Signatures of Supersymmetry at the LHC

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We briefly review the aspects of supersymmetry (SUSY) and SUSY breaking necessary for appreciating the prospect of discovering supersymmetric particles at the Large Hadron Collider (LHC), avoiding technicalities as far as practicable. After a brief survey of the generic signatures of SUSY at hadron colliders, we review the specific LHC signals of different models of SUSY breaking with emphasis on the *smoking gun* signals of each model, if any.

1. Introduction

Supersymmetry (SUSY) is a novel symmetry that relates bosons and fermions. Under a SUSY transformation a spin half electron, for example, changes into its superpartner: a spin-zero selectron. Historically SUSY was introduced purely for its theoretical elegance [1]. Quite remarkably subsequent researches revealed that it not only removes several shortcomings of the existing theories of particle physics but predicts a whole host of new phenomena which can be tested by experiments at high energy or by cosmological experiments [2]. The discovery of the superpartners as well as their properties is, therefore, a high priority programme for the physics at the Large Hadron Collider (LHC).

The Standard Model (SM) of particle physics (see below) has been very successful in explaining all the experimental data accumulated over the years. The list of such experiments is indeed impressive. It includes typical laboratory based low energy experiments like atomic parity violation [3] as well as accelerator-based experiments carried out at the highest energy attainable so far. The latter includes experiments at the Large Electron Positron (LEP) collider (CM energy approximately 200 GeV) at CERN which were completed a few years ago[1] and the ones at the $p\bar{p}$ collider Tevatron (CM energy approximately 2000 GeV or 2 TeV) which have been running at the Fermilab for the last 20 years or so [4].

One sector of the SM proposed by Glashow, Salam and Weinberg contains the unified theory of electroweak (EW) interactions. This sector known as the EW sector, contains three generations of spin-half quarks and leptons. The particles belonging to the first-generation are listed in Table 1. The symbols for the second-generation of particles are obtained by the replacements u (up) → c (charm), d (down) → s (strange) and e (electron) → μ (muon). The third-generation particles are denoted by u → t (top), d → b (bottom) and e → τ (tau-lepton).

Names	spin 0	spin 1/2
squarks, quarks	$(\widetilde{u}_L \;\; \widetilde{d}_L)$ $\widetilde{u}_R$ $\widetilde{d}_R$	$(u_L \;\; d_L)$ u_R d_R
sleptons, leptons	$(\widetilde{\nu}_e \;\; \widetilde{e}_L)$ $\widetilde{e}_R$	$(\nu_e \;\; e_L)$ e_R

Table 1. Quarks and leptons belonging to the first-generation of the standard model and their superpartners.

There are also spin-one force carriers or gauge bosons. They are the Ws $(W^{\pm}, W^0)$ and the B^0 (Table 2). The physical gauge bosons, which mediate EW interactions among the particles, are $W^{\pm}$, Z^0 and the photon (γ), the last two being orthogonal linear combinations of W^0 and B^0.

In addition there is the spin-zero Higgs boson[2] which generates the masses of all fermions and gauge bosons of the SM via spontaneous symmetry breaking. The masses of the heavier particles in the SM are $\sim$ 100 GeV. Qualitatively speaking, this mass or energy scale characterises the energy scale of the SM better known as the EW scale.

The remaining sector of the SM describes the strong

[1] For details see the article by A. Gurtu in this volume.

[2] For more about the Higgs bosons see the article by Djouadi and Godbole in this volume.

Names	spin 1/2	spin 1
gauginos, gauge bosons	$(\widetilde{W}^{\pm}\ \ \widetilde{W}^{0})$	$(W^{\pm}\ \ W^{0})$
	$\widetilde{B}$	B
	$\widetilde{g}$	g

Table 2. Gauge bosons of the standard model and their superpartners.

interaction among the quarks only. This interaction is mediated by a set of eight gauge fields collectively called the gluons (g). The theory of quarks and gluons is known as quantum chromodynamics (QCD) [3]

The supersymmetric partners of the particles belonging to the SM – popularly known as the superparticles or simply the sparticles are also shown in Tables 1 and 2. The Higgs sector of any supersymmetric model is more complicated than simply adding the superpartner of the single neutral Higgs boson in the SM. Since the Higgs sector is not our main concern, we refer the reader to the article by Djouadi and Godbole in this volume for the details.

The triumph of SUSY began by alleviating a pathological feature of the Higgs sector of the SM now known as the mass hierarchy problem. As in any quantum field theory, the mass of the Higgs boson in the SM receives quantum corrections called self energy corrections. It is well known that such corrections, except for very exceptional cases, turn out to be infinite because the self energy is determined by certain integrals which diverge. Of course in a renormalisable field theory, like the SM, a finite result can be obtained by the renormalisation prescription. The infinite self energy of the Higgs boson by itself is, therefore not a problem technically.

The problem begins when we realise that the SM cannot be the last word on nature. Surely at least the gravitational interactions, which cannot be embedded in the SM or in any renormalisable field theory for that matter are always there. The energy at which the gravitational interactions become strong is the Planck scale ($M_P \sim 10^{18}$ GeV). Since the SM certainly is not the only relevant theory at this very high energy, the field theoretic techniques applicable at lower energies are not reliable any more. The sensitivity of the magnitude of the Higgs boson self energy to the new energy scale can be estimated by putting M_P as the upper limit of the above divergent integral. The result is a disaster: the correction turns out to be proportional to M_P^2!

On the other hand although the Higgs boson is yet to be discovered there are strong indications from the LEP experiments that its mass cannot be much more than a few hundred GeV. This run away behaviour of the Higgs boson mass in the presence of new physics at a energy scale much higher than the characteristic energy scale of the SM, is known as the hierarchy problem [5]. It should be emphasised that the self energies of the fermions or the gauge bosons in the SM are free from this problem because of appropriate symmetries of the theory. On the other hand there is no corresponding symmetry which protects the Higgs boson mass. A solution of this hierarchy problem in terms of a new symmetry is, therefore very welcome. The self energy of the Higgs boson (ϕ) can be diagrammatically represented, for example, by the Feynman diagram in Fig. 1a, where a fermion (f), say a quark, circulates in the loop. Suppose there is an additional contribution to the self energy from a diagram with a spin-0 boson ($\tilde{f}$) having exactly the same mass as f circulating in the loop (Fig. 1b). Moreover let the couplings of the boson and the fermion with ϕ denoted by λ_s and λ_f respectively, be related in a certain way. A little experience in QFT would then tell us that the two contributions would neatly cancel and the dangerous contribution proportional to M_P^2 disappears. This is exactly what SUSY provides for us: corresponding to every fermionic (bosonic) particle there is a bosonic (fermionic) superpartner of exactly the same mass! Moreover SUSY relates the coupling strengths of any particle and its superpartner in such a way that each and every dangerous contribution to the self energy of ϕ in the SM is canceled. The solution of the hierarchy problem [6] triggered a renewed interest in SUSY. Realistic models of SUSY like the minimal supersymmetric extension of the SM (MSSM) were constructed and testable predictions of the models were computed with great enthusiasm.

SUSY cannot be an exact symmetry. Had there been a spin zero selectron as light as the electron it would have been discovered long before. Thus SUSY must be a broken symmetry with the sparticle masses significantly larger than the particle masses. However even in this case the hierarchy problem provides an important guideline. If f and $\tilde{f}$ have different masses, as they should if the symmetry is broken, the dangerous $\sim M_P^2$ terms in the Higgs boson self energy still cancel. But the remaining dominant contribution to the Higgs boson self energy is $\sim m_{\tilde{f}}^2 ln(M_P/m_{\tilde{f}})$, where $m_{\tilde{f}}$ is assumed to be much larger than m_f. The logarithmic dependence on M_P obviously softens the hierarchy problem. But if $m_{\tilde{f}}$ is much larger than a few TeV,

[3]See the article by Mathews and Ravindran in this volume for further details and references.

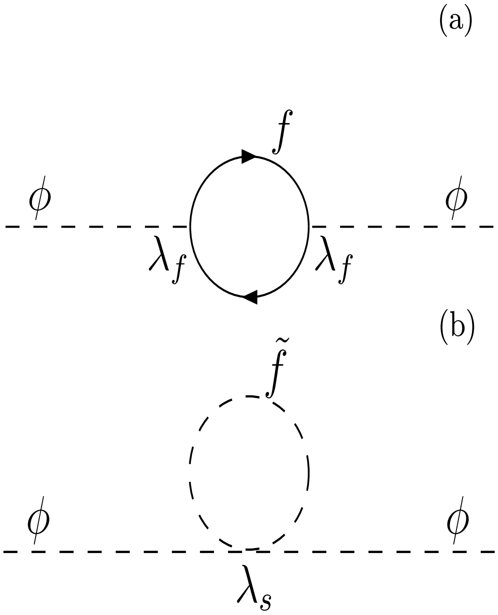

Figure 1. Cancellation of the dangerous contribution to the Higgs boson self energy in a supersymmetric theory (see the text for details)

one lands into a milder version of the hierarchy problem again and the Higgs boson mass tends to run away from the EW scale. This observation kindles the hope that if SUSY exists in nature, the sparticle masses cannot be much larger than the corresponding particle masses and hence the sparticles are likely to be accessible to the LHC with an unprecedented CM energy of 14 TeV.

The most general MSSM consistent with the symmetries of the SM turns out to be unacceptable since it contains baryon and lepton number violating interactions. Such interactions would spoil, for example, the stability of the proton. One can, however impose additional symmetries to get rid of these dangerous interactions. These symmetries introduce a new multiplicatively conserved quantum number called R-parity. It turns out that the SM particles have R-parity = 1 while the sparticles have R-parity = −1. A direct consequence of this new symmetry is that if a high energy collision of two particles (f_1 and f_2) produces sparticles, they must be produced in pairs, i.e.

$$f_1 + f_2 \rightarrow \tilde{f}_1 + \tilde{f}_2 \,, \tag{1}$$

where $\tilde{f}_1$ and $\tilde{f}_2$ are sparticles.

Another important consequence of the new symmetry is that any sparticle must decay into another sparticle along with one or more particles. This in turn implies that the lightest supersymmetric particle (LSP) must be stable. Now the LSP is very likely to be the superpartner of a neutral weakly interacting SM particle (Section 3.1). In that case this particle, if produced in colllider experiments, either directly or through the decays of heavier sparticles, will escape detection, resulting in apparent imbalance of energy and momentum. The presence of large missing energy in experiments involving the production and subsequent decay of sparticles is regarded as the hallmark of R-parity conserving SUSY. In this article we shall always assume R-parity conservation.

It is well-known that a large fraction of our universe is made of dark matter and dark energy. We know that the present universe is filled with left-overs from the evolution of the early universe following the Big Bang. The best known example of such relic particles is the cosmic microwave background. If the dark matter is also made of left-over particles of a certain type, they must be stable and weakly interacting (otherwise they would have been discovered long before). The LSP in the R-parity conserving SUSY appears to be tailor made as a dark matter candidate. Of course in order to account for the observed dark matter relic density quantitatively it must have mass and interaction strength in certain ranges. This pushes search for SUSY up by several rungs in the list of priorities for the LHC experiments. The discovery of SUSY at the LHC will be a great progress in our understanding of nature. The subsequent measurement of the properties of the sparticles - the LSP in particular - and quantitative verification of the conjecture of SUSY dark matter will be a cherry on the pudding[4].

The Achilles heel of any theory based on SUSY, however is the absence of a universally accepted guideline for SUSY breaking consistent with the presently available constraints on the sparticle masses. The implementation of spontaneous breakdown of SUSY in a realistic model without introducing additional unknown parameters would have been a great step forward. However, several attempts have failed to produce a consistent mass spectrum for the sparticles. For example, the SUSY breaking order parameter cannot belong to any of the MSSM supermultiplets, consisting of e.g. squark and quarks, because of certain mass sum rules. One such sum rule requires that the masses of the quarks and the squarks belonging to a generation satisfy the relation $m^2_{\tilde{u}_1} + m^2_{\tilde{u}_2} - 2m^2_u + m^2_{\tilde{d}_1} + m^2_{\tilde{d}_2} - 2m^2_d = 0$ (the

[4] For further details on dark matter and SUSY see the article by Baer and Tata in this volume.

notation will be further clarified in section 3). This is of course ruled out by experiments. However, this equation holds only at tree level and only for renormalisable theories. Then there exists a possibility that spontaneous SUSY breaking occurs in a sector which couples to the MSSM sector only via loops or via non-renormalisable interactions. These possibilities will be discussed in some details in the following sections.

At this stage the only option left is to add certain soft SUSY breaking terms to the theory without spoiling the symmetries of the SM and the renormalisability of the theory. Soft SUSY breaking masses - often called soft masses - which introduce the mass difference between particles and their superpartners are examples of such parameters. The number of new unknown parameters so added to the most general MSSM, however turns out to be more than hundred! The list is indeed too long for a theory claiming to be the fundamental description of nature.

Apart from the issue of aesthetics there are practical problems too. Arbitrary, random soft parameters in the MSSM Lagrangian potentially lead to flavour or CP violating processes with large probabilities. However, these processes are known to be severely suppressed even in comparison with the usual weak processes like μ decay and π decay. In order to understand this issue one should recall that once SUSY is allowed to be broken arbitrarily the masses of the sparticles also become arbitrary. This may lead to flavour-changing neutral current (FCNC) induced processes, such as $\mu \to e\gamma$, K^0–$\bar{K}^0$ mixing etc. The latter process, for example, can be suppressed if the masses of the $\tilde{u}$ and $\tilde{c}$ squarks happen to be degenerate to a very good approximation. Other FCNC processes can be similarly suppressed by requiring suitable pairs of other sparticles to be almost mass degenerate. Large CP violating effects can be avoided by assuming that the soft parameters do not introduce new complex phases.

From the physicist's point of view, however such degeneracies without an underlying symmetry is not satisfactory at all. One would like to find either a symmetry or a dynamical reason which would explain why apparently different parameters have nearly the same values. For example, these dangerous FCNC induced effects in the MSSM can be satisfactorily removed if one can show that the soft SUSY breaking parameters like the squark masses are universal because the dynamics at the SUSY breaking scale, which is much higher than the EW scale, is flavour blind (i.e. it does not distinguish between squarks of different types). It has also been noticed that if physics at some higher scale is indeed responsible for SUSY breaking, then a few other puzzles plaguing non-supersymmetic theories become non-issues in supersymmetric theories. Several models belonging to this class will be summarised in sections 2 and 4.

From the last paragraph it is clear that the sparticle mass spectrum and consequently, the production and decay of sparticles at the LHC are by and large model dependent. Thus the discovery of some generic SUSY-like signal, e.g. the observation of events having much larger missing energy than typical SM processes, though spectacular, will not automatically herald the nature of the underlying SUSY theory. Painstaking reconstruction of the sparticle masses, their spins and other properties will be required to identify the specific model of SUSY breaking. This makes the search for SUSY at the LHC even more exciting and challenging. Fortunately there are "smoking gun signals" of some of these models. The observation of these would strongly hint to the specific model without waiting for the full reconstruction of the model parameters. In section 3, we shall first summarise the generic SUSY signals. Then we shall take up the issue of signatures of specific models including the smoking gun signals, if any, in sections 3 and 4.

2. SUSY Breaking at a High Scale

In quantum field theory parameters like mass and charge (or any other coupling constants) are energy dependent. This prediction has indeed been verified by experiments. Low energy experiments, e.g. have measured the fine structure constant (α) very accurately and its value is well known ($\approx 1/137$). However the value of α as measured by the LEP experiments at a CM energy of approximately 100 GeV is significantly larger ($\approx 1/128$). Another well known example is the variation of the coupling strength of quark-gluon interactions in QCD. This coupling, however decreases with increasing energy and is rather small for quark-gluon interactions at very high energies (footnote 3).

The variation of the masses and coupling constants with energy in a particular theory can be theoretically studied by a set of coupled differential equations known as the renormalisation group (RG) equations [7]. The predictions of the RG equations of various theories - QCD in particular - have been verified by a large number of experiments operating at different energies.

That SUSY may be more attractive in presence of new physics at much higher energies is indicated by the supersymmetric generalisation of any grand unified theory (GUT) [8] of the simplest type [2].[5] The cou-

[5]These are the so called grand desert type theories in which there

pling constants of strong, electromagnetic and weak interactions measured at currently attainable energies of course have widely different magnitudes. However the weaker (stronger) couplings increase (decrease) with energy (see the examples discussed above). This suggests the interesting possibility that these interactions have a common strength at a much higher energy called the GUT scale ($M_G \sim 10^{16}$ GeV). This coupling constant unification, however fails in non-supersymmetric GUTs of the simplest kind. It is well-known that unification can be achieved in supersymmetric GUTs of the simplest variety provided the sparticle masses are ~ 1 TeV. The essential reason is that the sparticles contribute to the RG evolution of the couplings and these contributions are important to make them unify at the right scale (M_G).

The idea that the impact of physics at very high energies, much higher than the currently attainable energies, on physics at currently available energies can be studied via the RG equations, have led to several interesting theories of SUSY breaking. The main point is that certain interactions at very high energies may introduce a rather simplified pattern of SUSY breaking involving only a few free parameters. The prediction for the low energy spectrum may then be obtained by the RG equations. Alternatively one may determine the masses of the sparticles from experiments at a lower energy. One can then evolve these masses to a high scale via the RG equations and check whether they exhibit any special feature. This information may eventually reveal the underlying SUSY breaking model.

The models of high scale SUSY breaking have a generic feature. Each consists of a 'hidden sector' which does not interact with the 'observable sector' consisting of the particles belonging to the MSSM. The hidden sector particles are assumed to be too heavy to be observed in accelerator based experiments in the near future. SUSY is broken in the hidden sector spontaneously. The central question of SUSY has now been changed : "How does the *observable sector* know about SUSY breaking?" The mechanism of mediating SUSY breaking from the hidden sector to the observable sector via some 'messenger fields' which couples to both the hidden and the observable sectors are different in different models. This leads to characteristic MSSM soft terms and, consequently, sparticle masses and collider signals in different models.

The most popular example of mediating SUSY breaking from the hidden to the visible sector is via gravitational interactions which play the role of the messenger. Here the new physics enters near the Planck Scale (M_P). If SUSY is broken in the hidden sector by a vacuum expectation value (VEV) $\langle F \rangle$, where F is a hidden sector field, then the soft terms in the visible sector should be $m_{\rm soft} \sim \langle F \rangle / M_{\rm P}$, by dimensional analysis. Hence, if we demand $m_{\rm soft}$ of the order of a TeV, then the scale associated with the origin of SUSY breaking in the hidden sector should be $\sqrt{\langle F \rangle} \sim 10^{10}$ or 10^{11} GeV.

One can argue, somewhat naively, that since gravitational interactions does not distinguish between different scalar particles in the MSSM, all such particles (the squarks, sleptons and the Higgs bosons) acquire the same mass (m_0) at a high scale via this mechanism. This is already encouraging because a common mass of the sfermions belonging to the first two generations strongly hints towards a natural mechanism for the suppression of the dangerous FCNC processes discussed in the introduction. A similar argument would indicate that the soft breaking masses for all the spin-1/2 gauge fermions or the gauginos would be the same ($m_{1/2}$). Admittedly this argument is oversimplified. In order to realise universal masses for the scalars and the gauginos one needs additional simplifying assumptions for hidden sector interactions. This model is popularly known as the minimal supergravity model or the mSUGRA model [9].

In the mSUGRA model, the free parameters at high energy, commonly chosen as the GUT scale are

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

where m_0 and $m_{1/2}$ are the common mass of the scalars and gauginos respectively. The parameter A_0 is another soft SUSY breaking parameter known as the trilinear coupling. The ratio of the VEVs of the two neutral Higgs bosons (footnote 2) is denoted by tan β. The parameter μ respects SUSY and is commonly referred to as the Higgsino mass parameter as it contributes to the masses of the superpartners of the Higgs boson. The masses of sparticles and couplings at the energy scale of experimental interest can be obtained via the RG evolutions of these parameters [10]. Some example of masses at the EW scale are

$$m_{\tilde{g}} \approx 2.7 m_{1/2}\,, \tag{2}$$

where $M_{\tilde{g}}$ is the mass of the gluino, the superpartner of the gluon. The squark masses of the first two generations are

$$m^2_{\tilde{u}_L} \approx m_0^2 + 4.9 m^2_{1/2}, \quad m^2_{\tilde{u}_R} \approx m_0^2 + 4.6 m^2_{1/2}\,. \tag{3}$$

It turns out that the up-type squarks belonging to the second generation ($\tilde{c}_L$ and $\tilde{c}_R$) have exactly the same

is no new physics between the EW scale and the GUT scale M_G.

mass. The masses of the down type squarks ($\tilde{d}_L$, $\tilde{d}_R$, $\tilde{s}_L$ and $\tilde{s}_R$) have the same mass as the corresponding up type squark except for relatively small EW symmetry breaking corrections. The mass degeneracy between u and c type squarks and that between d and s type squarks are sufficient to suppress the FCNC processes. This is one of the attractive features of the mSUGRA model.

The numerical coefficients in Eqs. (2) and (3) are determined by the gauge interactions of the sparticle concerned and the universal strong interaction determines their magnitudes to a large extent although there are sub-dominant EW contributions. This in fact is the main reason for the near degeneracy of all squarks belonging to the first two generations.

In principle the Yukawa interactions of the fermions in the SM also contribute to the RG evolutions of the soft masses of their superpartners. But these interactions proportional to the fermion masses are practically negligible for the light quarks of the first two generations. However the situation is very different for the third-generation. For example, large m_t (the top quark mass) ensures that the contribution of the top quark Yukawa coupling to the evolution of the masses, e.g. of the top squarks ($\tilde{t}_L$ and $\tilde{t}_R$) may make these sparticles significantly lighter than the squarks belonging to the first two generations. The experimental constraints on the FCNC processes mediated by the $\tilde{t}$ squarks are not very stringent and relatively light $\tilde{t}_{L,R}$ are consistent with the data.

We recall that in the SM a negative mass squared parameter of the Higgs boson is introduced rather arbitrarily. This leads to the spontaneous breakdown of EW symmetry which yields positive masses to all particles of the SM including the Higgs boson. In mSUGRA the Higgs mechanism occurs naturally. The mass squared of the Higgs boson (m_0^2) is positive at the GUT scale. However the subsequent RG evolution makes it negative at the EW scale. This triggers spontaneous break down of EW symmetry. This radiative EW symmetry breaking mechanism [11] adds one more feather to the cap of the mSUGRA model.

In the next section we shall review the expected signatures of the mSUGRA model at the LHC and in section 4 we shall do the same for others models of SUSY breaking.

3. SUSY Searches at the LHC

Two collaborations ATLAS and CMS[6] will independently look for signals of SUSY and other new physics candidates, if any, with all purpose detectors. At the LHC, two protons each of energy 7 TeV, will collide head on at a centre of mass energy $\sqrt{S}$ = 14 TeV and they will break up into constituents of the proton, namely, gluons, quarks and antiquarks, collectively known as partons. It is mainly the strong interaction of the partons via QCD, which leads to observable processes with large cross sections. However as we shall see below, the EW interactions of quarks and antiquarks also produce important signals al beit with smaller cross sections.

At the LHC the particles of the SM like top quarks, gauge bosons (W,Z), etc. will be produced in very large numbers. The sparticles, if they exist will also be produced, but with somewhat smaller cross sections since they are likely to be significantly heavier than the corresponding particles. Separating the signatures of sparticle production from the huge SM background is the main challenge confronting the SUSY searchers at the LHC.

In order to fully appreciate the mass reach of the LHC, let us briefly recollect the sparticle mass limits from LEP and Tevatron Run II – the best limits obtained to date. The highest CM energy attained by LEP was 209 GeV and the mass limts on most of the sparticles accessible to LEP were extended almost to the kinematic limit (i.e. ≈ 100 GeV) in the clean environment of an $e^+ - e^-$ collider[7] [12,13].

The CDF and D0 collaborations[8] have been looking for the sparticles since the dawn of the Tevatron experiments nearly 20 years ago [13]. As noted earlier these limits are usually model dependent. Assuming that there are five flavours of squarks of L and R type and each has approximately the same mass as the gluino ($m_{\tilde{q}} \approx m_{\tilde{g}} = \tilde{m}$) the CDF collaboration obtained the limit $\tilde{m} > 392$ GeV. For heavier squarks $m_{\tilde{q}} = 600$ GeV, the gluino mass limit is $m_{\tilde{g}} > 280$ GeV [14]. The D0 collaboration obtained similar limits [15].

3.1. Sparticle Production at the LHC

If the mass of the strongly interacting sparticles (the gluinos ($\tilde{g}$) and the squarks ($\tilde{q}$), the superpartners of gluons and quarks respectively) are within the kinematic reach of the LHC, which roughly corresponds to $m_{\tilde{g}}, m_{\tilde{q}} \leq 2.5$ TeV, they will produce observable signals. As already noted in the introduction these sparticles

[6] A large number of Indian physicists from different research institutes and universities are members of the CMS collaboration.

[7] The limits on the LSP and the sneutrino masses are more model dependent and may be well below the kinematic limit in some scenarios.

[8] A large number of Indian physicists from different universities and institutes have been participating in the D0 experiment.

will be produced in pairs according to R-parity conserving SUSY models. At the partonic level some typical reactions are

$$q\bar{q},\ gg \to \tilde{g}\tilde{g}, \tilde{q}\tilde{q}^*, \tag{4}$$

$$qq \to \tilde{q}\tilde{q}, \tag{5}$$

$$qg \to \tilde{q}\tilde{g}, \tag{6}$$

$$q\bar{q},\ gg \to \tilde{t}_1\tilde{t}_1^*. \tag{7}$$

Since the dynamics of QCD at high energies is fairly well known the cross sections of the above processes are essentially controlled by the masses of the final state sparticles [16]. These masses, in turn are determined by the SUSY breaking mechanism. Signals of sparticle production is, therefore somewhat model dependent. In this section we shall, however concentrate on some generic features of sparticle production which are fairly model independent.

The inclusion of the process in Eq. (7), which is nothing but a special example of $\tilde{q}\tilde{q}^*$ production, deserves further clarifications. In the MSSM, the third-generation squarks ($\tilde{q}_L, \tilde{q}_R$; Q = t,b) require special attention. After soft SUSY breaking, many of the sparticles in Tables 1 and 2 do not represent the physical states. In fact suitable mixtures of the states listed in the above tables correspond to the observable states expected to be produced at the LHC or other future colliders. Consider the L and R type sfermions $\tilde{f}_L$ and $\tilde{f}_R$, which are the superpartners of the chiral fermions f_L and f_R, where f represents any quark or lepton flavour in the SM. These sfermions can mix via the mass squared matrix

$$M^2_{\tilde{f}} = \begin{bmatrix} m^2_{LL} & m_f\tilde{A}_f \\ m_f\tilde{A}_f & m^2_{RR} \end{bmatrix}, \tag{8}$$

where the matrix elements are given by

$$\begin{aligned} m^2_{LL} &= m^2_f + m^2_{\tilde{f}_L} + (I^3_f - q_f\sin^2\theta_W)\cos 2\beta\, M^2_Z \\ m^2_{RR} &= m^2_f + m^2_{\tilde{f}_R} + q_f\sin^2\theta_W\ \cos 2\beta\, M^2_Z \\ \tilde{A}_f &= A_f - \mu(\tan\beta)^{-2I^3_f}. \end{aligned} \tag{9}$$

Here $m_{\tilde{f}_L,\tilde{f}_R}$ are the soft masses for the L and R type sfermions, m_f is the mass of the corresponding fermion (e.g. a quark or a lepton) and A_f is another SUSY breaking parameter. The parameter μ and tan β have already been defined. The other SM parameters appearing in Eqs. (8) and (9) are not particularly relevant for our present discussion. The effects of mixing depend on m_f in the off diagonal terms. These terms are therefore, much smaller than $m_{\tilde{f}_L,\tilde{f}_R}$ for the first two generations. Only for the mixing between the superpartners of the top quark the off diagonal terms in the mass matrix could be important due to large m_t (the soft breaking term A_t may also be large). This may lead to sizable splitting between two physical states $\tilde{t}_1$ and $\tilde{t}_2$.[9] In mSUGRA type models the top squarks may also be light due RG evolution of their masses (Section 2). For a large region of the parameter space the lighter physical state ($\tilde{t}_1$) can be significantly lighter than the other squarks due to these two effects. Therefore light $\tilde{t}_1$ pairs can be produced with substantial rates at the LHC.

In Fig. 2 the lowest order cross sections for gluino and squark pair production in various combinations are shown along with that for $\tilde{t}_1 - \tilde{t}_1^{\,*}$ pair production [17]. In this figure $\tilde{q}$ collectively represent degenerate L and R type squarks of all flavours except $\tilde{t}$. It is also assumed that $m_{\tilde{g}} = m_{\tilde{q}}$. Notice that the cross sections fall rapidly with the masses. For $m_{\tilde{g}}, m_{\tilde{q}}$ = 1 TeV, e.g. the total squark-gluino pair production cross section is $\sim$ 4–5 pb. Thus for an integrated luminosity of $\mathcal{L}$ = 10 fb^{-1} about 40,000–50,000 squark-gluino pairs will be produced.[10] For the same $m_{\tilde{t}_1}$, about 100 stop squark pairs are expected.

The existence of a light top squark may lead to several other important consequences. For example in some models $\tilde{t}_1$ may turn out to be the next-to-lightest supersymmetric particle (NLSP). This leads to the interesting possibility that the $\tilde{t}_1$ - LSP annihilation might have been the dominant dark matter producing process in the early universe (see the article by Baer and Tata in this volume for further details). Other important consequences of the light $\tilde{t}_1$ scenario will be reviewed later.

In addition to the dominant production processes, Eqs. (4, 5, 6 and 7), the EW gauginos - two charginos ($\tilde{\chi}^{\mp}_i$; i = 1–2) and four neutralinos ($\tilde{\chi}^0_j$; j = 1–4) may also be produced with smaller but observable rates. The indices i and j increase in the order of increasing mass of the sparticle. The charginos are linear

[9] The mixing between the superpartners of the bottom quark and the τ lepton may also be large especially if tan β is large.

[10] Luminosity is an important parameter of any collider experiment. The product of the cross section of a process and luminosity gives the number of events produced per second. Integrated luminosity is the luminosity accumulated over some interval of time. At the LHC $\mathcal{L}$ = 10 fb^{-1} will be accumulated in about two years (1 fb^{-1} = 1000 pb^{-1}) inspite of the low luminosity phase during the first two-three years. This should be compared with the rather optimistic goal of collecting 9 fb^{-1} at the Tevatron Run II by the end of 2010 after approximately nine years of running. In the high luminosity phase the LHC is expected to collect 100 fb^{-1} per year.

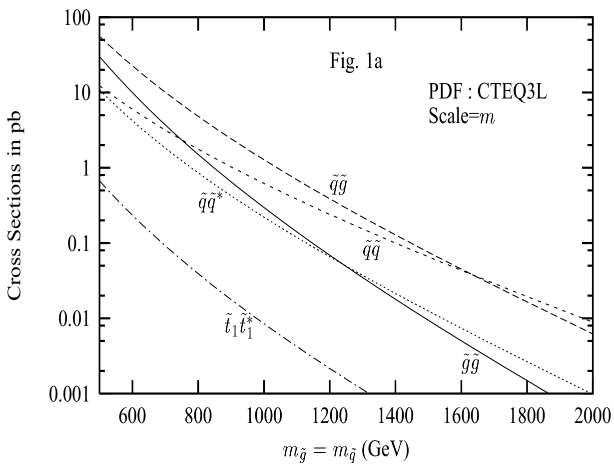

Figure 2. Gluino and squark pair production cross sections at the LHC, from [17]

combinations of $\tilde{W}^\pm$ (the superpartners of the charged W bosons) and $\tilde{H}^\pm$ (the superpartners of the charged Higgs bosons). On the other hand the neutralinos are linear combinations of $\tilde{W}_3$ (the superpartners of the neutral W boson), $\tilde{B}$ (the superpartners of the neutral U(1) gauge boson B) and $\tilde{H}_{1,2}$ (the superpartners of the two neutral Higgs bosons). In most models the lightest neutralino ($\tilde{\chi}_1^0$) is assumed to be the LSP.

They are produced by quark and anti-quark annihilations, mediated by photons, W/Z gauge bosons and squarks. The initial quark and, in particular, the antiquark flux in the proton are rather low at the LHC. Moreover the underlying interactions are EW in nature. As a result the production rates are comparatively low at the LHC. Although, in principle, all combinations of charginos and neutralions can be produced, the dominant ones are

$$q\bar{q}' \to \tilde{\chi}_1^\mp \tilde{\chi}_1^\pm, \quad \tilde{\chi}_1^\mp \tilde{\chi}_2^0, \quad \tilde{\chi}_1^0 \tilde{\chi}_2^0 \,. \tag{10}$$

It should be noted that the cross sections of the above processes are not controlled by the masses of the final state particles alone. The compositions of the charginos and neutralinos also determine the magnitude of cross sections. These compositions in turn are controlled by the parameters

$$M_1, M_2, \mu, \tan\beta \,, \tag{11}$$

where M_2, M_1 - the soft masses for SU(2) and U(1) gauginos respectively, μ and $\tan\beta$ are already defined. Obviously the cross sections for these processes are more parameter space dependent than that of squark-gluino production. As an example we note that, the cross section for the $\tilde{\chi}_1^\mp \tilde{\chi}_2^0$ chargino-neutralino pair production is at the sub pb level for $m_{\tilde{\chi}_1^\mp}, m_{\tilde{\chi}_2^0} \sim$200 GeV.

The associated production of $\tilde{g}$ and $\tilde{q}$ with charginos and neutralinos are also possible. For example,

$$gq \to \tilde{q}\tilde{\chi} \tag{12}$$

is an allowed process, where $\tilde{\chi}$ stands for either charginos ($\tilde{\chi}_i^\mp$; i = 1–2) or neutralinos ($\tilde{\chi}_j^0$; j = 1–4). The cross sections for the associated production are also dependent on the masses of produced sparticles and as well as the SUSY parameters in Eq. (11) due to the presence of one EW vertex. Typical cross sections of these processes are at a few pico barn level for the range of masses ∼200–300 GeV.

Finally it would be incomplete unless we mention the pair production of sleptons via the Drell Yan like mechanisms

$$q\bar{q} \to \tilde{\ell}\tilde{\ell}^*, \quad \tilde{\nu}_\ell \tilde{\nu}_\ell^* (\ell = \mathrm{e}, \mu, \tau) \tag{13}$$

$$q\bar{q}' \to \tilde{\nu}_\ell \tilde{\ell}^* + \tilde{\nu}_\ell^* \tilde{\ell} \tag{14}$$

mediated by Z/γ/W through s-channel. Because of the low initial quark flux and energy suppression as these interactions are s-channel processes, the production cross sections are rather tiny ($\sim \mathcal{O}$(fb)) for $m_{\tilde{\ell}}$, $m_{\tilde{\nu}} \sim$ 200 GeV.

One important point to be noted here is that the leading order (LO) cross sections as shown in Fig. 1 have been obtained in the lowest order in perturbation theory. They involve significant theoretical uncertainties due to the choice of the QCD scale and the parameterisation of parton distributions inside the proton. One can reduce some of these uncertainties by including more terms in the perturbative series for the cross section.

Since the dominant sparticle productions take place via strong interaction these higher order corrections are expected to be substantial. Generally, the relative importance of these corrections is given by the K-factor defined by the ratio K = $\frac{\sigma_{NLO}}{\sigma_{LO}}$, where σ_{NLO} is the cross section including the next-to-leading order terms. The K factors can be substantially large ranging from ∼ 1.1–1.6 [18]. In predicting the sparticle production rates one should take care of this K-factor appropriately.

3.2. SUSY Signals at the LHC

Once SUSY particles, dominantly gluinos and squarks are produced, they eventually decay to almost massless leptons and quarks along with the lightest neutralino ($\tilde{\chi}_1^0$) - the LSP - via a long decay chain involving the charginos and the neutralinos in the intermediate states. The long cascade decay processes, of course, depend on the relative magnitudes of $m_{\tilde{g}}$ and $m_{\tilde{q}}$. Accordingly, there are two possibilities: (a) $m_{\tilde{g}} \gtrsim m_{\tilde{q}}$ (b) $m_{\tilde{g}} \lesssim m_{\tilde{q}}$. The decay chain for case (a) is

$$\tilde{g} \to q\tilde{q}; \quad \tilde{q} \to q\tilde{\chi}_i^{\mp}, q\tilde{\chi}_j^0 \tag{15}$$

and in case (b) we have

$$\tilde{q} \to q\tilde{g}; \quad \tilde{g} \to q\bar{q}'\tilde{\chi}_i^{\mp}, q\bar{q}\tilde{\chi}_j^0. \tag{16}$$

Finally, the chargino and neutralino states decay to ordinary light fermions and the lightest neutralino

$$\tilde{\chi}_i^{\pm} \to f_1\bar{f}_2\tilde{\chi}_1^0, \tag{17}$$

$$\tilde{\chi}_j^0 \to f_3\bar{f}_3\tilde{\chi}_1^0. \tag{18}$$

Here f_i's stand for appropriate quarks and leptons. In Eqs. (15, 16) the $\tilde{g}/\tilde{q}$ decay branching ratios (BRs) depend on their masses and on the SUSY parameters space Eq. (11). Typically the BRs of the gluino are $\tilde{g} \to q\bar{q}'\tilde{\chi}_1^{\pm}$ (50–60%), $q\bar{q}\tilde{\chi}_2^0$ (35–30%), $q\bar{q}\tilde{\chi}_1^0$ (15–10%) for a wide region of the parameter space. Squark decays to charginos and neutralinos have almost the same relative rates. The lighter chargino ($\tilde{\chi}_1^{\pm}$) decays, Eq. (17), follow roughly the same branching ratio as W decays but the BRs of $\tilde{\chi}_2^0$ decay depend sensitively on the SUSY parameter space. For example, if the sleptons are light, i.e if $m_{\tilde{l}}$ is small and $\tilde{\chi}_2^0$ is gaugino dominated then, this neutralino will have relatively large BR to leptonic decay channels. The heavier chargino and neutralino states also decay to lighter particles accompanied by Higgs scalars [17]

$$\begin{aligned} \tilde{\chi}_i^0 &\to \phi\tilde{\chi}_k^0, (i > k)(\phi = h, H, A,), \\ \tilde{\chi}_2^{\pm} &\to H^{\pm}\tilde{\chi}_1^0 \\ &\to W^{\pm}\tilde{\chi}_1^{\mp}, \end{aligned}$$

where h, H and A are the three neutral Higgs bosons and $H^{\pm}$ is the charged Higgs boson in the MSSM. The sleptons which are produced via Eq. (14) also end up with light fermions and $\tilde{\chi}_1^0$, $\tilde{\ell} \to \ell\tilde{\chi}_1^0$.

As has already been noted the quarks and leptons in the final state will indicate an apparent imbalance of energy-momentum. This happens because the LSP escapes detection. In hadron colliders the momentum imbalance in the longitudinal direction, i.e. along the beam direction, cannot be tested as the fragments of colliding hadrons follow that direction and remain undetected. However in the transverse direction, i.e. perpendicular to the beam direction, the missing transverse energy ($\not{E}_T$), which is an experimental observable related to the missing transverse momentum, can be measured. This momentum imbalance is a characteristic signature of SUSY.

The quarks in the final stage of SUSY cascade decays cannot be observed directly. In fact each of them goes through non-perturbative processes like fragmentation and hadronisation and end up into a narrow stream of charged and neutral hadrons. These particles enter into a part of the detector called the hadron calorimeter and deposit energy in a cluster like forms. These clusters are called jets. These jets are the important observables in collider experiments. Photons and electrons (muons) are detected in the electromagnetic calorimeter (muon chambers).

Generically, any sparticle pair production followed by the cascade decays (Eqs. (15–18)) of the members of the pair at hadron colliders lead to the signal

$$\mathrm{n-leptons} + \mathrm{m-jets} + \not{E}_T; n, m = 0, 1, 2, 3..... \tag{19}$$

The long sequence of steps involved in production and detection of the SUSY signals are studied by Monte Carlo simulation which is implemented by packages called event generators [19]. Some of the commonly used and publicly available generators which can simulate a large number of SM or SUSY processes are HERWIG [20], ISAJET [21], PYTHIA [22] etc. In addition there are generators dedicated to specific processes. The prospects of detecting SUSY of various shades and stripes at the LHC experiments have been assessed by simulations using these generators. It should, however be noted that the non-perturbative processes mentioned above cannot be computed by using the standard techniques of field theory. Various models have been developed to study them. The event generators employ these models. Although many of the generators have so far been successful in explaining data up to Tevatron Run II, the reliability of these generators at the LHC has to be tested against the data on well understood SM processes collected during the early stages of the experiment. Only after a generator has been validated in the above way it can be used to study signals of new physics with confidence.

There are experimental challenges as well. For example, measuring $\not{E}_T$ in an experiment with the desired accuracy is a non-trivial task. Moreover, there are various sources of fake $\not{E}_T$. For example, detector effects like instrumental noise, hot or dead channels or cracks in between different parts of the detectors give some amount of fake $\not{E}_T$. Therefore, in order to establish SUSY signals based on $\not{E}_T$ convincingly, it is crucial to understand very precisely the fake $\not{E}_T$ and to develop a method for eliminating it.

In the LHC experiment a huge number of events will be produced. Only a small fraction of them will be of interest and will be stored. These are known as physics events. The ratio of the events stored to the total number of events produced is $\sim 10^{-10}$–10^{-9}. Even the physics events will mostly consist of the known particles like the W, Z, top quark, bottom quark, etc. If we are lucky and nature is kind to us, then some new particles may be produced with comparatively low rates. It is indeed a very challenging task to pick up the events signaling new physics out of the debris, called background or noise, mainly due to the SM processes.

The standard technique is to apply some selection criteria based on the kinematic properties of the signals from the new particles. As for example, a lepton, electron or muon, as well as a jet is selected off-line provided it transverse momentum (P_T) or energy (E_T) is above a certain minimum value which is quite common in the SM processes. This value is often referred to as a kinematical cut or simply a cut. As already discussed a strong cut on $\not\!\!E_T$ is a powerful weapon in the arsenal of a SUSY searcher. Of course, the magnitude of these cuts depend on the nature of the signal and and the SM background events. In addition there are many other kinematic observables, like angular correlation among different final state particles, sum of their energies, etc. Suitable cuts are applied on these observable with a goal to optimise the signal to background ratio. One should also worry about the instrumental backgrounds and design suitable selection cuts to eliminate them. The dominant SM backgrounds which can give the same type of event topology as the SUSY signals are

$$pp \to W/Z + jets\,, \tag{20}$$

$$pp \to t\bar{t} + jets\,, \tag{21}$$

$$pp \to QCD\,(jets) \tag{22}$$

with leptonic or hadronic decay of W/Z and top quark decay, $t \to bW \to bf\bar{f}'$(f=ℓ, q). Moreover, there are other sub-dominant SM process e.g. WW/WZ/ZZ + jets, etc. which can mimic the SUSY signals. For the clean detection of signal it is often needed to suppress the SM backgrounds by factor of $\sim 10^4$–10^5. We will see in a later section that using some experimental techniques and some kinematic selection cuts, the level of SM backgrounds can indeed be brought down to a negligible level.

The discovery potential of SUSY is usually studied in the literature on the basis of a variety of models predicting different type of final states (Sections 2 and 4). However for illustration we will discuss in this section SUSY signals at the LHC within the framework of the mSUGRA model. The SUSY signals in other scenarios will be taken up in a later section.

In the mSUGRA model, as discussed in section 2, there are only five free parameters specified at some unification scale and the relevant parameters at the energy of experimental interest can be obtained via the RG equations [10]. A set of these parameters, therefore completely determines the SUSY signals corresponding to a point in the mSUGRA parameter space. Several representative points, better known as the benchmark points are chosen such that the members of the set reflects different characteristics of SUSY signals. Both the CMS and the ATLAS collaborations have performed Monte Carlo studies on the feasibility of discovering SUSY with their detectors in this way. In the following we shall present results mainly based on the analysis by the CMS group [23].

We present in Fig. 3 the simulated $\not\!\!E_T$ distribution for SUSY events (the upper black curve) with only jets in the final state corresponding to one such bench mark point (LM1). This plot is made by full simulation[11] [23]. Along with the signal distribution, the distributions of the same observable for several SM backgrounds, Eqs. (20–22) are also shown. The signal is found to be well above the total background.

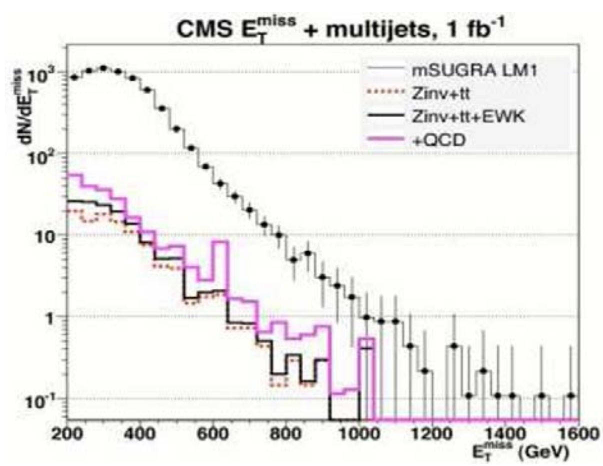

Figure 3. Missing E_T distribution for SUSY events and SM backgrounds with purely hadronic final states [23]

[11] In full simulation all possible detector effects are taken into account.

Recall that due to the long cascade decay chains in the signal events comparatively large multiplicities of jets and leptons appear. They are also harder as they originate from comparatively heavy sparticles like $\tilde{g}$, $\tilde{q}$. Therefore, a variable, called the effective mass, defined by the scalar sum of $\not{E}_T$ and the E_T of the four leading jets

$$M_{eff} = \sum_{j=1}^{4} |E_T^j| + |\not{E}_T| \tag{23}$$

will show a striking difference in its distribution in SUSY signal and SM background events. In Fig. 4 we present the distribution of M_{eff} obtained by the ATLAS collaboration [24] for illustration. In this figure

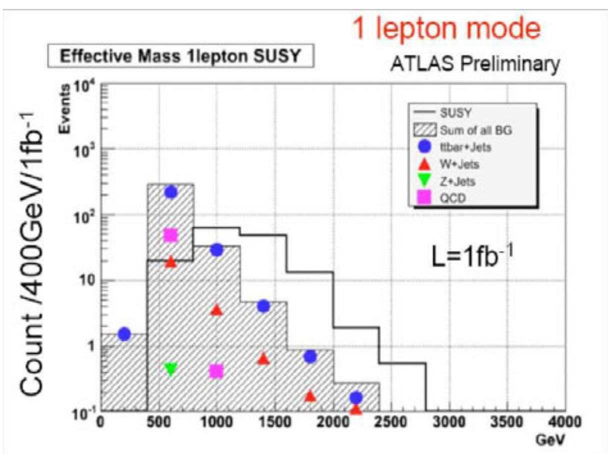

Figure 4. Effective mass M_{eff} distribution for SUSY events and SM backgrounds with 1 lepton + jets in the final states [24]

both the signal and the total background contributions are shown. Clearly the $\not{E}_T$ and as well as M_{eff} distributions will indicate unambiguously the existence of SUSY provided we have a complete understanding of the sources of fake $\not{E}_T$.

We present in Fig. 5, the CMS discovery reach of SUSY in the m_0–$m_{1/2}$ plane for different final state topologies corresponding to an integrated luminosity of 1 fb^{-1} [23] which will be accumulated during the first few months of data taking (footnote 10). The search channels with promising discovery potentials are now described briefly.

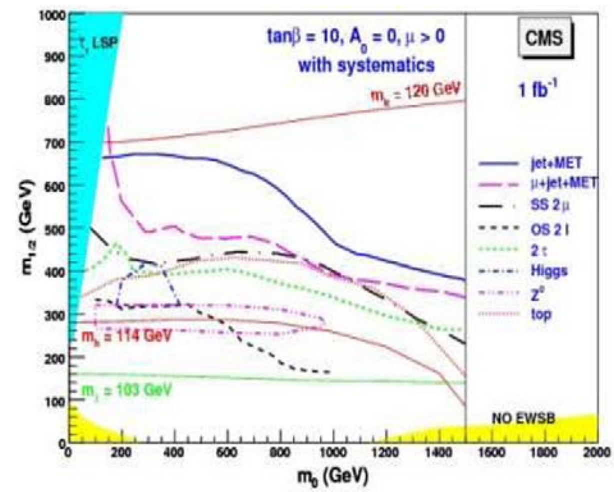

Figure 5. Discovery potential of the CMS experiment for 1 fb^{-1} luminosity [23]

• Jets + missing transverse energy ($\not{E}_T$).
This inclusive signal corresponding to the best discovery reach (see the solid line in Fig. 5) is due to gluino and squark production followed by their cascade decays to purely hadronic final states (Eqs. (15–18)). The dominant SM background is due to the Z + jets events with $Z \to \nu\nu$ giving missing energy like signal events. Other sources of SM backgrounds are $t\bar{t}$, WW, single top and pure QCD events. These backgrounds are suppressed by judiciously chosen kinematic cuts. From the squark and gluino mass formulae in terms of m_0 and $m_{1/2}$ (Section 2), it follows that for $m_0 \lesssim 1$ TeV, gluinos with mass $m_{\tilde{g}} \sim$ 1300–1400 GeV can be observed even at this low luminosity. However for $m_0 \gtrsim 1$ TeV, the squark mass becomes large yielding low rates leading a drop in the discovery reach.
• Single μ + Jets + $\not{E}_T$
In this signal final states containing at least one muon are considered. This muon may arise in the cascade decay chain due to the leptonic decay of a chargino Eq. (17). or a neutralino Eq. (18), where muons may also come from W or Z decays appearing in the decay cascade. Because of the relatively low BR of leptonic decays of $\tilde{\chi}_1^{\mp}$, which is about 11% for each species of letpons, the signal rates are depleted to some extent. However due to the presence of a high p_T muon the SM backgrounds, especially the background from pure QCD jets are small and consequently, the signal to background ratio large. In Fig. 5 the discovery reach

for this channel is shown by the long dashed line. At low values of m_0 squarks are light making $\tilde{\chi}_1^{\mp}\tilde{\chi}_1^{\pm}$ and $\tilde{\chi}_1^{\mp}\tilde{\chi}_2^0$ production rates large. They contribute dominantly to this final state along with $\tilde{g}$ and $\tilde{q}$ pair production making this channel the main discovery channel. On the whole this channel gives the best discovery reach after jet + $\not{E}_T$ channel.

In summary squark and gluino masses of about 1.5 TeV are within the striking range of the LHC for a nominal integrated luminosity of 1 fb^{-1}. This is already far above the mass reach of the Tevatron, the most powerful collider before the advent of the LHC. The above reach can be further extended to 2 TeV for an integrated luminosity of about 10 fb^{-1} [23]. With accumulating integrated luminosity larger regions of the SUSY parameter space will be accessible to the LHC experiments. For 300 fb^{-1} squark gluino masses upto 2.5 TeV can be probed but the precise mass reach is rather model dependent.

It is also seen from the Fig. 5 that several other search channels have reasonable reaches although they do not look like discovery channels. In fact large regions of the m_0–$m_{1/2}$ parameter space are covered by more than one search channel. The observation of the signal in multiple channels may provide additional information about the underlying theory. For example, the existence of a Z or the lightest Higgs boson in the final state would strongly suggest the presence of a $\tilde{\chi}_2^0$ in the decay cascade.

So far we have focused on generic SUSY signals which are events of the type n-lepton + m-jet + $\not{E}_T$. It is encouraging to note that the relative sizes of the signals corresponding to different m and n can distinguish between different regions of the mSUGRA parameter space. For example, the region corresponding to low m_0 allowed by dark matter data has three distinct regions. These regions can be distinguished by the above characteristics [25]. Since the ratio of the number of events in two different channels originating from the same production processes are practically independent of major theoretical uncertainties like the QCD scale dependence, this approach can indeed be helpful in obtaining additional information.

The particles belonging to the third-generation lead to special collider signatures. For example, the jets from the decay of B-hadrons (hadrons containing a b-quark) do not point to the collision vertex since the B, with a relatively large life time, travels a measurable distance away from this point before it decays. The jets from the hadronic decays of a τ contains very few hadrons compared to other jets. The jets coming from B or τ decays can be identified by the LHC experiments with large efficiency. Using the b-jet tagging or τ tagging facilities at the LHC, important additional information can be extracted from SUSY signals. For example $\tilde{\chi}_2^0$ can decay into a τ^+–τ^- pair with very large BR particularly for large tan β. In this case the τ slepton turns out to be much lighter than the other sleptons due to the RG effects and mixing as discussed in sections 2 and 3. Thus $\tilde{\chi}_2^0$ decays mediated by this sparticle dominantly goes to di-tau pairs. The members of the τ-pair are often polarised. This polarisation can be further exploited to improve the prospect of SUSY search at the LHC [26].[12]

CMS has also looked for SUSY signals in the di-tau channel and have found a reasonable reach. However even for moderately large tan β some of the low m_0 regions of the parameter space allowed by the dark matter data discussed above can lead to final states with many tagged τ and b-jets [25,27]. The importance of flavour tagging in identifying SUSY signals has been noted by several groups [28].

Another interesting region of the parameter space in the mSUGRA model is the 'focus point' region [29]. This is one of the few regions in the mSUGRA model consistent with the dark matter data (footnote 4). In this region characterised by large m_0, the squarks and the sleptons are too heavy to be detected at the LHC. However the gluino may well be within the reach of the LHC. Even in this region the top squark can be relatively light due to RG evolution (Section 2)[13]. As a result the gluino decays, though mediated by all squarks in principle, will be dominantly mediated by the lighter top squark. Thus they primarily decay into channels involving large number of t and b quarks. Some of these b quarks come directly from gluino decays and others from t decays. Consequently b-jet tagging can again be employed to distinguish the focus point signals from the ones from other mSUGRA regions as well as from the SM background [30].

3.3. Measurements of Masses at the LHC

So far we have discussed the prospects of discovering SUSY via different signals at the LHC. However the mere existence of these signals is not a conclusive evidence for SUSY. It is necessary to measure the mass, charge and other quantum numbers of each superparticle to identify the SUSY model chosen by nature. In this section we describe briefly the first few steps of achieving this goal. In our example, the second lightest

[12] See the article by Guchait and Roy in this volume for further details.

[13] Mixing effects (Section 3) are unlikely to be very significant here since $m_0 >> m_t$.

neutralino in a squark decay cascade decays to dileptons

$$\tilde{\chi}_2^0 \to \ell^+\ell^-\tilde{\chi}_1^0$$

which is a generic decay in many SUSY models (Fig. 6). This decay may occur via a two body decay mode if $m_{\tilde{\chi}_2^0} > m_{\tilde{\ell}}$ $(m_Z + m_{\tilde{\chi}_1^0})$, $\tilde{\chi}_2^0 \to \tilde{\ell}\ell \to \ell\ell\tilde{\chi}_1^0$ ($\tilde{\chi}_2^0 \to Z\tilde{\chi}_1^0 \to \ell\ell\tilde{\chi}_1^0$), otherwise it goes into three body modes mediated by an off-shell slepton or Z boson. In the

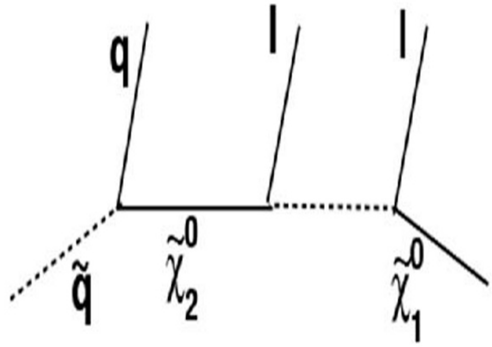

Figure 6. A simple decay chain containing a $\tilde{\chi}_2^0$ decay into the dilepton channel

case of three body decay, the distribution of dilepton invariant mass $m_{\ell\ell}$ shows a sharp edge and the position of the end point is exactly equal to $m_{\tilde{\chi}_2^0} - m_{\tilde{\chi}_1^0}$. At the end points the dilepton system and $\tilde{\chi}_1^0$ are at rest in the $\tilde{\chi}_2^0$ rest frame, so from the measurements of dilepton four vectors, it is possible to determine the four vector of $\tilde{\chi}_1^0$ and hence the four vectors of $\tilde{\chi}_2^0$. The mass of $\tilde{\chi}_1^0$ can also be determined from the kinematics. And eventually adding the four momentum of the correct jet in the event with that of $\tilde{\chi}_2^0$ it is possible to reconstruct the mass of the origin of the $\tilde{\chi}_2^0$ (either a $\tilde{g}$ or a $\tilde{q}$).

In case of two body decay of $\tilde{\chi}_2^0$, the edge of dilepton invariant mass distribution can be expressed in terms of $m_{\tilde{\ell}}$ and the neutralino masses

$$m_{\ell\ell}^{max} = m_{\tilde{\chi}_2^0}\sqrt{1 - \frac{m_{\tilde{\ell}}^2}{m_{\tilde{\chi}_2^0}^2}}\sqrt{1 - \frac{m_{\tilde{\chi}_1^0}^2}{m_{\tilde{\ell}}^2}}\,.$$

In case of $\tilde{q} \to q\tilde{\chi}_2^0$ decay, there are also similar end points for the $q\ell$ and qll. These end pints are expected to be measured within a good accuracy $\sim$ 10% [23]. Once these measurements are performed, it is possible to determine the masses of $m_{\tilde{\chi}_1^0}, m_{\tilde{\chi}_2^0}, m_{\tilde{\ell}}, m_{\tilde{q}}$. Recently, quite a few new ideas have been proposed to reconstruct the sparticle masses from different observables [31].

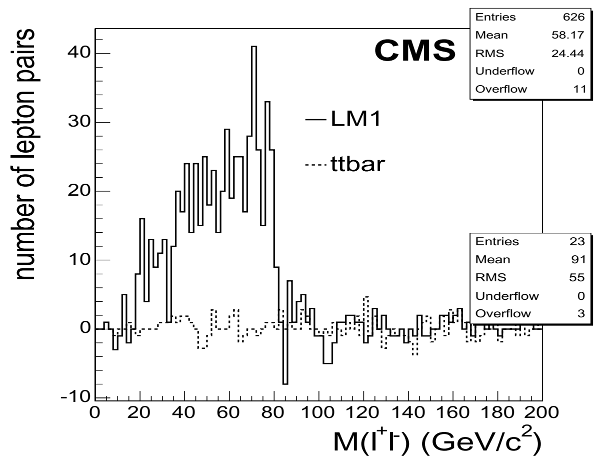

Figure 7. Invariant mass distribution of same flavour distributions [23]

Figure 7 presents the distribution of dilepton invariant mass of same flavour leptons after subtracting the background due to SUSY and SM. It is shown after a full simulation that end points can be measured within an uncertainty of $\sim$ 1 GeV [23].

4. Other Models of Mediation of SUSY Breaking

We have seen in the previous sections that the phenomenology of minimal supergravity has been very well studied with the lightest neutralino as the LSP over most of the parameter space. In this section we review some other mechanisms of SUSY breaking where the modes of communicating this breaking from the hidden sector to the observable sector MSSM are different. In particular, we discuss the LHC signatures of gauge-mediated SUSY breaking (GMSB) and anomaly-mediated SUSY breaking (AMSB) as well as signatures of models which are generalisation of mSUGRA, like models with non-universal scalar and gaugino masses and split-SUSY.

4.1. Gauge Mediated Supersymmetry Breaking

In gauge mediated SUSY breaking [32] the basic idea is to introduce new supermultiplets, not included in the

MSSM, as the messenger fields (Section 2).[14] They couple to the hidden sector which is the source of SUSY breaking. These messengers carry SM quantum numbers and, consequently couple directly to the SM gauge fields and the corresponding gauginos. Although the messenger fields do not couple to the sfermions in the lowest order, these couplings arise in higher orders of perturbation theory. The SUSY breaking in the hidden sector is thus communicated to the observable sector and soft masses and other SUSY breaking terms are generated.

One of the attractive features of gauge mediation is that the soft masses of squarks and sleptons depend only on their gauge quantum numbers. Thus for example, squarks of the same type (say, $\tilde{u}$ and $\tilde{c}$) are mass degenerate. In this way mass degeneracies required for suppressing FCNC effects in all sectors are guarantied.

For the sake of completeness we now review very briefly the theoretical ingredients of GMSB. The readers mostly interested in the LHC phenomenology may skip the next two paragraphs and directly move to the sparticle spectrum and collider signatures. In the simplest case, let us assume that the messenger fields are a set of left-handed chiral supermultiplets $\hat{\psi}$, $\hat{\bar{\psi}}$ and they have specific transformation properties under the SM gauge group. They are taken to be vectorlike with respect to the SM gauge interaction (i.e. $\hat{\bar{\psi}}\hat{\psi}$ are SM singlet). We assume that they couple to a gauge-singlet superfield S through a superpotential: $W_{\rm mess} = yS\hat{\psi}\hat{\bar{\psi}}$. Both scalar and auxiliary components of the superfield S acquires VEVs, denoted by $\langle S\rangle$ and $\langle F_S\rangle$ respectively. This way the fermionic and scalar component of the messengers get different masses: $m^2_{\rm fermions} = |y\langle S\rangle|^2$, $m^2_{\rm scalars} = |y\langle S\rangle|^2 \pm |y\langle F_S\rangle|$. In the messenger spectrum $\langle F_S\rangle \neq 0$ leads to SUSY violation and since the messenger fields are charged under the SM gauge groups, the gauginos of the MSSM can receive masses at one-loop. The scalars of the MSSM get leading contribution to their masses at the two-loop level.

If we assume that the messengers come in complete multiplets of the $SU(5)$ global symmetry, and are very close in mass, then approximate unification of gauge couplings will occur at the GUT scale M_U. In such a case if we consider the messengers to consist of N_5 copies of the $\mathbf{5} + \bar{\mathbf{5}}$ of $SU(5)$, then we have for the gaugino masses
$M_a = \frac{\alpha_a}{4\pi}\Lambda N_5 \;\; (a = 1, 2, 3)$.
Here we have introduced one more mass scale $\Lambda \equiv \langle F_S\rangle/\langle S\rangle$. Here $\alpha_a = g_a^2/4\pi$ and the g_a are the three gauge couplings of the SM. Theses gaugino masses are the running gaugino masses at an RG scale Q_0 corresponding to the average mass of the heavy messenger particles: $Q_0 \sim M_{\rm mess} \sim y\langle S\rangle$. For the MSSM scalars we have
$m^2_{\phi_i}(Q_0) = 2\Lambda^2 N_5 \sum_{a=1}^{3} \left(\frac{\alpha_a}{4\pi}\right)^2 C_a(i)$,
where $C_a(i)$ are the quadratic Casimir invariants of the representation of ϕ_i under the SM group. One should note here that the gaugino and scalar masses are in the same order in α. The requirement of gauge coupling unification demands that for messenger masses of order 10^6 GeV or less, one needs $N_5 \leq 4$. The trilinear soft terms are much smaller and one can assume them to be zero at the messenger scale Q_0. The LEP constraints on the sparticle masses and the requirement that the gluino mass is $\lesssim$ 1 TeV restrict Λ in the range 30 TeV $\lesssim \Lambda \lesssim$ 120 TeV.

In the GMSB spectrum the strongly interacting sparticles, the squarks and the gluino, are the heaviest since their soft masses are generated by the strong gauge interactions with the messenger sector. The soft masses of the sleptons and the EW gauginos are generated by weaker interactions. These sparticles are therefore, somewhat lighter. However the most distinctive feature of GMSB is that the gravitino ($\tilde{G}$) is the LSP. This can have very important consequences for collider physics [33] (see below). The mass of the gravitino is roughly in the range 1 eV $\lesssim m_{3/2} \lesssim$ 1 GeV. Since gravity is too weak, the direct production of gravitino in experiments is highly suppressed. However since this is the LSP, all the other SUSY particles will eventually decay into final states containing the gravitino leading to missing energy.

The decay width of a sparticle $\tilde{X}$ into its SM partner X and a gravitino (more precisely the longitudinal goldstino component) is

$$\Gamma(\tilde{X} \to X\tilde{G}) = \frac{m^5_{\tilde{X}}}{16\pi F^2}\left(1 - m^2_X/m^2_{\tilde{X}}\right)^4 . \qquad (24)$$

Here $\sqrt{F}$ is the fundamental scale of SUSY breaking and is typically ~ 100 TeV. The decay width (the decay length) of $\tilde{X}$ is larger (smaller) for smaller F.

The collider signals in GMSB depends sensitively on the nature of the next-to-lightest SUSY particle

[14] A supermultiplet consists of a particle and its superpartner (e.g. the u_L quark and the squark $\tilde{u}_L$ form a supermultiplet of MSSM). More formally a spin-1/2 particle (a quark or a lepton) and the corresponding sparticle can be looked upon as the components of a 'matter' superfield or a chiral superfield of L or R type. However a superfield also consists of unphysical particles called auxiliary components. Similarly a SM gauge boson, the corresponding gaugino and unphysical particles are components of a gauge superfield or a vector superfield. Usually the MSSM Lagrangian is written interms of these superfields, since a superfield as a whole has simple properties under SUSY transformations.

(NLSP). As we have discussed in an earlier paragraph, the answer to this question is model dependent. Nevertheless, there are only a few distinct possibilities. We have noted above that the gaugino masses scale like N_5, while the scalar masses like $\sqrt{N_5}$. As a result the sleptons will tend to be lighter than the EW gauginos for larger values of N_5. Thus for lower values of N_5, say, $N_5 = 1$, the lightest neutralino, which is usually bino like is the LSP over most of the parameter space. The NLSPs can also be the right (R)-type sleptons ($\tilde{e}_R$, $\tilde{\mu}_R$ and $\tilde{\tau}_R$). This possibility is favoured if N_5 is larger. For large value of $\tan\beta$ the stau ($\tilde{\tau}$) turns out to be lighter than the other two R-sleptons due to mixing and RG effects and becomes the sole NLSP in most cases. There is another possibility called the neutralino-stau co-NLSP scenario where signatures of neutralino NLSP and stau NLSP are both present.

The decay length of the (Eq. (24)) NLSP depends on the fundamental SUSY breaking scale and so a measurement of the decay length gives a measurement of the fundamental SUSY breaking scale $\sqrt{F}$. Depending on the decay length, the NLSP can decay inside the detector or outside. In the former case it may either decay promptly with a small life time or its life time may be large enough to produce an observable secondary vertex which can be reconstructed.

If the NLSP is the lightest neutralino then the dominant decay mode is $\tilde{\chi}_1^0 \to \gamma + \tilde{G}$. For a neutralino NLSP, its prompt decay into a photon and a gravitino (which escapes detection) can produce the inclusive $\gamma\gamma + \not{E}_T$ events at the LHC [34] and at the Tevatron [35]. For relatively light squarks and gluinos the production of these strongly interacting sparticles dominantly contribute to the signal. If on the other hand the squarks and gluinos are very heavy, the dominant contributions come from chargino and neutralino production. It has been shown that with an integrated luminosity of 10 fb^{-1}, the LHC experiments can probe values of $\Lambda \lesssim$ 400 TeV, corresponding to $m_{\tilde{g}} \leq 2.8$ TeV. Here $m_{\tilde{g}}$ is computed from its theoretical mass formula. This mass reach is certainly comparable to what we have in mSUGRA.

If the neutralino NLSP decays inside the detector with a long enough decay length then displaced secondary vertices appear. As a result the highly energetic photon from the neutralino NLSP decay is in general non-pointing (i.e. it may not be pointing to the interaction point). In this case the signal is again $\gamma\gamma + \not{E}_T$ events along with measurable decay lengths. This could be considered as the 'smoking gun' signature of GMSB.

The masses of the slepton ($\tilde{\ell}$) and the neutralino can be determined at the LHC from events with a lepton and a non-pointing photon arising from the cascade decay $\tilde{\ell} \to \ell\tilde{\chi}_1^0 \to \ell\gamma\tilde{G}$ [36]. The slepton $\tilde{\ell}$ may be copiously produced at the LHC from the gluino or squark cascades involving $\tilde{\chi}_2^0$ and $\tilde{\chi}_1^\pm$. A Monte Carlo simulation with the ATLAS detector shows that the masses could be measured with an error of 3% for $\mathcal{O}(100)$ $\ell\gamma$ pairs. The fitted masses are shown in Fig. 8. The fit results are $m_{\tilde{\ell}} = 162.1$ GeV and $m_{\tilde{\chi}_1^0} = 117.0$ GeV, while the input values are 161.7 GeV and 117.0 GeV, respectively. In Ref. [37] it has been shown that the

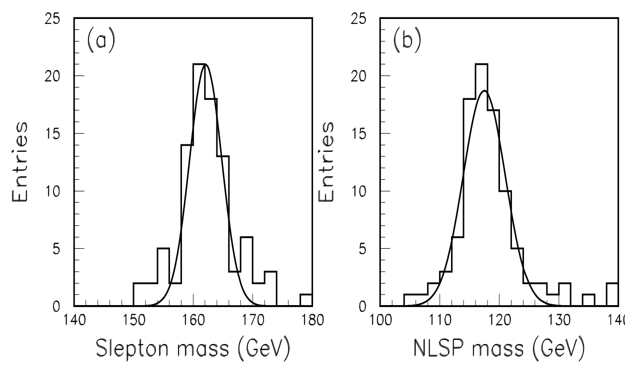

Figure 8. Distributions of the fit results of (a) the slepton mass $m_{\tilde{\ell}}$ and (b) the neutralino mass $m_{\tilde{\chi}_1^0}$. Results of Gaussian fitting are also shown. Taken from Ref. [36]

well studied $\gamma\gamma$ inclusive Higgs signal can be used at the LHC to test GMSB models in which a heavy neutral Higgs boson decays into two light neutralinos, the latter yielding two photons and missing energy.

If on the other hand a right handed slepton happens to be the NLSP it decays by $\tilde{\ell}_R \to \ell + \tilde{G}$. At the LHC pair production of NLSP sleptons with prompt decay to gravitino can give rise to final states $\ell^+\ell^- + \not{E}_T$. However, this signal suffers from large irreducible backgrounds. On the other hand, production of heavier states cascading to $\tilde{\ell}_R$ can give clean signatures with multiple leptons. For example, $\tilde{\ell}_L\tilde{\ell}_L$ pair production followed by the cascade decays $\tilde{\ell}_L^\pm \to \tilde{\ell}_R^+\ell^-\ell^\pm$ can produce final states $6\ell + \not{E}_T$ which do not suffer from SM backgrounds [38].

If the charged slepton decay takes place inside the detector with a measurable decay length, then one can observe very clear heavily ionising charged tracks with a kink. The "kinks" are sudden turns in the track of a charged particle. In this case the kinks appear where the charged lepton and the gravitino are being emmitted. As in the case of neutralino NLSP, a measurement of the decay length distribution can give a direct measure of the fundamental SUSY breaking scale. If the decay $\tilde{\ell} \to \ell + \tilde{G}$ takes place outside the detector, then the signature of GMSB would be two heavily ionizing charged tracks without missing energy. This is a very non-standard signature of SUSY and is easily detectable [39].

4.2. Anomaly-mediated Supersymmetry Breaking Models

Anomaly-mediated supersymmetry breaking (AMSB) involves a higher dimensional supergravity theory where the hidden sector and visible sector superfields are localised on two distinct parallel three-branes separated by a distance $\sim r_c$ (r_c is the compactification radius) in the extra dimension [40]. Assuming one single extra dimension, the flavour-violating terms are suppressed by factors like $e^{-r_c M_5}$, where M_5 is the 5-dimensional Planck scale. So as long as $r_c M_5 \gg 1$, the dangerous FCNC terms are absent.

The AMSB scenario can be described in terms of a 4-dimensional effective theory below the scale μ_c ($\sim r_c^{-1}$) where only 4-dimensional supergravity fields are assumed to propagate in the bulk. The whole process can be studied with the help of a "compensator" superfield Φ whose scalar and auxiliary components get VEVs according to $\langle\Phi\rangle = 1 + \langle F_\Phi\rangle\theta^2$. The auxiliary field acquires VEV ($\langle F_\Phi\rangle$) through its coupling to SUSY breaking sector. It also couples to the visible sector fields due to an anomalous violation of super-conformal invariance at the quantum level (super-Weyl anomaly). This causes SUSY breaking to show up in the visible sector and hence the name anomaly-mediated SUSY breaking. Denoting the VEV of the auxiliary component as F_Φ from now on, the resulting soft terms are

$$M_a = \frac{\beta_{g_a}}{g_a} F_\Phi; \quad a_y = \frac{\beta_y}{y} F_\Phi;$$
$$m_{\tilde{f}}^2 = -\frac{1}{4}\left(\frac{\partial\gamma}{\partial g}\beta_g + \frac{\partial\gamma}{\partial y}\beta_y\right)|F_\Phi|^2\,, \tag{25}$$

where appropriate beta functions ($\beta_{g,y}$) and anomalous dimensions (γ) are to be considered. It should also be noted that $F_\Phi \sim m_{3/2}$, the gravitino mass.

The analytic expressions for the scalar and gaugino masses are renormalisation group invariant, and thus, can be computed at the low-energy scale. However at low energies, it predicts the existence of tachyonic sleptons. Several solutions to this problem exist but here we will consider the minimal AMSB model wherein a constant term m_0^2 is added to all the scalar squared masses thereby making the slepton squared masses sufficiently positive. As a consequence, the RG invariance is lost and one needs to consider the expressions in Eq. (25) as the boundary conditions at the unification scale (M_U).

The minimal AMSB (mAMSB) model is described by just four parameters: $m_{3/2}$, m_0, $\tan\beta$ (the ratio of two Higgs VEVs) and sign(μ). The magnitude of μ is determined by requiring correct EW symmetry breaking. A particularly interesting feature of the mAMSB model is that the ratios of the respective $U(1)$, $SU(2)$ and $SU(3)$ gaugino mass parameters M_1, M_2 and M_3, at low energies turn out to be $|M_1| : |M_2| : |M_3| :: 2.8 : 1 : 7.1$. An immediate consequence is that the lighter chargino and the lightest neutralino are both almost exclusively a Wino and, hence, nearly degenerate in mass (each with mass approximately equal to M_2). A very small mass splitting ($\Delta M_{\tilde{\chi}_1} \sim 200$ MeV) is generated at the tree level as well as from the one-loop corrections. The gravitino mass is much larger than the MSSM soft terms. Left and right charged sleptons of first two-generations are nearly mass degenerate. The staus are somewhat split in mass with almost maximal mixing. In this model, the squarks are always much heavier than the sleptons due to larger contributions from the strong coupling constant.

The lighter chargino is long-lived with the dominant decay mode $\tilde{\chi}_1^\pm \to \tilde{\chi}_1^0 + \pi^\pm$. This would typically result in a short heavily ionising charged track and/or a characteristic soft pion in the detector [41]. This is the 'smoking gun' signature of anomaly mediation with a Wino LSP, and is applicable beyond the minimal model. These charginos can be triggered on when produced in association with jets. It has been shown that for Tevatron Run II with 2 fb^{-1} luminosity, one could discover chargino masses up to 300 GeV [41]. In order to distinguish this scenario from GMSB with long-lived sleptons, correlations between particle masses and cross sections may be used. A similar analysis has been performed also at the LHC for the high luminosity run [42].

It has been shown that with an integrated luminosity of 10 fb^{-1}, the dilepton plus jets plus $\not\!\!E_T$ channel offers the reach at the LHC to values of the gluino mass ($m_{\tilde{g}}$) ~ 2.5 TeV for low values of m_0. For large m_0, the reach is $m_{\tilde{g}} \sim 1.3$ TeV where the best signature is typically 0 or 1 isolated lepton plus jets plus $\not\!\!E_T$ [43]. The presence of terminating tracks due to the long-lived chargino in the signal events could help to distinguish the mAMSB model from other SUSY breaking mediation model. The discovery potential of the LHC has also been investigated for the mAMSB model, using the ATLAS fast detector simulator, including track reconstruction and particle identification. It has been found that with 100 fb^{-1} of integrated luminosity the search will reach upto 2.8 TeV in the squark mass and 2.1 TeV in the gluino mass for a 5σ discovery [44]. In Fig. 9, the number of $\tilde{\chi}_1^+$s which would be produced at the LHC and decay within a fiducial volume in the active material of the ATLAS tracker is plotted as a function of m_0 and $m_{3/2}$.

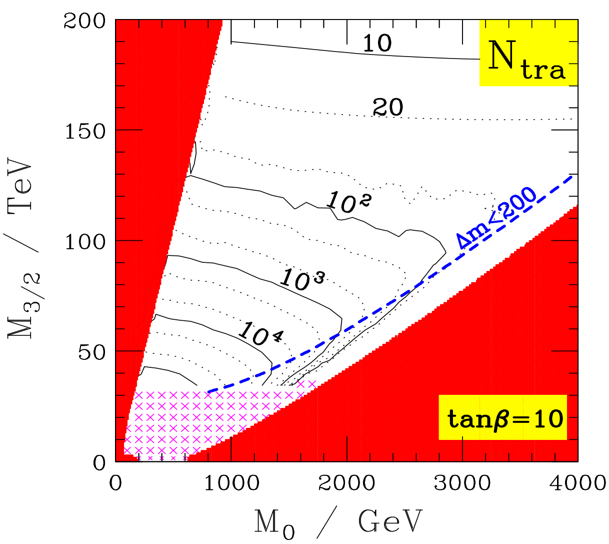

Figure 9. The number of $\tilde{\chi}_1^{\pm} \rightarrow \tilde{\chi}_1^0$ decays expected within the central region of the ATLAS detector ($|\eta| < 2$) with transverse decay vertices between 100 mm and 800 mm from the interaction point for integrated luminosity of 100 fb^{-1}. The initial track from the chargino is required have $p_T > 10$ GeV. $\Delta M_{\tilde{\chi}_1}$ is < 200 MeV. Figure is taken from Ref. [44]

4.3. Non-universal Gaugino and Scalar Masses

Most studies on the signatures of SUSY at the LHC, are based on the assumption that the gaugino masses are universal at the unification scale as in the mSUGRA model. There is, however no compelling theoretical reason for such a choice. In a supersymmetric grand unified theory (SUSYGUT) like the SU(5) grand unified model, gaugino masses are generated by the VEV of the auxiliary scalar component F_Φ of a superfield Φ in the gauge kinetic function $f(\Phi)$ (see below). When the scalar field is chosen to be a singlet under the GUT group the gaugino masses turn out to be universal usualy parametrised by a single parameter $m_{1/2}$. If on the other hand a GUT non-singlet scalar is chosen non-universal gaugino masses emerge but these masses are related by calculable group theoretic factors (Table 3).

For the more theoretically oriented readers we summarise in the next few lines the basic mechanism of generating non-universal gaugino masses. The readers only interested in the LHC signatures can directly go to Table 3 and the discussions following Eq. (28).

The function $f(\Phi)$ is an analytic function of the chiral superfields Φ in the theory [45]. It should be noted that the chiral superfields Φ consist of a set of gauge singlet superfields Φ^s and gauge non-singlet superfields Φ^n, respectively, under the grand unified group. If the auxiliary part F_Φ of a chiral superfield Φ in the $f(\Phi)$ gets a VEV, then gaugino masses arise from the coupling of $f(\Phi)$ with the field strength superfield W^a. The Lagrangian for the coupling of gauge kinetic function with the gauge field strength is written as

$$\mathcal{L}_{gk} = \int d^2\theta f_{ab}(\Phi) W^a W^b + H.c. \,, \tag{26}$$

where a and b are gauge group indices [for example, $a, b = 1, 2, ..., 24$ for SU(5)] and repeated indices are summed over. The gauge kinetic function $f_{ab}(\Phi)$ is

$$f_{ab}(\Phi) = f_0(\Phi^s)\delta_{ab} + \sum_n f_n(\Phi^s)\frac{\Phi^n_{ab}}{M_P} + \cdots \,, \tag{27}$$

where as described above the Φ^s and Φ^n are the singlet and non-singlet chiral superfields, respectively. Here $f_0(\Phi^s)$ and $f_n(\Phi^s)$ are functions of gauge singlet superfields Φ^s and M_P is some large scale. When F_Φ gets a VEV $\langle F_\Phi \rangle$, the interaction (26) gives rise to gaugino masses

$$\mathcal{L}_{gk} \supset \frac{\langle F_\Phi \rangle_{ab}}{M_P} \lambda^a \lambda^b + H.c. \,, \tag{28}$$

where $\lambda^{a,b}$ are gaugino fields.

We now illustrate the generation of non-universal gaugino masses in the framework of the group SU(5). Since the gauginos belong to the **24**-dimensional adjoint representation of the gauge group, F_Φ can belong to any of the following representations appearing in the symmetric product of the two **24**-dimensional representations of SU(5):

$$(\mathbf{24} \otimes \mathbf{24})_{Symm} = \mathbf{1} \oplus \mathbf{24} \oplus \mathbf{75} \oplus \mathbf{200}. \tag{29}$$

In the minimal case, F_Φ is assumed to be in the singlet representation of SU(5). This corresponds to equal gaugino masses at the GUT scale. However, Φ can belong to any of the non-singlet representations **24**, **75**, and **200** of SU(5). In that case, the gaugino masses are unequal but related to one another via the representation invariants [46]. It should be kept in mind that an

F_Φ	M_1^G	M_2^G	M_3^G	$M_1^{\rm EW}$	$M_2^{\rm EW}$	$M_3^{\rm EW}$
1	1	1	1	0.14	0.29	1
24	–0.5	–1.5	1	–0.07	–0.43	1
75	–5	3	1	–0.72	0.87	1
200	10	2	1	1.44	0.58	1

Table 3. *Ratios of gaugino mass parameters at the GUT scale in the normalisation $M_3^{\rm G}$ = 1 and at the electroweak scale in the normalisation $M_3^{\rm EW}$ = 1 at the one-loop level.*

arbitrary combination of these different representations is also allowed.

In Table 3, we display the ratios of the resulting gaugino masses at tree level for F_Φ belonging to different representations of SU(5). Clearly, the non-singlet representations have characteristic mass relationships for the gauginos at the GUT scale. The resulting relations at the EW scale, using the renormalisation group (RG) evolution (Section 2) at the one-loop level are also displayed.

The phenomenology of supersymmetric models with non-universal gaugino masses has been considered e.g. in [47–52] and very recently in [53–56] in the context of the LHC. The phenomenology of supersymmetric models depends crucially on the compositions of neutralinos and charginos. Hence, it is extremely important to investigate the changes in the experimental signatures with the changes in the composition of neutralinos and charginos which arise because of the non-universal gaugino masses at the GUT scale.

In Ref. [56] a multichannel analysis of SUSY signals has been carried out at the LHC for a number of non-universal representations breaking the SU(5) and SO(10) GUT groups. These channels include *jets* + $\not{E}_T$, *same-sign-dileptons* (SSD), *opposite-sign-dileptons* (OSD), *trileptons + jets +* $\not{E}_T$ and *single lepton + jets +* $\not{E}_T$. The results have been compared with those coming from the universal gaugino masses. It has been found that the most useful way to discriminate among the various cases is to look for the ratios of events rates in various channels. Hadronically quiet trilepton channels were studied in [54].

If the squarks and gluinos are light enough, their production cross sections are large at the LHC. The light neutralinos $\tilde\chi_{1,2}$ are typical decay products of $\tilde{\rm g}$ and $\tilde{\rm q}$. Following [17] it was already noted in section 3.2 that the neutral Higgs bosons can be copiously produced in the decay of $\tilde\chi_2$, if the mass difference between $\tilde\chi_2$ and $\tilde\chi_1$ is large enough. At the LHC heavy higgs (H/A) search may be difficult for low and medium $\tan\beta$ values [17,57]. The cross section for H/A production decreases rapidly for masses around ~ 280 GeV/c^2. Thus the discovery region is not expected to cover these high Higgs mass values even with high luminosity unless the cross section is boosted by high $\tan\beta$. As the squark, gluino production rates are largely independent of the value of $\tan\beta$, Higgs production via $\tilde\chi_2$ decays have been found to be particularly interesting. It is, therefore, worthwhile to revisit the prospect of Higgs search in these channels in models with non-universal gaugino masses.

Production of Higgs bosons via $\tilde\chi_2 \to \tilde\chi_1$ h/H/A in models where gaugino mass patterns are as in the singlet and **24** representation (Table 3) was studied in [47]. It was found for sample parameters that only the lightest Higgs could be produced in the model with singlet, while also the heavy Higgses, H and A, could be produced from the cascade in the model with **24** representation. Recently the role of gaugino mass non-universality (in a model independent framework) in the context of Higgs productions has been discussed in [55].

In [53] the discovery potential of the neutral Higgs bosons in the representation 24 was studied and the results were compared with those obtained with universal gaugino masses. The event selection is based on the requirement of four energetic jets, large missing $\rm E_T$, separation of the jets into two hemispheres and the reconstruction of the Higgs boson mass from two jets tagged as b jets. The event generation and simulation were performed and the results were made public in the CMS framework [58]. The discovery potential in the (m_A–$\tan\beta$) plane is shown in Fig. 10, where m_A is the mass of the CP-odd Higgs boson. The discovery region extends to H/A masses of $\sim$210 GeV/c^2 around $\tan\beta = 10$ and to $\sim$190 GeV/c^2 around $\tan\beta = 30$. For $\sim$ 190 GeV/c^2 the low $\tan\beta$ region was covered with $> 5\sigma$ significance down to $\tan\beta = 2$.

There is also scope for non-universality in scalar masses which can have interesting implications for the signatures at the LHC. It is known for quite some time that non-universal soft SUSY breaking terms for the scalars at a very high scale, may arise naturally within the supergravity framework [59–61]. For example, non-universal soft SUSY breaking terms may arise at the GUT scale due to renormalisation group evolutions from the Planck scale to the GUT scale [60]. On the other hand, if the rank of a GUT group is reduced by spontaneous symmetry breaking, one may obtain D-term contributions to scalar masses [61]. The size of these new contributions to scalar masses can be comparable to the universal mass m_0 in mSUGRA. These

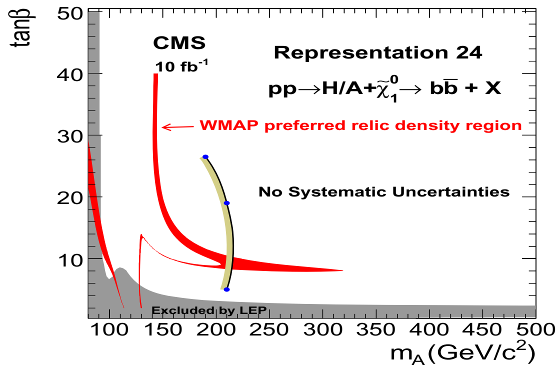

Figure 10. CMS discovery potential at 5σ level for $\tilde{q}, \tilde{g} \to \tilde{\chi}_1 H/A \to \tilde{\chi}_1 b\bar{b}$ in the representation **24** with an integrated luminosity of 10 fb^{-1}. No systematic uncertainties are included

contributions will in general have different values for different members of the same GUT multiplet which can lead to non-universal squark and slepton masses at the GUT scale. These non-universal terms are generation independent and do not lead to additional problems related to FCNCs. The low energy spectrum in these cases evolves from these high scale inputs. Because of these non-universalities of the scalar masses at a very high scale, the low energy sparticle spectrum can be very different from the universal scalar mass (m_0) case and the signatures at the LHC can also become significantly distinct. Non-universalities in the scalar masses may affect the SUSY signals through the gluino branching ratios as well as through the total squark-gluino production cross section [62]. Recently, the signatures at the LHC for the cases with high-scale scalar non-universality have been studied in various channels such as like and opposite sign dileptons, inclusive and hadronically quiet trileptons, inclusive jets, etc. [54,63]. It has been observed that these non-universalities affect the ratios of various final states produced through the decay cascades of sparticles. A more detailed discussion of these issues is beyond the scope of this article.

4.4. Split Supersymmetry

Let us discuss a more recent and very interesting scenario, namely split SUSY which assumes a very large splitting between the soft scalar and fermion masses in the MSSM. In order to understand this scenario better, let us first recall that the naturalness criterion has been one of the guiding principles in the formulation of the MSSM. This criterion demands that the masses of the superpartners should be somewhere around 1 TeV or so. However, the naturalness criterion is incompatible with the tiny cosmological constant. An important question, therefore arises. Is it possible to abandon the principle of naturalness and maintain the nice phenomenological aspects of the MSSM at the same time?

It has been argued [64] that the successful unification of gauge couplings of the MSSM can be retained even when all the scalars of the theory, except one fine-tuned light Higgs boson lie far above the EW scale. In this scenario, despite the loss of the original motivation to cure the hierarchy problem, one can still find a supersymmetric theory which is free of many of the undesirable features of the MSSM such as the flavour and CP problem, fast proton decay via dimension five operators, a tightly constrained mass of the lightest Higgs, etc. The gauginos and Higgsinos of this theory are chosen to lie near the TeV scale to ensure gauge coupling unification at $M_U \sim 10^{16}$ GeV as well as a stable LSP in the desirable mass range. These features describe the scenario of split SUSY.

Split SUSY is an effective theory in which the heavy scalars are integrated out and assumed to have degenerate mass $\tilde{m}$. The coupling constants at the scale $\tilde{m}$ are obtained by matching the Lagrangian describing the effective theory with the interaction terms of the SUSY Higgs doublets H_u and H_d. The Higgs doublet $h = -cos\beta\epsilon H_d^* + sin\beta H_u$ is fine-tuned to have a small mass term. One can identify a minimal split SUSY model described by four parameters: (1) the common mass $\tilde{m}$, (2) $\tan\beta$, (3) the Higgsino mass parameters $\mu(M_U)$ at the GUT scale and (4) the gluino mass $m_{\tilde{g}}$. It has been shown that certain special constraints [65] are imposed on the parameter space of the minimal split SUSY model by the infrared fixed point of the top Yukawa coupling. In order to study the RG evolution of various masses and couplings from the unification scale down to the EW scale, one should remember that between the unified scale M_U and the scale of heavy scalars $\tilde{m}$, the theory is described by the MSSM fields and we should use the MSSM RG equations in that region. Below the scale $\tilde{m}$, one should use the spectrum of split SUSY with gauginos and higgsinos included in the two-loop evolution.

Phenomenologically the most interesting feature of this scenario is that the gluinos are long-lived due to large squark masses which mediate their decays. Negative searches for anomalously heavy isotopes suggest that $\tilde{m} \lesssim 10^{13}$ GeV for 1 TeV gluino. In colliders these long-lived gluinos can produced displaced vertices. These gluinos can also hadronise into colour singlet states (called R-hadrons). If these states are neutral then they lose energy through hadronic interactions and if they are charged then they can deposit energy in the form of ionisation [66]. These types of signatures of the long-lived gluions are the 'smoking gun' signatures of split-SUSY. Collider phenomenology of split-SUSY models were also considered in [67]. It was found that the long-lived gluino can be discovered at the LHC even if their masses are above 2 TeV.

5. Conclusion

In this work we have reviewed how various SUSY breaking mechanisms lead to different soft masses of the sparticles at very high energies (Sections 2, 4.1–4.4). These masses in turn lead to model dependant sparticle spectrum at the energy scale of experimental interest (the EW scale). The nature of the lightest supersymmetric particle which carries the missing energy as well as the decay cascades of the heavier sparticles are also different in various SUSY breaking scenarios of contemporay interest. This has very interesting implications for cosmology and collider physics.

The expected SUSY signatures at the LHC have been studied by Monte carlo simulations. The generic $\mathrm{n-leptons} + \mathrm{m-jets} + \not{E}_T$; n, $m = 0, 1, 2, 3..$ signatures are primarily due to squark-gluino production with large cross sections followed by cascade decays of these sparticle. Even if the squarks and gluinos have masses as large as 2.5 TeV, they are likely to be within the striking range of the LHC in most models and in most cases the jets + $\not{E}_T$ signal appears to be the most potent search channel. However, a characteristic signature of GMSB (Section 4.1), which distinguishes it from other SUSY breaking mechanisms, is the additional abundance of photons in the final states - a smoking gun signal.

The other search channels with different values of m and n may provide complementary information. The flavour tagging facilities at the LHC like b-jet tagging and τ-jet tagging may be helpful in identifying specific regions of the parameter space (Section 3.2 and references there in).

In order to establish SUSY beyond doubt it is essential to reconstruct the masses and other quantum numbers of the sparticles - a daunting task indeed. Mass reconstruction is possible in some favourable scenarios (Section 3.3 and 4.1 and references there in).

Several mechanisms of SUSY breaking predicts long lived sparticles providing 'smoking gun' signals of the underlying models. Examples are the next-to-lightest-supersymmetric particle of the GMSB model (either the lightest neutralino or a τ-slepton), the lighter chargino of the AMSB model and the gluino in split SUSY. The charged tracks or displaced vertices associated with these long lived spartcles will be tantalising to observe.

Physics at the LHC promises to be stimulating for many years to come!

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REFERENCES

1. P. Ramond, Phys. Rev. D3 (1971) 2415; A. Neveu and J.H. Schwarz, Nucl. Phys. B31 (1971) 86; J.L. Gervais and B. Sakita, Nucl. Phys. B34 (1971) 632; Yu. A. Gol'fand and E.P. Likhtman, JETP Lett.13 (1971) 323; D.V. Volkov and V.P. Akulov, Phys. Lett. B46 (1973)109; J. Wess and B. Zumino, Nucl. Phys. B70 (1974) 39
2. For textbooks see, e.g., *Perspectives in Supersymmetry,* ed. G.L.Kane, World Scientific(1998); *Theory and Phenomenology of sparticles*, M. Drees, R.M. Godbole and P.Roy; World Scientific(2004); *Weak scale Supersymmetry: From superfields to scattering events,*, H. Baer and X. Tata, Cambridge UK, Univ. Press(2006). For reviews, see, e.g., H. P. Nilles, Phys. Rept. **110** (1984) 1; H. E. Haber and G. Kane, Phys. Rept. **117** (1985) 75; S.P. Martin, hep-ph/9709356 (published in the first reference, ed. G. L. Kane)
3. For recent reviews see, e.g., A. Derevianko, S. G. Porsev, Eur.Phys.J.A32 (2007); J. Guena, M. Lintz and M. Bouchiat, Mod.Phys.Lett.A20 (2005) 375
4. for details see Tevatron Electroweak Working Group http:tevewwg.fnal.gov/; see also Precision electroweak physics at the Tevatron, by CDF and D0 Collaborations (Eric B. James for the collaboration), hep-ex/0701003
5. S. Weinberg, Phys. Rev. D13 (1976) 974; Phys. Rev. D19 (1979) 1277; E. Gildener, Pjhys. Rev. D14 (1976) 1667; L. Susskind, Phys. Rev. D20 (1979) 2619; G. 't Hooft, in *Recent developments in gauge theories*, Proceedings of the NATO Advanced Summer Institute, Cargese 1979, (Plenum, 1980)
6. S. Dimopoulos and S. Raby, Nucl. Phys. B192 (1981) 353; E. Witten, Nucl. Phys. B188 (1981) 513; M. Dine, W. Fischler, and M. Srednicki, Nucl. Phys. B189 (1981) 575; S. Dimopoulos and H. Georgi, Nucl. Phys. B193 (1981) 150; N. Sakai, Z. Phys. C11 (1981) 153; R. K. Kaul, Phys.Lett.B109 (1982) 19; R.K. Kaul and P. Majumdar, Nucl. Phys. B199 (1982) 36
7. See any modern text book on quntum field theory like *An Introduction to Quantum Field Theory*, M. E. Peskin and D. V. Schroeder, Westview Press(1995). First Indian edition Levant Books, Kolkata (2005)
8. Some standard text books are *Grand Unified Theo-*

ries, G. G. Ross, Benjamin-Cummins (1984); *Unification and Supersymmetry*, R. N. Mohapatra, Springer Verlag (1986)
9. A.H. Chamseddine, R. Arnowitt, and P. Nath, Phys. Rev. Lett. 49 (1982) 970; R. Barbieri, S. Ferrara, and C.A. Savoy, Phys. Lett. B119 (1982) 343; L. J. Hall, J. Lykken and S. Weinberg, Phys. Rev. **D27** (1982) 2359; N. Ohta, Prog. Theor. Phys. **70**, 542 (1983).
10. K. Inoue, A. Kakuto, H. Komatsu and S. Takeshita, Prog. Theor. Phys. **67** (1982) 1889; ibid. 68 (1982) 927; Erratum-ibid. **70** (1983) 330
11. L. E. Ibanez and G. G. Ross, Phys. Lett. B **110**, (1982) 215; K. Inoue, A. Kakuto, H. Komatsu and S. Takeshita, Prog. Theor. Phys. **68**, (1982) 927; Erratum-ibid. **70**, (1983) 330; J. R. Ellis, J. S. Hagelin, D. V. Nanopoulos and K. Tamvakis, Supergravity," Phys. Lett. B **125**, (1983) 275; L. Alvarez-Gaume, J. Polchinski and M. B. Wise, Nucl. Phys. **B221**, (1983) 495
12. For the latest limits on the sparticle masses from LEP experiments see: http://lepsusy.web.cern.ch/lepsusy/
13. For a recent review with almost upto date experimental limits on sparticle masses see, L. Pape and D. Treille, Rep. Prog. Phys. **69** 2843
14. For the latest results from CDF see http://www-cdf.fnal.gov/physics/exotic/exotic.html
15. For D0 results on new phenomena see http://www-d0.fnal.gov/Run2Physics/WWW/results/np.html
16. G.L. Kane and J.P. Leveille, Phys. Lett. B112 (1982) 227; P.R. Harrison and C.H.Llewellyn-Smith, Nucl. Phys. B213 (1983) 223 (Erratum: Nucl. Phys. B223 (1983) 542); S. Dawson, E. Eichten and C. Quigg, Phys. Rev. D31, (1985) 1581; E. Reya and D.P. Roy, Phys. Lett. **B141** (1984)442; Phys. Rev. D32 (1985) 645
17. Aseshkrishna Datta, A. Djouadi, M. Guchait, F. Moortgat, Nucl.Phys. B681 (2004) 31
18. W. Beenaker, R. Hopker, M. Spira and P. M. Zerwas, Nucl.Phys. B492 (1917) 51
19. For a review see, e.g., T. Sjostrand, hep-ph/0611247
20. G. Corcella *et al*, J. High. Energy Phys. **JHEP01** (2001) 010
21. F. E. Paige, S. D. Protopopescu, H. Baer and X. Tata, hep-ph/0312045
22. T. Sostrand *et al*, Comput. Phys. Commoun. , **135** (2001) 238
23. The CMS collaboration, CERN-LHCC-2006-021(2006); J. Phys G: Nucl. Part. Phys. 34 (2007) 995
24. ATLAS Collaboration, Expected Performance of the ATLAS Experiment, Detector, Trigger and Physics, CERN-OPEN-2008-020, Geneva, 2008, to appear
25. N. Bhattacharyya, Amitava Datta and S. Poddar, arXiv:0807.0278 (to appear in Phys. Rev. D)
26. M.Guchait and D.P.Roy, Phys. Lett. B541(2002)356
27. U. Chattopadhyay, D. Das, Amitava Datta, S. Poddar, Phys.Rev.D76 (2007) 055008
28. R.H.K. Kadala, P.G. Mercadante, J.K. Mizukoshi, Xerxes Tata, Eur.Phys.J. **C56** (2008) 511; Daniel Feldman, Zuowei Liu, Pran Nath, Phys.Rev.D78 (2008) 083523
29. K.L.Chan, U. Chattopadhyay and P.Nath, Phys. Rev. D58 (1998) 096004; J. Feng, K. Matchev and T. Moroi, Phys. Rev. Lett. 84 (2000) 2322, Phys. Rev D61 (2000) 075005
30. U. Chattopadhyay, Amitava Datta, Aseshkrishna Datta, Anindya Datta and D.P.Roy, Phys. Lett. B493 (2000) 127; P.G.Mercadante, J.K.Mizukoshi and X. Tata, Phys. Rev D72 (2005) 035009; H. Baer, V. Barger, G. Shaughnessy, H. Summy and L.-t. Wang, Phys.Rev. **D75** (2007) 095010; S.P.Das, Amitava Datta, M.Guchait, M. Maity, S. Mukherjee, Eur. Phys. J C54 (2008) 645
31. B.C.Allanach, C.G.Lester, M. A. Parker and B.R.Webber, JHEP 0009 (2000) 004; M. M. Nojiri,Y. Shimizu, S. Okada and K. Kawagoe, JHEP 806 (2003) 035. W.S. Cho, K. Choi, Y.G.Kim, C. B. Park, JHEP 0802 (2008) 035; G.G.Ross, M. Serna, Phys. Lette B665 (2008) 212; D. J. Miller, P. Osland and A.R.Raklev, JHEP 03 (2006) 034; M. M. Nojiri, G. Polesello and D.R.Tovey, JHEP 0805 (2008) 14
32. M. Dine and W. Fischler, Phys. Lett. B**110**, (1982) 227; C.R. Nappi and B.A. Ovrut, Phys. Lett. B**113**, (1982) 175; L. Alvarez-Gaumé, M. Claudson and M.B. Wise, Nucl. Phys. B **207**, (1982) 96; For reviews, see, e.g. C. Kolda, ArXiv: hep-ph/9707450; G.F. Giudice and R. Rattazzi, Phys. Rept. **322**, (1999) 419
33. See, e.g., S. Dimopoulos, S. Thomas, and J.D. Wells, Phys. Rev. D **54**, (1996) 3283; Nucl. Phys. B**488**, (1997) 39; B. Mukhopadhyaya and S. Roy, Phys. Rev. D **57**, (1998) 6793; J.A. Bagger *et al.*, Phys. Rev. D**55**, (1997) 3188
34. H. Baer *et al.*, Phys. Lett. B **435**, (1998) 109
35. S. Ambrosanio *et al.*, Phys. Rev. D **54**, (1996) 5395
36. K. Kawagoe *et al.*, Phys. Rev. D**69**, (2004) 035003
37. J.L. Diaz-Cruz, D.K. Ghosh, and S. Moretti, Phys. Rev. D**68**, (2003) 014019
38. S. Dimopoulos, M. Dine, S. Raby, S. Thomas, J.D. Wells, Nucl. Phys. B (Proc. Suppl.) **52A**, (1997) 38
39. See the second reference in [34]
40. L. Randall and R. Sundrum, Nucl. Phys. B**557**, (1999) 79; G.F. Giudice *et al.*, J. High Energy Phys. **12**, (1998) 027
41. J.L. Feng *et al.*, Phys. Rev. Lett. **83**, (1999) 1731; S. Roy, Mod. Phys. Lett. **A19**, (2004) 83 and references therein. T. Gherghetta, G.F. Giudice, and J.D. Wells, Nucl. Phys. B **559**, (1999) 27
42. M. Ibe, T. Moroi, and T.T. Yanagida, Phys. Lett. B**644**, (2007) 355; S. Asai *et al.*, Phys. Lett. B**653**, (2007) 81
43. H. Baer, J.K. Mizukoshi, and X. Tata, Phys. Lett. B**488**, (2000) 367
44. A.J. Barr, B.C. Allanach, C.G. Lester, M.A. Parker, and P. Richardson, J. High Energy Phys. **03**, (2003) 045
45. E. Cremmer, S. Ferrara, L. Girardello, and A. Van Proeyen, Phys. Lett. B**116**, (1982) 231

46. J. Ellis, K. Enqvist, D.V. Nanopoulos, and K. Tamvakis, Phys. Lett. B**155**, (1985) 381; M. Drees, Phys. Lett. B**158**, (1985) 409; G. Anderson *et al* hep-ph/9609457
47. K. Huitu, J. Laamanen, P.N. Pandita, and S. Roy, Phys. Rev. D**72**, (2005) 055013
48. G. Anderson, H. Baer, C.H. Chen, and X. Tata, Phys. Rev. D**61**, (2000) 095005
49. K. Huitu, Y. Kawamura, T. Kobayashi, and K. Puolamaki, Phys. Rev. D**61**, (2000) 035001
50. G. Belanger, F. Boudjema, A. Cottrant, A. Pukhov, and A. Semenov, Nucl. Phys. **B706**, (2005) 411
51. A. Djouadi, Y. Mambrini, and M. Muhlleitner, Eur. Phys. J. **C20**, (2001) 563
52. S.I. Bityukov and N.V. Krasnikov, Phys. Atom. Nucl. **65**, (2002) 1341; Yad. Fiz. **65**, (2002) 1374
53. K. Huitu, R. Kinnunen, J. Laamanen, S. Lehti, S. Roy, and T. Salminen, arXiv:0808.3094 [hep-ph]
54. S. Bhattacharya, Aseshkrishna Datta, and B. Mukhopadhyaya, arXiv:0809.2012 [hep-ph]
55. P. Bandyopadhyay, Aseshkrishna Datta, and B. Mukhopadhyaya, arXiv:0806.2367 [hep-ph]
56. S. Bhattacharya, Asehkrishna Datta, and B. Mukhopadhyaya, J. High Energy Phys. **10**, (2007) 080
57. H. Baer, M. Bisset, X. Tata, and J. Woodside, Phys. Rev. D**46**, (1992) 303
58. CMS PRS, http://cmsdoc.cern.ch/cms/PRS/results/susybsm/rep24/rep24.html
59. See, e.g., M. Olechowski and S. Pokorski, Phys. Lett. B**344**, (1995) 201; T. Kobayashi, D. Suematsu, K. Yamada, and Y. Yamagishi, Phys. Lett. B**348**, (1995) 402
60. P. Moxhay and K. Yamamoto, Nucl. Phys. **B256**,(1985) 130; B. Gato, Nucl. Phys. **B278**, (1986) 189; N. Polonsky and A. Pomarol, Phys. Rev. D**51**, (1995) 6532
61. M. Drees, Phys. Lett. B**181**, (1986) 279; J.S. Hagelin and S. Kelly, Nucl. Phys. **B342**, (1990) 95; Y. Kawamura, H. Murayama, and M. Yamaguchi, Phys. Rev. D**51**, (1995) 1337
62. Amitava Datta, Aseshkrishna Datta, and M.K. Parida, Phys. Lett. B**431**, (1998) 347; Amitava Datta, Aseshkrishna Datta, M. Drees, and D.P. Roy, Phys. Rev. D**61**, (2000) 055003
63. S. Bhattacharya, Aseshkrishna Datta, and B. Mukhopadhyaya, Phys. Rev. D**78**, (2008) 035011
64. N. Arkani-Hamed and S. Dimopoulos, J. High Energy Phys. **06**, (2005) 073; G.F. Giudice and A. Romanino, Nucl. Phys. B**699**, (2004) 65; *Erratum-ibid.* B **706**, (2005) 487
65. K. Huitu, J. Laamanen, P. Roy, and S. Roy, Phys. Rev. D**72**, (2005) 055002
66. J.L. Hewett *et al.*, J. High Energy Phys. **09**, (2004) 070
67. W. Kilian, T. Plehn, P. Richardson, and E. Schmidt, Eur. Phys. J. **C39**, (2005) 229