1^0$ dominate unless – due to large mixing effects – the two-body modes $\tilde{g} \to b\tilde{b}$, $\tilde{t}\tilde{t}$ become kinematically accessible; vice versa, if the gluino is heavier than squarks, then the two-body decay $\tilde{g} \to q\tilde{q}$ dominates. This simple argument shows that gluino decays in $bb$ and $tt$ (the latter being penalized by the large top quark mass) are relevant whenever $t_1$ and $b_1$ are light.
which pushes them in a back-to-back configuration\textsuperscript{29}, yielding to a lower event $E_T$ due to mutual compensation.

Events are preselected by requiring $E_T > 35$ GeV and at least three $E_T > 15$ GeV jets within $|\eta| < 2.0$; at least one of these jets is required to be central. Most of the preselected events are yielded by QCD multijet events, with the $E_T$ arising from mismeasurements of the jet energies. This contribution can be reduced by requiring the angular distance in the plane transverse to the beamline between the leading three jets and the $E_T$ being greater than $40^\circ$.

Background modelling is checked in kinematic regions (control regions) in which signal contribution is expected to be negligible in order to avoid statistical biases. Since gluino events are characterized by the absence of high-$p_T$ isolated leptons and large $E_T$, three control regions can be built by reversing one at a time the two relations defining the signal region. As a result, the region with moderate $E_T$ and a high-$p_T$ isolated lepton is populated by QCD, top and W/Z+jets events, the region with moderate $E_T$ and no reconstructed lepton is dominated by QCD production and the region with high $E_T$ and at least one high-$p_T$ lepton by $t\bar{t}$.

Finally, signal events are divided in two disjoint classes: the exclusive single b-tag and the inclusive double b-tag; the $E_T$ distribution of each of them is shown in figure 14. In order to increase the signal sensitivity, a tighter $E_T > 80$ GeV. A counting experiment is performed on each categories, comparing the number of expected events with the number of observed events in 156 pb$^{-1}$ of data, which are found to be in good agreement. 95\% confidence level limits are then set on the cross-section for gluino pair production; this result can then be used for excluding regions of the $(M_{\tilde{g}}, M_{\tilde{b}_1})$ plane, as shown by figure 15, where a fixed neutralino mass has been used.

3.4. Search for chargino-neutralino production
Charginos and neutralinos, which are admixtures of the supersymmetric partners of Higgs and SU(2)$_L \times$ U(1)$_Y$ gauge bosons\textsuperscript{30}, can be directly produced at hadron colliders in electroweak

\textsuperscript{29}It is understood that a real back-to-back topology is achieved only within the $\tilde{g}\tilde{g}$ rest frame; alternatively, the back-to-back configuration is accomplished in terms of vectorial transverse momenta.

\textsuperscript{30}Chargino are generated from charged wino and higgsino components, while neutralinos include bino and neutral wino and higgsino components.
pair production processes; due to the nature of the interaction, the cross-section is rather small. Nevertheless, in SUSY-breaking scenarios inspired to minimal supergravity (mSUGRA), charginos and neutralinos can be much lighter than squarks and gluinos and could be the only SUSY particles accessible at the Tevatron. Besides, under the assumption of $R$–parity conservation, a very distinct signature can be achieved in the decay pattern of a $\tilde{\chi}^\pm_1 \tilde{\chi}^0_2$ chargino-neutralino pair: if the chargino decays according to $\tilde{\chi}^\pm_1 \rightarrow \tilde{\chi}^0_1 W^\pm, \nu \tilde{\ell}^\pm, \ell^\pm \tilde{\nu}_\ell$, and the neutralino as $\tilde{\chi}^0_2 \rightarrow Z\tilde{\chi}^0_1, h\tilde{\chi}^0_1, \ell^\pm \ell^\mp \rightarrow \ell^+ \ell^- \chi^0_1$, the final state is characterized by three charged leptons and $E_T$ from the escaping neutrino and neutralino LSPs. If any two-body decay channel is not kinematically allowed, the corresponding decay occurs through a three-body process: in this case, the reaction involves a virtual (vector or Higgs) boson or a virtual slepton. If sleptons are much heavier than $\tilde{\chi}^\pm_1$ and $\tilde{\chi}^0_2$, then chargino and neutralino decays to virtual vector bosons are dominant; virtual Z and W bosons eventually behave as on-shell particles in their decays. However, different behaviours are expected for charginos and neutralinos as a consequence of their different couplings to vector bosons: while both components of the chargino couple to charged W bosons, only the higgsino component of the neutralino couples to the neutral Z boson; this means that the rate of neutralino decays involving a Z boson is proportional to the higgino component of the neutralino. Vice versa, if sleptons are lighter than $\tilde{\chi}^\pm_1$ and $\tilde{\chi}^0_2$, the contribution of virtual sleptons can modify substantially the previous pattern: as a result, specific decay channels can be enhanced or suppressed in a model-dependent way.

Two concurrent processes contribute to chargino-neutralino production\textsuperscript{33}: a virtual W boson s-channel exchange and a virtual squark t-channel exchange, with the two contributions interfering destructively. If the case when all squarks are heavy, the s-channel contribution dominates; conversely, when squarks are sufficiently light, both s-channel and t-channel

\textsuperscript{31} Where Higgs bosons are involved, the corresponding decay in $\tau^+ \tau^-$ or $\tau^\pm \nu$ is understood.

\textsuperscript{32} Note that $M_{\tilde{\chi}^\pm_1} \simeq M_{\tilde{\chi}^0_2}$ in mSUGRA.

\textsuperscript{33} $\tilde{\chi}^\pm_1 \tilde{\chi}^0_2$ and $\tilde{\chi}^\pm_1 \tilde{\chi}^\mp_1$ production processes are also permitted; however, their observation is compromised by W+jets and WW production, which lead to similar signatures and are characterized by cross-sections of the order of 3 nb and 10 pb, large if compared to signal ($\sigma \sim 1$ pb). Note that production cross-sections are not simple functions of chargino-neutralino masses, but depend also on their mixings and on squark masses [21].
contributions are relevant (and so is the interference term).

Figure 16. Chargino-neutralino associated production with subsequent decay into three-lepton final states: cross-section as a function of tan\(\beta\) for two mSUGRA points in typical Tevatron Run II scenarios [22].

The yield of \(\tilde{\chi}_1^\pm\tilde{\chi}_2^0\) events into a trilepton final state is determined mainly by the branching ratio of the chargino-neutralino pair, which – as anticipated – is highly model-dependent; figure 16 shows the lepton flavour composition of the trilepton sample as a function of tan\(\beta\) for two typical SUSY scenarios. The rise of tauonic channels for increasing values of tan\(\beta\) is evident.

The efficiency of the trilepton selection depends crucially on the mass splitting between the \(\tilde{\chi}_1^\pm\) and \(\tilde{\chi}_2^0\) and the LSP (\(\tilde{\chi}_1^0\)): if this splitting is small, one or more leptons can be emitted at low momentum and fail to be reconstructed. Therefore, key features of the analysis are the charged lepton efficiency and acceptance, with their product entering as a third power in the overall signal efficiency.

The DØ analysis [23] defines four different selections based on the lepton content of the final state: two electrons or two muons plus an additional lepton (ee\(\ell\) and \(\mu\mu\ell\) respectively), two muons with equal charge (\(\mu^+\mu^-\)) and an electron, a muon plus a third lepton (e\(\mu\ell\); in all cases, the third lepton can be of any flavour (therefore including tau leptons). A similar approach is followed by CDF [24]. Each selection requires two identified isolated leptons with minimum transverse momenta (\(p_T > 11 \div 12\) GeV/c for leading and \(p_T > 5 \div 8\) GeV/c for next-to-leading leptons). DY and WZ diboson backgrounds, where di-electrons and di-muons are typically back-to-back in the transverse plane and with an invarian mass close to \(M_Z\), are suppressed by means of cuts on the di-lepton invariant mass and azimuthal opening angle. QCD multijet events, which contribute to background due to lepton misidentification ("fake" leptons), are rejected if they contain jets and have a small \(E_T\) significance\(^{34}\) or if the \(E_T\) direction is aligned with leptons; when this happens, the \(E_T\) is most likely due to an incorrect determination (underestimation) of the lepton energy, which results also in a smaller value of the di-lepton transverse mass. \(t\bar{t}\) events can be rejected with a cut on \(H_T\) (scalar sum of the transverse momenta of jets).

A third lepton requirement is then applied for further signal discrimination: the selection is relaxed to an isolated track\(^{35}\) originating from the same vertex of the two leading leptons; no additional lepton identification selection is performed in order to maximize the acceptance.

\(^{34}\text{Missing }E_T\text{ significance is defined as the ratio between }E_T\text{ and its uncertainty; small values of this ratio may be symptom of energy mismeasurements.}\)

\(^{35}\text{The isolation criteria is optimized in order to be efficient in selecting tau leptons also.}\)
requirement proves to be very efficient in the reduction of background events, where isolated tracks originating from the underlying events or jets are typically characterized by low transverse momenta (below 10 GeV/c).

Figure 17. 95% confidence level upper limits for chargino-neutralino associated production in the three-lepton final states compared to theoretical predictions in a scenario with no slepton mixing. Left: DØ result based on 320 pb⁻¹ of data [23]. Right: CDF result based on 310 ± 750 pb⁻¹ [24].

CDF cross-checks the background modeling in statistically unbiased samples, finding data and expectations in agreement within the experimental uncertainties in all control regions. The number of events selected agrees with the SM expectation also in the signal region; 95% confidence level upper limits on the cross-section for the process \( \sigma(p\bar{p} \rightarrow \widetilde{\chi}^0 \rightarrow 3\ell + X) \) can be set under the assumption that no signal is present. This is eventually translated into a (model-dependent) lower limit on the chargino mass. Results are collected in figure 17. DØ tests three different scenarios: “heavy–squarks”, where the scalar mass unification is relaxed (model-dependent) lower limit on the chargino mass. Results are collected in figure 17. DØ tests three different scenarios: “heavy–squarks”, where the scalar mass unification is relaxed in order to enhance the production cross-section, “3ℓ–max”, corresponding to a maximal branching fraction, and “large M0”, where a high slepton mass forces the chargino-neutralino decay through vector boson exchange, resulting in relatively small leptonic branching fractions. The CDF limit is obtained in the “3ℓ–max” scenario.

3.5. Search for anomalous events with di-photon and missing transverse energy
In low-scale gauge-mediated SUSY-breaking scenarios, the LSP role is played by \( \tilde{G} \), a light gravitino; in this case, all superpartners of the MSSM acquire an additional decay mode, namely \( X \rightarrow \tilde{G} \). However, if the SUSY-breaking occurs beyond the electroweak scale, ordinary decays mediated by strong and electroweak interactions – when kinematically accessible – dominate over decays through gravitinos. Therefore, heavy SUSY particles rapidly cascade into the lightest SM superpartner, which – if \( R \)-parity is conserved – is stable with respect to

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**Footnotes:**

36 Control regions have been defined according to \( E_T \), leading di-lepton invariant mass and jet multiplicity.

37 The t-channel squark exchange being suppressed by high squark masses.

38 More precisely, the spin-\( \frac{1}{2} \) longitudinal components of the gravitino – the superpartner of the graviton – acquire mass by absorbing the Goldstino, which is the relic of the spontaneous braking of SUSY; the Goldstino, being very light, is therefore the actual LSP. However, since the 3/2-helicity states of the gravitino couple only with gravitational strength and are never relevant in collider phenomenology, only the essentially massless Goldstino-gravitino can be considered. For this reason, the two terms gravitino and Goldstino are used indifferently in literature.
all the MSSM couplings; only at this point the decay to the gravitino LSP takes place. The lightest SM superpartner is therefore the next-to-lightest SUSY particle (NLSP). If the NSLP is the lightest neutralino, then the process $\tilde{\chi}^0_1 \rightarrow \tilde{G} \gamma$ can conclude the decay cascade to the LSP\textsuperscript{39}. Therefore, pair production and subsequent decay of SUSY particles lead to a $\gamma\gamma + \not{E}_T$ final state signature. In the considered scenario, $\tilde{\chi}^\pm_1 \tilde{\chi}^\mp_1$ and $\tilde{\chi}^\pm_1 \tilde{\chi}^0_2$ production are the dominating mechanisms yielding to events with di-photons and large $\not{E}_T$; due to small SM contributions, the search for anomalous two-photon production becomes an attractive probe for new physics.

\begin{figure}[h]
\centering
\includegraphics[width=\textwidth]{diagram.png}
\caption{Search for anomalous production of diphoton events. Top: missing $E_T$ spectra observed by CDF with 202 pb$^{-1}$ (left) \cite{26} and by DØ with 263 pb$^{-1}$ (right) \cite{27}. Bottom: CDF and DØ combined 95\% confidence level upper limit on the diphoton plus missing $E_T$ production cross-section \cite{28}.}
\end{figure}

Both CDF and DØ analyses \cite{26, 27} select events by requiring two highly energetic central photons compatible with originating from the same primary vertex; photons are identified primarily by energy depositions in the calorimeter. In order to remove contamination from hadronic jets, photons need a considerable fraction ($\gtrsim 95\%$) of their energy to be of electromagnetic origin; further separation from jets is achieved by requiring photons to be

\textsuperscript{39}Being the neutralino a superposition of bino and neutral wino and higgsinos, other possibilities include decays to a neutral vector or Higgs boson; however, for the energies reached at the Tevatron, the decay to photon is favoured.
isolated, both in the calorimeter and in the tracking volume. Contamination by $\pi^0 \rightarrow \gamma \gamma$ is reduced by requiring the transverse shower profile to be compatible with a single photon hypothesis. Electrons are rejected by vetoing tracks that extrapolate to the electromagnetic cluster. In order to avoid missing $E_T$ mismeasurements, events where the direction of $E_T$ is back-to-back to a jet are removed. The $E_T$ spectra observed by CDF and DØ are shown in figure 18; both experiments require a $E_T > 40$ GeV cut.

Apart from small physical background contributions from $W(\rightarrow e\nu)+$jets, $\gamma W(\rightarrow e\nu)+$jets and $\gamma Z(\rightarrow e^+e^-\nu)+$jets production with an electron misidentified as a photon, the dominant background comes from QCD events with either real photons or jets misidentified as photons. The shape of the QCD contribution is determined from data events selected with looser photon isolation and identification criteria; the normalization is obtained from the number of data events in the $E_T < 20$ GeV portion of the spectrum, where no other background is expected to contribute.

The number of observed events is in good agreement with the predictions; therefore, an upper limit on $\sigma(p\bar{p} \rightarrow \tilde{\chi}^\pm_1 \tilde{\chi}^\mp_1, \chi^0_2 \rightarrow \gamma \gamma /E_T + X)$ is obtained in the hypothesis of absence of signal by means of a Bayesian approach based on a flat prior cross-section distribution; this can be translated in a 95% confidence level lower limit on the chargino mass. The combined result for CDF and DØ [28], which improves the limits set individually by the two experiments, is shown in figure 18; lower limits for the chargino and neutralino masses are also shown.

4. Supersymmetry with violated $R$–parity

In the previous session, supersymmetric processes have been discussed under the assumption of $R$–parity conservation, which was introduced in order to avoid phenomenological problems induced by lepton and baryon number violations; however, there are no reasons other than simplicity why it should be $R$–parity and not some other symmetry which suitably suppresses $B$ and $L$ violations.

In case of violated $R$–parity, the superpotential terms described in equation (2) are not forbidden any longer: if non-vanishing values of $\mu', \lambda, \lambda'$ and $\lambda''$ come into play, a new variety of signatures can arise [29]. However, the topology of the events that could be observed in the eventuality of $R$–parity violation depends crucially on the strength of the $R$–parity violating couplings: weak values are likely to be observed mostly in the decays of SUSY particles, especially of the LSP, which is not any longer constrained to be stable; conversely, increasingly stronger couplings can contribute to the indirect and direct production of single supersymmetric particles. The common practice to front a potentially chaotic situation due to a huge number of free parameters entering the $R$–parity violating Lagrangian is to assume a strong (arbitrary) hierarchy among the couplings; the simplest implementation of such practice is to postulate – for a given search strategy – the existence of a single dominant $R$–parity violating coupling.

4.1. $R$–parity violating scalar top quark decays

In the hypothesis of a light top squark, the $R$–parity violating $\tilde{t}_1 \rightarrow b\tau^+\tau^-$ channel ($\lambda_{333}' \neq 0$) can become dominant over the three-body $t_1 \rightarrow b\ell^+\bar{\nu}$ and one-loop $t_1 \rightarrow c\chi^0_1$ decay modes. In this scenario, pair-produced top squarks can lead to two $b$ quarks and two tau leptons.

The CDF analysis searches for stop pair production in final states characterized by either an electron or a muon from the $\tau \rightarrow \ell\nu\nu_\tau$ decay, one semi-hadronically decaying tau and two or more jets [30]. Events are preselected by requiring at least one isolated central lepton identified as an electron or muon with $p_T > 10$ GeV/$c$; events with the primary lepton being associated to either a photon conversion or a cosmic ray are vetoed. The tau candidate decaying in semi-hadronic mode is extracted from the chargino mass by exploiting the mass relations embedded in the model.
is reconstructed by means of the procedure described on page 4. Jets \( (E_T > 20 \text{ GeV} \text{ and } |\eta| < 2.4) \) are required to be well separated from all reconstructed leptons. In order to reduce the dominant \( Z + \text{jet} \) background, events are rejected if the invariant mass of the primary electron (muon) and a second, loosely identified electron (muon) with opposite charge lies in the \( Z \) window mass \((76, 106) \text{ GeV}/c^2\); the same cut is applied to the primary lepton-tau candidate pair (regardless of their charge) if the two leptons are collinear in the transverse plane. Further suppression of the \( Z \rightarrow \tau^+\tau^- + \text{jets} \) background is obtained by setting a cutoff on \( Y_T(\equiv |p_T^\ell | + |p_T^\tau | + E_T) \).

**Figure 19.** Search for pair production of scalar top quarks both decaying to a \( \tau \) lepton and a \( b \) quark in 322 pb\(^{-1}\) of CDF data [30]. Left: number of jets. Right: transverse mass for the lepton-\( E_T \) system.

**Figure 20.** Upper limit on the pair production of scalar top quarks both decaying to a \( \tau \) lepton and a \( b \) quark from 322 pb\(^{-1}\) of CDF data: limit lowers to \( M_{\tilde{t}_1} = 151 \text{ GeV}/c^2 \) if theoretical uncertainties on parton distribution functions and factorization scale are taken into account [30].

Signal is searched in the region with two or more jets and \( M_T(\ell, E_T) < 35 \text{ GeV}/c^2 \); figure 19 illustrates the distribution for these quantities in signal, background and data samples. Background predictions have been checked against data in control regions defined on the basis of jet multiplicity and \( M_T(\ell, E_T) \): the number of data events in all regions is found in good agreement with the SM predictions. With no excess in the signal region, a 95% confidence level
upper limit is set on the stop pair production cross-section by means of a likelihood method in the hypothesis of $BR(\tilde{t}_1 \rightarrow b\tau^+) = 100\%$; the result is shown in figure 20. Correspondingly, a lower limit of 155 GeV/c$^2$ on the lightest stop can be set.

4.2. $R$–parity violating decay of the Lightest Supersymmetric Particle

When $R$–parity is violated, the associated production of chargino-neutralino or chargino-chargino can lead to even more striking signatures than the trilepton topology considered in § 3.4: if a non trivial value of $\lambda_{ijk}$ is assumed, then the neutralino can decay into two charged leptons and one neutrino (see figure 21). Under the assumption that both chargino and neutralino decay under $R$–parity conservation$^{41}$, four charged leptons are then yielded in the final state by the decay of the two neutralinos.

*Figure 21. Decay of the neutralino through the non trivial $R$–parity violating $\lambda_{1jk}$ trilinear coupling.*

The DØ search$^{31}$ is focussed on processes where $\lambda_{1jk} \neq 0$; three different couplings are considered: $\lambda_{121}$, $\lambda_{122}$ and $\lambda_{133}$, where only one coupling is assumed to be dominant at a time$^{42}$. It should be noted that in this situation the branching ratios of the neutralino is not influenced by the magnitude of the non trivial $R$–parity violating parameter, which, on the other hand, affects the lifetime of the neutralino; only neutralinos decaying within 1 cm from the primary interaction vertex are considered in the analysis. The search is interpreted within a MSSM scenario characterized by heavy squarks and sleptons (1 TeV/c$^2$), $\tan \beta = 5$ and a large higgsino mixing term $\mu = 1$ TeV. In this framework, the dominant contribution to the signal is given by $\tilde{\chi}_1^0 \tilde{\chi}_2^0$ and $\tilde{\chi}_1^\pm \tilde{\chi}_1^\mp$.

As the neutralino can be light, the charged leptons generated in its decay can have small transverse momenta and therefore be difficult to reconstruct; for this reason, only three charged leptons are required with flavour combinations $e\ell$, $\mu\mu \ell$ and $e\ell$ ($\ell = e, \mu$), corresponding to a non trivial value of $\lambda_{121}$, $\lambda_{122}$ and $\lambda_{133}$ respectively. Electrons and muons (with minimum $p_T$ ranging between 8 and 20 GeV/c, depending on the search channel) are required to be isolated from each other and from any hadronic jet; tau leptons are considered only in the semi-hadronic decay mode: they are identified as isolated pencil-like jets in a similar fashion to the procedure described on page 4. In the $\mu\ell\ell$ analysis, looser identification criteria are applied to the least energetic lepton in order to increase the acceptance. Background contributions from di-lepton resonances and DY production are suppressed by means of two-dimensional cuts on the ($E_T, M_{\ell\ell}$) and ($E_T, \Delta\phi(\ell\ell)$) planes. Electrons consistent with being originated in a photon conversion are not considered in the selection.

$^{41}$ The weak limit of $R$–parity violation is considered.

$^{42}$ A similar search has been performed by CDF$^{32}$. 
Background contributions from QCD production are estimated from data events selected with loose lepton identification criteria and cross-checked, together with all other background predictions, in a variety of control samples. The number of events observed in data is in good agreement with the expectations from the SM in all control and signal regions. A 95% confidence level upper limit on $\sigma(p\bar{p} \rightarrow \tilde{\chi}^\pm_1 \tilde{\chi}^0_0, \tilde{\chi}^\pm_1 \tilde{\chi}^{\mp}_1 \rightarrow \ell\ell, \mu\mu\ell, \ell\nu\ell + X)$ is extracted by means of a modified frequentist approach in the hypothesis of absence of signal. The three search analyses are combined for each coupling in order to improve the limit; for the combination phase, events selected by multiple analyses are assigned only to the search channel with the largest signal to background ratio and removed from all the other analyses. The results of this procedure are summarized in figure 22; the visible gap at low neutralino masses is due to the explicit cut on the mean decay length and on the reference values used as Monte Carlo inputs for the $R$–parity violating couplings ($0.01$ for $\lambda_{121}$, $\lambda_{122}$ and $0.003$ for $\lambda_{133}$), which are still related to the lifetime of the neutralino.

4.3. $R$–parity violating production and decays of resonant scalar leptons
Under the assumption that only $\lambda'_{211} \neq 0$, at $p\bar{p}$ colliders an initial $u\bar{d}$ or $d\bar{d}$ can produce – via a dominant s-channel process – a single scalar muon $\tilde{\mu}$ or a scalar muon neutrino $\tilde{\nu}_\mu$. The smuon can then decay either in a muon-neutralino ($\tilde{\mu}^\pm \rightarrow \mu^\pm \tilde{\chi}^0_0$) or in a neutrino-chargino ($\tilde{\mu}^\pm \rightarrow \nu_\mu \tilde{\chi}^\pm$) pair without violating the $R$–parity; the muon sneutrino, on the other hand, decays either in neutrino-neutralino ($\tilde{\nu}_\mu \rightarrow \nu_\mu \tilde{\chi}^0_0$) or in muon-chargino ($\tilde{\nu}_\mu \rightarrow \mu^\pm \tilde{\chi}^\mp$). The non trivial $\lambda'_{211}$ coupling finally allows the lightest neutralino to decay in two quarks of the first generation and a muon or muon neutrino. Therefore, several different signatures can emerge from the decay of a resonant muon or muon sneutrino; however, in order to achieve maximum discrimination with respect to QCD background, only events with at least two muons are considered. This signature essentially accepts contributions from three channels: $\tilde{\mu}^\pm \rightarrow \mu^\pm \tilde{\chi}^0_1, \tilde{\mu}^\pm \rightarrow \mu^\pm \tilde{\chi}^0_{2,3,4}$ and $\tilde{\nu}_\mu \rightarrow \mu^\pm \tilde{\chi}^{\mp}_1$: channels

The DØ analysis [33] selects events if two isolated muons are identified in the central part of the detector, with transverse momenta exceeding 15 and 8 GeV/$c$. In all channels, jets from the fragmentation of light quarks are expected: a minimum of two $p_T > 15$ GeV/$c$ jets within $|\eta|=2$ is required in order to increase the signal sensitivity with respect to the dominant SM
background from di-muon DY production. The invariant mass of the di-muons and the two leading jets – which corresponds to the mass of the resonant slepton in the $\mu^± \rightarrow \mu^± \tilde{\chi}_1^0$ channel – is shown in figure 23; note that the remaining signal channels, being characterized by cascade decays, contribute on average to lower invariant mass values. Background from QCD multijet production is extracted from data by means of looser muon isolation requirements.

Figure 23. Search for a resonant scalar muon using 380 pb$^{-1}$ of data collected by the DØ detector. Left: invariant mass of the two muons and two leading jets of the event. Right: 95% confidence level exclusion contours in the $(M_{\tilde{\chi}_1^0}, M_{\tilde{\ell}})$ plane for a $\tan \beta = 5$, $\mu < 0$ and $A_0 = 0$ mSUGRA scenario with [33].

Good agreement between the SM expectations and the observed number of events is found at all stages of the analysis. In the absence of an excess in the data, a 95% confidence level upper limit is set to $\sigma(p\bar{p}\rightarrow \mu^±, \tilde{\nu}_\mu \rightarrow 2\mu + \geq 2$ jets) by means of a modified frequentist approach. The $\tilde{\chi}_1^0$ decay branching fractions are predicted within a mSUGRA scenario with $\tan \beta = 5$, $\mu < 0$ and $A_0 = 0$; the same model is used to interpret the results in the $(M_{\tilde{\chi}_1^0}, M_{\tilde{\ell}})$ plane, as shown in figure 23. Note that in this case the numerical values assumed by the $R$–parity violating coupling affect the slepton (or sneutrino) lifetime$^{43}$ and its production cross-section; for this reason exclusion contours are functions of $\lambda^\prime_{211}$.

For light sleptons and sneutrinos, the phase space allowed to the decay products of the $R$–parity conserving processes $\tilde{\ell}^\pm \rightarrow \ell^\pm \tilde{\chi}_1^0$, $\nu_l \tilde{\chi}_1^±$ and $\tilde{\nu}_l \rightarrow \nu_l \tilde{\chi}_1^0$, $\ell^± \tilde{\chi}^0_1$ is small; as a result, cascade decays leading to the LSP with subsequent $R$–parity violating decays of the LSP lead to the production of leptons and jets with low $p_T$, which are likely to fail identification and reconstruction criteria. In this case, which is of particular interest for third-generation sleptons ($\tilde{\tau}$ and $\tilde{\nu}_\tau$) affected by large mixing, leptonic decays of the slepton or sneutrino are possible if a second $R$–parity violating coupling, $\lambda_{ijk}$, assumes a non trivial value$^{44}$. Direct production of a single stau or tau sneutrino can be achieved in $R$–parity violating mode with non trivial $\lambda^\prime_{211}$; in this case, a charged lepton pair is required in the final state in order to suppress the abundant $W$+jets background. This request effectively rules out the production of tau sleptons, always leading to a $\tau^± + \vec{E}_T$ signature. Among the $R$–parity violating decay modes of a light sleptons.

$^{43}$ Throughout all this analysis $\lambda^\prime_{211} > 0.03$, which grants prompt decays of the lightest neutralino.

$^{44}$ Hadronic decays of the slepton or sneutrino are indeed possible via a non trivial $\lambda^\prime_{ijk}$ coupling; however, these decays are affected by large backgrounds at hadron collider.
tau sneutrino, the most appealing combination of oppositely-charged leptons is given by an electron and a muon (corresponding to $\lambda_{132} = 0$), due to the low SM process contributions to this channel\(^{45}\).

The CDF analysis [35] investigates the production of a single tau sneutrino decaying to an electron-muon pair; $R$–parity is considered violated by the $\lambda'_{311}$ and $\lambda_{132}$ couplings only. Events are selected by requiring an isolated central electron-muon pair with $E_T(p_T) > 20 \text{ GeV}/c$; the two leptons are required to be compatible with originating from the same primary vertex. The dominant $Z \rightarrow \tau^+ \tau^-$ background can be removed by defining a signal region for $M_{e\mu} > 100 \text{ GeV}/c^2$; at this stage, major background contributions come from di-boson and $t \bar{t}$ production. SM expectations are then compared to data in the $50 < M_{e\mu} < 100 \text{ GeV}/c^2$ control region – where signal is excluded by previous searches – and found to be in good agreement.

**Figure 24.** Search for a resonant scalar neutrino using 344 pb\(^{-1}\) of data collected by the CDF detector. 95\% confidence level upper limits on $\lambda'_{311}$ (left) and $\lambda_{132}$ (right) as functions of $M_{e\mu}$ [35].

A Bayesian approach is used to set an upper limit on $\sigma(p\bar{p} \rightarrow \tilde{\nu}_\tau \rightarrow e\mu)$ in the hypothesis of absence of signal; the results are then translated in the 95\% confidence level exclusion contours in the $(M_{e\mu}, \lambda'_{311})$ and $(M_{e\mu}, \lambda_{132})$ planes shown in figure 24.

## 5. Indirect searches

With no significant deviations from the SM in any of the direct searches, signs of SUSY can be investigated in processes where new particles enter only as virtual states; an excellent opportunity for these studies is offered by processes that are rare within the SM.

In the SM, flavour changing neutral currents are strongly suppressed and can occur only through higher order diagrams; the decay rates of the processes $B_s \rightarrow \mu^+ \mu^-$ and $B_d \rightarrow \mu^+ \mu^-$ are proportional to $|V_{ts}|^2$ and $|V_{td}/V_{ts}|^2$ respectively. The SM predicts $BR(B_s \rightarrow \mu^+ \mu^-) = (3.42 \pm 0.54) \times 10^{-9}$ and $BR(B_d \rightarrow \mu^+ \mu^-) = (1.00 \pm 0.14) \times 10^{-10}$, two orders of magnitude smaller than the current experimental sensitivity. However, new physics may significantly increase these decay rates: SUSY, for instance, predicts a possibly strong enhancement of the branching ratio of $B_s,d \rightarrow \mu^+ \mu^-$ by introducing additional contributions mediated by neutral Higgs bosons (depicted in figure 25) and whose decay rates are proportional to $\tan^6 \beta$ [36]. This can result in branching ratios that can be three orders of magnitude larger than the SM prediction; in this perspective, the observation of such decays would be a strong – although indirect – evidence supporting SUSY.

\(^{45}\) The oppositely-charged, same flavour di-lepton final states are investigated by searches targeting the production of heavy resonances; the results of these (essentially model-independent) searches are interpreted also in terms of production of spin-0 particles (like sneutrinos) and used to set limits on the corresponding cross-sections [34].
Figure 25. Supersymmetric correction to the process $B_s \rightarrow \mu^+ \mu^-$.  

The CDF search for the $B_{s,d} \rightarrow \mu^+ \mu^-$ rare decays [37] has interested the first 364 pb$^{-1}$ of Run II data$^{46}$. Events are preselected by identifying two central muons of opposite charge with $p_T > 2.2$ GeV/c; these are used for reconstructing the B hadron, which is required to have a transverse momentum of at least 4 GeV/c and a proper decay length significance$^{47}$ of at least 2. A likelihood function is used to separate di-muons originated from the decay of a particle with lifetime from prompt di-muons; suitable discriminating variables are identified as the invariant mass of the di-muons, the proper decay length $\lambda$ of the B hadron, its isolation$^{48}$ and the three-dimensional angle between the B hadron momentum and its decay axis.

The $B_s \rightarrow \mu^+ \mu^-$ branching ratio is normalized to the number of $B^+ \rightarrow J/\Psi K^+ \rightarrow \mu^+ \mu^-$ decays in order to disentangle the analysis from the $b$ production cross-section. Background estimations are cross-checked by means of suitable control data samples; no statistically significant discrepancies are observed. No signal events are selected within the signal boxes shown in figure 26; this allows to exclude $BR(B_s \rightarrow \mu^+ \mu^-) > 2.0 \times 10^{-7}$ and $BR(B_d \rightarrow \mu^+ \mu^-) > 5.1 \times 10^{-8}$ at the 95% confidence level.

The CDF and DØ analyses have been combined [39], leading to the upper limits $BR(B_s \rightarrow \mu^+ \mu^-) < 1.5 \times 10^{-7}$ and $BR(B_d \rightarrow \mu^+ \mu^-) < 4.0 \times 10^{-8}$ at the 95% confidence level. According to [36], mSUGRA models predicting $BR(B_s \rightarrow \mu^+ \mu^-) > 10^{-7}$ require $\tan \beta > 40$;

46 A similar analysis has been performed by DØ with 300 pb$^{-1}$ [38].
47 Defined as $\lambda/\sigma$, where $\sigma$ is the uncertainty on $\lambda$ ($\lambda = c \cdot \tau$).
48 Defined as the fractional $p_T$ of the $B^0 \rightarrow \mu^+ \mu^-$ within a $\Delta R = 1$ cone around the direction of flight of the B hadron.
therefore, the search for the $B_s \rightarrow \mu^+ \mu^-$ rare decays sets the upper limit $\tan \beta \lesssim 40$ if a mSUGRA scenario is assumed.

6. Conclusions

Despite the excellent agreement between experiments and Standard Model predictions, the road to a unified quantum field theory describing the fundamental particles and their interactions is still long and uncertain. No direct evidence of new phenomena has been observed so far. However, indirect indications based on experimental data and supported by theoretical arguments suggest that the we may be close to the frontier of a new world; if these conjectures were true, then the discovery of new physics should be imminent. Several possible scenarios of what could look like the universe of particle physics beyond the electroweak scale have been proposed in the past few years: among them, Supersymmetry still provides one of the most credited grounds to extend the Standard Model to higher energies.

The lecture focuses on the many challenges that the Tevatron collaborations are facing in the search for Supersymmetry; the indications suggesting that new physics phenomena might arise at relatively low energies and be even connected with the $SU(2)_L \times U(1)_Y$ symmetry breaking make the challenge even more enticing. This document is a review of the most recent analyses performed by CDF and DØ in their quest for Supersymmetry and is aimed at showing how experiments contribute to the development of theoretical models by providing increasingly tighter constraints on the behaviour of Nature. Particular attention has been devoted to the explanation of the experimental signatures and to the description of the techniques that are commonly used by these analyses.

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