Abstract

The motivation for introduction of supersymmetry in high energy physics as well as a possibility for supersymmetry discovery at LHC (Large Hadronic Collider) are discussed. The main notions of the Minimal Supersymmetric Standard Model (MSSM) are introduced. Different regions of parameter space are analyzed and their phenomenological properties are compared. Discovery potential of LHC for the planned luminosity is shown for different channels. The properties of SUSY Higgs bosons are studied and perspectives of their observation at LHC are briefly outlined.
1 Introduction

Supersymmetry or symmetry between bosons (particles with integer spin) and fermions (particles with half-integer spin) has been introduced in theoretical papers nearly 30 years ago [1]. Since that time there appeared thousands of papers, all quantum field theory models were supersymmetrized, new mathematical tools were derived that allow one to work with anticommuting variables. The reason for this remarkable activity is the unique mathematical nature of supersymmetric theories, possible solution of various problems of the Standard Model of fundamental interactions within its supersymmetric extensions as well as the opening perspective of unification of all interactions in the framework of a single theory [2].

Supersymmetry today is the main candidate for a unified theory beyond the Standard Model. Search for various manifestations of supersymmetry in Nature is one of the main tasks of numerous experiments at colliders and in non-accelerator experiments of the last decade. Unfortunately, the result is negative so far. There are no any direct indications on existence of supersymmetry in particle physics though existing supersymmetric models satisfy all theoretical and experimental requirements. Remarkably that the scale of supersymmetry breaking, or as it is often said the scale of new physics, is about 1 TeV what is 10 times bigger than the electroweak symmetry breaking scale at which the LEP accelerator was adjusted. And it is this energy scale that the LHC accelerator will explore. It is assumed that at LHC the TeV energy range will be examined in detail, the Higgs boson will be found and supersymmetry will be discovered.

Supersymmetry is the challenge for the world physics community which was accepted with construction of LHC. Thus, high energy physics approaches the crucial moment when low energy supersymmetry will be either discovered or abandoned. One has to be ready for such circumstances and clearly realize which signatures of supersymmetry one can expect and how to extract them from that sea of data which will be obtained at the two main detectors of LHC: ATLAS and CMS.

2 Motivation of Supersymmetry

Recall what are the main arguments in favour of supersymmetric extension of the Standard Model of fundamental interactions. Though these arguments are not new their attractiveness does not weaken with time. They include

- unification with gravity. This is perhaps the main argument in favour of supersymmetry within the unification paradigm. The point is that SUSY algebra being a generalization of Poincaré algebra links together representations with different spins. The key relation is given by the anticommutator

\[ \{ Q_\alpha, \bar{Q}_{\dot{\alpha}} \} = 2 \sigma_\mu \alpha \bar{\alpha} P_\mu. \]

Taking infinitesimal transformations \( \delta_\epsilon = e^\alpha Q_\alpha, \ \bar{\delta}_{\dot{\epsilon}} = \bar{Q}_{\dot{\alpha}} \bar{\epsilon}^\dot{\alpha} \), one gets

\[ \{ \delta_\epsilon, \bar{\delta}_{\dot{\epsilon}} \} = 2(\epsilon \sigma_\mu \bar{\epsilon}) P_\mu, \] (1)
where $\epsilon$ is a transformation parameter. Choosing $\epsilon$ to be local, i.e. a function of a space-time point $\epsilon = \epsilon(x)$, one finds from eq.(11) that an anticommutator of two SUSY transformations is a local coordinate translation. And a theory which is invariant under the general coordinate transformation is General Relativity. Thus, making SUSY local, one obtains General Relativity, or a theory of gravity, or supergravity [3].

• unification of gauge couplings. According to hypothesis of Grand Unification gauge symmetry increases with energy. All known interactions are the branches of a single interaction associated with a simple gauge group which includes the group of the SM as a subgroup. Unification (or splitting) occurs at very high energy ($10^{15} \div 10^{16}$ GeV).

To reach this goal one has to examine how the coupling change with energy. This is described by the renormalization group equations. In the leading order of perturbation theory solutions take a simple form:

$$\frac{1}{\alpha_i(Q^2)} = \frac{1}{\alpha_i(\mu^2)} - b_i \log \left( \frac{Q^2}{\mu^2} \right),$$

where index $i$ refers to the gauge groups $SU(3) \times SU(2) \times U(1)$, and for the SM one has $b_i = (41/10, -19/6, -7)$. Result is demonstrated in Fig. 1 where the evolution of the inverse couplings is shown as functions of log of energy. In the left part of Fig. 1 one can see that in the SM unification of the gauge couplings is impossible. In the supersymmetric case the slopes of RG curves are changed, for the minimal supersymmetric extension of the SM one has $b_i = (33/5, 1, -3)$. It happens that in supersymmetric model one can achieve perfect unification as it is shown in the right part of Fig. 1. Fitting the curves one can get the scale of SUSY breaking $M_{SUSY} \sim 1 T$ [4].

• solution of the hierarchy problem. The appearance of two different scales in Grand Unified theories, namely $M_Z \ll M_{GUT}$, leads to the very serious problem which is called the hierarchy problem. First, this is the very existence of the hierarchy. Second, is the preservation of a given hierarchy in presence of the radiative corrections. These corrections, proportional to the mass of a heavy particle, inevitably destroy the hierarchy unless they are cancelled. The only way to get this cancellation of quadratic mass terms (also known as the cancellation of quadratic divergences) is supersymmetry. Moreover, supersymmetry automatically cancels all quadratic corrections in all orders of perturbation theory due to the contributions of superpartners of the ordinary particles. The contributions of the boson loops are cancelled by those of fermions due to additional factor $(-1)$ coming from Fermi statistic. This cancellation is true up to the SUSY breaking scale, $M_{SUSY}$, since

$$\sum_{bosons} m^2 - \sum_{fermions} m^2 = M^2_{SUSY},$$  

(3)
Unification of the Coupling Constants in the SM and the minimal MSSM

![Graph showing the evolution of the inverse gauge couplings in the SM (left) and in the MSSM (right).]

Figure 1: Evolution of the inverse gauge couplings in the SM (left) and in the MSSM (right).

which should not be very large ($\leq 1$ TeV) to make the fine-tuning natural. Indeed, let us take the Higgs boson mass. Requiring for consistency of perturbation theory that the radiative corrections to the Higgs boson mass do not exceed the mass itself gives

$$\delta M_h^2 \sim g^2 M_{SUSY}^2 \sim M_h^2.$$ (4)

Thus, we again get the same rough estimate of $M_{SUSY} \sim M_Z/g \sim 10^3$ GeV as from the gauge coupling unification above. Two requirements match together.

The origin of the hierarchy is the other part of the problem. We show below how SUSY can explain this part as well.

- radiative electroweak symmetry breaking. The "running" of the Higgs masses leads to the phenomenon known as **radiative electroweak symmetry breaking**. Indeed, one can see from Fig. 2 that the mass parameters from the Higgs potential $m_1^2$ and $m_2^2$ (or one of them) decrease while running from the GUT scale to the scale $M_Z$ may even change the sign. As a result for some value of the momentum $Q^2$ the potential may acquire a nontrivial minimum. This triggers spontaneous breaking of $SU(2)$ symmetry. The vacuum expectations of the Higgs fields acquire nonzero values and provide masses to quarks, leptons and $SU(2)$ gauge bosons, and additional masses to their superpartners. Thus, the breaking of the electroweak symmetry is not introduced by a brute force as in the SM, but appears naturally from the radiative corrections. In this way
one also obtains the explanation of why the two scales are so much different. Due to the logarithmic running of the parameters, one needs a long "running time" to get $m^2_2$ (or both $m^2_1$ and $m^2_2$) to be negative when starting from a positive value of the order of $M_{SUSY} \sim 10^2 \div 10^3$ GeV at the GUT scale.

- Dark matter in the Universe. The visible (or shining) matter is not the only matter in the Universe. Considerable amount of matter is the so-called Dark matter. Direct indication on the existence of the Dark matter are rotation curves of spiral galaxies. To explain these curves one usually assumes the existence of a galactic halo consisting of non-shining matter which takes part in gravitational interaction. According to recent data \[6\], the matter content of the Universe is the following:

$$\Omega_{total} = 1.02 \pm 0.02$$

$$\Omega_{vacuum} = 0.73 \pm 0.04, \quad \Omega_{matter} = 0.23 \pm 0.04, \quad \Omega_{barion} = 0.044 \pm 0.004\%,$$

i. e. Dark matter makes up a considerable part exceeding the visible barionic matter by the order of magnitude.

There are two possible kinds of nonbarionic Dark matter: hot DM, consisting of light relativistic particles, and cold DM, consisting of weakly interacting massive particles (WIMPs). Hot DM might consist of neutrino, however, this is problematic from the point of view of large structure formation in the Universe. Besides, neutrinos are too light to produce enough DM. As for the cold DM, in the SM there are no appropriate particles. At the same time, supersymmetry provides an excellent candidate for this role, namely, neutralino, the lightest
superparticle. It is stable, so that the relic neutralinos might survive in the Universe since the Big Bang.

3 MSSM: the field content and Lagrangian

Despite complexity of the mathematical structure of supersymmetric gauge theories, any supersymmetric extension of the Standard Model has some general simple features which do not depend on a particular model. This is, first of all, the doubling of particles: each particle of the SM, quark or lepton or the gauge boson like photon, gluon or intermediate weak boson, has a partner with the same quantum numbers but with the spin differing by $1/2$. These particles are called superpartners. Note that the usual particles of the SM can not be partners of each other since in the SM one has no particles with the same quantum numbers and different spin.

The field content of the Minimal Supersymmetric Standard Model (MSSM) is shown in Table 1 (hereafter the tilde over the symbol of a particle denotes a superpartner of a usual particle).

| Superfield | Bosons | Fermions | $SU_c(3)$ | $SU_L(2)$ | $U_Y(1)$ |
|------------|--------|----------|-----------|-----------|----------|
| $G^a$      | gluon  | $g^a$    | gluino    | $\tilde{g}^a$ | 8        | 1        | 0         |
| $V^k$      | Weak   | $W^k$    | wino, zino| $\tilde{w}^k$ (\$\tilde{w}^\pm, \tilde{z}\$) | 1        | 3        | 0         |
| $V'$       | Hypercharge | $B(\gamma)$ | bino | $\tilde{b}(\gamma)$ | 1        | 1        | 0         |

| Matter     |        |          |           |           |          |
|------------|--------|----------|-----------|-----------|----------|
| $L_i$      | sleptons | $\bar{L}_i = (\bar{\nu}, \bar{e})_L$ | leptons   | $L_i = (\nu, e)_L$ | 1        | 2        | $-1$      |
| $E_i$      |         | $\bar{E}_i = \bar{e}_R$ |           | $E_i = e_R$ | 1        | 1        | 2         |
| $Q_i$      |         | $\bar{Q}_i = (\bar{u}, \bar{d})_L$ |           | $Q_i = (u, d)_L$ | 3        | 2        | $1/3$     |
| $U_i$      | squarks | $\bar{U}_i = \bar{u}_R$ | quarks    | $U_i = u_R$ | $3^*$    | 1        | $-4/3$    |
| $D_i$      |         | $\bar{D}_i = \bar{d}_R$ |           | $D_i = d_R$ | $3^*$    | 1        | $2/3$     |

| Higgs      |        |          |           |           |          |
|------------|--------|----------|-----------|-----------|----------|
| $H_1$      | Higgses | $H_1$    | higgsinos | $\tilde{H}_1$ | 1        | 2        | $-1$      |
| $H_2$      |         | $H_2$    |           | $\tilde{H}_2$ | 1        | 2        | 1         |

The labels L or R for squarks or sleptons do not mean that they are left or right handed. Being spin zero particles they have no handedness. This is used to mark that they are superpartners of left or right handed quarks and leptons.

The presence of the additional Higgs boson is a generic property of the supersymmetric theory. In the MSSM there are two doublets of scalar fields with quantum numbers $(1, 2, -1)$ and $(1, 2, 1)$. 


Table 1 does not contain gravitational fields. In the simplest version of supergravity one has to add to the set of the MSSM fields the pair of graviton and gravitino, the particle with spin 3/2.

The other important feature is the breaking of supersymmetry. If this does not happen the superpartners would be degenerate in masses with the ordinary particles, which is not observed. Due to supersymmetry breaking this degeneracy disappears and superpartners acquire large masses that explains their non-observation at the moment. However, the trace of supersymmetry should remain in relations between the amplitudes of various processes (with participation of the usual particles and superpartners) and in contributions of superpartners to the radiative corrections below the threshold. The concrete predictions depend on the details of supersymmetry breaking mechanism which is not known yet.

The Lagrangian of the MSSM consists of two parts; the first part is SUSY generalization of the Standard Model, while the second one represents the SUSY breaking as mentioned above. The supersymmetric part of the Lagrangian consists of the gauge invariant kinetic terms corresponding to the $SU(3)$, $SU(2)$, $U(1)$ gauge groups depending on 3 gauge couplings as in the Standard Model and of the superpotential. Usually the superpotential is chosen in the form repeating that of the Yukawa interaction in the SM

$$W = \epsilon_{ij} (y^U_{ab} Q_a^i U^c_b H^i_2 + y^D_{ab} Q_a^i D^c_b H^i_1 + y^L_{ab} L^i_a E^c_b H^i_1 + \mu H^i_1 H^j_2),$$

where $i, j = 1, 2, 3$ are the $SU(2)$ and $a, b = 1, 2, 3$ are the generation indices; colour indices are suppressed, $y^{U,D,L}$ are the Yukawa couplings. This part of the Lagrangian almost exactly repeats that of the SM except that the fields are now the superfields rather than the ordinary fields of the SM. The only difference is the last term which describes the Higgs mixing. It is absent in the SM since there is only one Higgs field there.

In principle the superpotential can contain other interactions:

$$W_{NR} = \epsilon_{ij} (\lambda^L_{a b d} L^i_a L^j_b E^c_d + \lambda^L_{a b d} L^i_a Q^j_b D^c_d + \mu' L^i_a H^j_2 + \lambda^B_{a b d} U^c_a D^c_b D^c_d).$$

These terms are absent in the SM. The reason is very simple: one can not replace the superfields in eq. (5) by the ordinary fields like in eq. (5) because of the Lorentz invariance. These terms have a different property; they violate either lepton (the first three terms in eq. (5)) or baryon number (the last term). Since both effects are not observed in Nature, these terms must be suppressed or excluded. One can avoid such terms by introducing a special symmetry called the $R$-parity defined by

$$R = (-1)^{3(B-L)+2S}$$

where $B$ — is the baryon number, $L$ — is the lepton number, and $S$ — is the spin of the particle. Conservation of the $R$-parity has important phenomenological consequences: superparticles are created in pairs and the lightest supersymmetric particle (LSP) is absolutely stable. This makes the LSP an excellent candidate
for the Dark matter particle that is one of attractive features of supersymmetric extension of the SM.

Since none of the fields of the MSSM can have nonzero vacuum expectation value, needed for the SUSY breaking, without violating the gauge invariance, it is assumed that spontaneous breaking of supersymmetry takes place with the help of some other fields. The most popular scenario of getting low-energy broken supersymmetry is the so-called hidden sector scenario \[\text{s}\]. According to it there are two sectors: the usual matter belongs to the "visible" sector, while the second, "hidden" sector, contains the fields that break supersymmetry. These two sectors interact with each other by exchanging some fields called messengers. They transport supersymmetry breaking from the hidden sector to the visible one. The messengers may be various fields: gravitons, gauge bosons, etc.

SUSY breaking terms of the Lagrangian are often called the soft terms since they are the operators of dimension 2 and 3. They contain a vast number of free parameters which spoils the predictive power of the model. To reduce their number, we adopt the so-called universality hypothesis, i.e., we assume the universality or equality of various soft parameters at the high energy scale, namely, we put all the spin 0 particle masses to be equal to the universal value \(m_0\), all the spin 1/2 particle (gaugino) masses to be equal to \(m_{1/2}\) and all the cubic and quadratic terms repeat the structure of the Yukawa superpotential \[\text{h}\]. This is an additional requirement motivated by the supergravity mechanism of SUSY breaking. Universality is not a necessary requirement and one may consider nonuniversal soft terms as well. However, it will not change the qualitative picture presented below; so for simplicity, in what follows we consider the universal boundary conditions. In this case, the soft terms take the form

\[
L_{\text{Breaking}} = m_0^2 \sum_i |\varphi_i|^2 + \left( \frac{1}{2} m_{1/2} \sum_a \lambda_a \lambda_a \right)
\]

\[
+ A[|y^{U}_{ab} \tilde{Q}_a \tilde{U}_b^c|H_2 + |y^{D}_{ab} \tilde{Q}_a \tilde{D}_b^c|H_1 + |y^{L}_{ab} \tilde{L}_a \tilde{E}_b^c|H_1] + B[|\mu H_1 H_2| + h.c.) ,
\]

where \(\varphi\) denote the scalar fields of squarks, sleptons and the Higgs bosons, \(\lambda\) — are the gauginos, the spinor superpartners of the gauge fields, \(A\) and \(B\) — are the new parameters of dimension of a mass.

Thus, the Minimal Supersymmetric Standard Model has the following free parameters: i) three gauge couplings \(\alpha_i\); ii) three matrices of the Yukawa couplings \(y^i_{ab}\), where \(i = L, U, D\); iii) the Higgs field mixing parameter \(\mu\); iv) the soft supersymmetry breaking parameters. Compared to the SM there is an additional Higgs mixing parameter, but the Higgs self-coupling, which is arbitrary in the SM, is fixed by supersymmetry. The main uncertainty comes from the unknown soft terms.

With the universality hypothesis one is left with the following set of 5 free parameters defining the mass scales

\[
\mu, \ m_0, \ m_{1/2}, \ A \text{ and } B \leftrightarrow \tan \beta = \frac{v_2}{v_1}
\]

8
Instead of parameter $B$ one usually uses the parameter $\tan \beta \equiv v_2/v_1$ equal to the ratio of v.e.v's of the Higgs fields. Choosing the values of these free parameters one can predict the mass spectrum of superpartners and the cross-sections for their production.

4 Supersymmetry breaking: the parameter space

To reduce arbitrariness in the choice of the MSSM parameters and to make more definite predictions one usually imposes several constraints which also serve as a consistency checks of the model. As it happens, in the MSSM one can simultaneously fulfil several such constraints:

- Gauge coupling constant unification. This is one of the most restrictive constraints. It fixes the scale of SUSY breaking of the order of 1 TeV.

- $M_Z$ from electroweak symmetry breaking:
  Radiative EW symmetry breaking defines the mass of the Z-boson
  \[
  \frac{M_Z^2}{2} = \frac{m_1^2 - m_2^2 \tan^2 \beta}{\tan^2 \beta - 1} = -\mu^2 + \frac{m_{H_1}^2 - m_{H_2}^2 \tan^2 \beta}{\tan^2 \beta - 1}. \tag{9}
  \]
  This condition determines the value of $\mu^2$ for given values of $m_0$ and $m_{1/2}$.
  The sign of $\mu$ remains undefined, it can be fixed from the other constraints.

- Yukawa coupling constant unification. The masses of top, bottom and $\tau$ can be obtained from the low energy values of the running Yukawa couplings via
  \[
  m_t = y_t \ v \sin \beta, \quad m_b = y_b \ v \cos \beta, \quad m_\tau = y_\tau \ v \cos \beta. \tag{10}
  \]
  They can be translated to the pole masses taking into account the radiative corrections, which restricts possible solutions in the Grand Unified Theories.

- Precision measurement of decay rates. Radiative corrections due to superpartners may essentially influence the decay rates under the threshold. The typical example is the branching ratio $BR(b \to s\gamma)$ which has been measured by BaBar, CLEO and BELLE collaborations [9] and yields the world average of $BR(b \to s\gamma) = (3.43 \pm 0.36) \cdot 10^{-4}$. The Standard Model contribution to this process gives slightly lower result, thus leaving a window for SUSY. This requirement imposes severe restrictions on the parameter space, especially for the case of large $\tan \beta$.

- Anomalous magnetic moment of muon. Recent measurement of the anomalous magnetic moment indicates small deviation from the SM of the order of $2 \sigma$ [10]:
  \[
  \Delta a_\mu = a_{\mu}^{exp} - a_{\mu}^{theor} = (27 \pm 10) \cdot 10^{-10}. \]
  The deficiency may be easily filled with SUSY contribution, which is proportional to $\mu$ and $\tan \beta$. This requires positive sign of $\mu$ and kills a half of the parameter space of the MSSM [11].
- Experimental lower limits on SUSY masses. SUSY particles have not been found so far and from the searches at LEP one knows the lower limit on the charged lepton and chargino masses of about one half of the centre of mass energy \[12\]. The lower limit on the neutralino masses is smaller. There exist also limits on squark and gluino masses from the hadron colliders \[13\]. These limits restrict the minimal values for the SUSY mass parameters.

- Dark Matter constraint. Recent precise astrophysical data restrict the amount of the Dark Matter in the Universe to \(23 \pm 4\%\). Assuming that the Hubble constant is \(h_0 \approx 0.7\) one finds that the contribution of each relict particle \(\chi\) has to obey the constraint \(\Omega_\chi h_0^2 \sim 0.12 \pm 0.02\) and serve as a very severe bound on SUSY parameters \[14\] that leaves a very narrow band of allowed region in parameter space.

Requirement of simultaneous fulfilment of these constraints defines the allowed regions of parameter space. However, not all of the above mentioned parameters are equally important. Besides, some of them are practically not free, since they are severely constrained. For example, as we already mentioned, the Higgs mixing parameter \(\mu\) is related to \(m_0\), \(m_{1/2}\) and \(Z\)-boson mass. The triple coupling \(A\) in many cases is inessential and its value at the GUT scale is often chosen to be \(A_0 = 0\). The requirement of Yukawa coupling unification restricts the value of \(\tan \beta\).

There are two possible scenaria: scenario with small \(\tan \beta\) (\(\tan \beta \approx 1 \div 3\)) and scenario with large \(\tan \beta\) (\(\tan \beta \approx 30 \div 70\)) \[15\]. These scenaria are rather different from phenomenological point of view since the allowed regions of parameter space are different. Unfortunately, the recent LEP data practically exclude the small \(\tan \beta\) scenario since the mass of the lightest Higgs boson happens to be below the experimental limit. Besides, the astrophysical data are also in favour of large \(\tan \beta\).

Thus, from the set of free parameters of the model we basically have two independent ones: \(m_0\) and \(m_{1/2}\). It is very useful therefore to use the \((m_0, m_{1/2})\) plane to present the theoretical and experimental constraints that we discussed above. Since the scale of supersymmetry breaking is of the order of 1 TeV, the masses of superpartners should be in the same region, which defines the range of parameters \(m_0\) and \(m_{1/2}\).

We consider further how each of the above mentioned constraints cuts out the allowed regions in the parameter space in the plane \((m_0, m_{1/2})\). In Fig. 3 these regions are shown for two fixed values of \(\tan \beta = 35\) and \(\tan \beta = 50\) without account of astrophysical data yet for the values of \(m_0\) and \(m_{1/2}\) in the interval from 200 to 1000 GeV. We start with the radiative electroweak symmetry breaking. It happens so that for very large \(m_0\) and small \(m_{1/2}\) the mass parameters of the scalar potential, which start from \(m_0\), do not have enough ”time” to run to negative values when the conditions of non-trivial minima of the potential are satisfied. Therefore the left low corner of \((m_0, m_{1/2})\) plane is usually excluded by this requirement.

Similar thing happens with the constraint related to the non-observation of the Higgs boson, which forbids the region of small values of \(m_{1/2}\) practically independently of \(m_0\) for fixed \(\tan \beta\). As for the \(\tan \beta\) dependence it is the following: the smaller the value of \(\tan \beta\) the large the excluded region. This fact basically excludes
Figure 3: The allowed regions of parameter space for the large tan $\beta$ scenario for different constraints (left – tan $\beta = 35$, right – tan $\beta = 50$).

Figure 4: Constraints on the parameter space from the requirement of right amount of the Dark matter (left - tan $\beta = 35$, right - tan $\beta = 50$).

the small tan $\beta$ scenario as incompatible with the LEP lower limit on the Higgs boson mass $m_h \geq 114.3$ GeV.

The small values of $m_{1/2}$ do not satisfy also the constraint coming from the rate of the rare decay $BR(b \rightarrow X_s\gamma)$. However, in this case the dependence on tan $\beta$ is opposite to that of the previous case. In the case of tan $\beta = 35$ only a small part of the parameter space is forbidden ($m_0, m_{1/2} \leq 300$ GeV), while for the large value of tan $\beta = 50$ the forbidden region is $m_{1/2} \leq 300 \div 400$ GeV for any values of $m_0$. This constraint happens to be more restrictive for large tan $\beta$ than the one related to the Higgs boson mass.
Constraint on experimental value of the anomalous magnetic moment of muon leaves the allowed band in the \((m_0, m_{1/2})\) plane the width of which depends on \(\tan \beta\). For \(\tan \beta = 35\) the excluded regions are the left low corner \((m_0, m_{1/2} \leq 300 \text{ GeV})\) (as for the \(BR(b \to s\gamma)\) constraint) and almost all right upper part of the plane. In this case we get a restriction from above on the masses of superpartners. For \(\tan \beta = 50\) small values of \(m_0\) and \(m_{1/2}\) are also excluded, however, the upper bound is shifted to the region above 1 TeV, allowing heavier superpartners.

In the case of \(R\)-parity conservation the lightest superparticle (LSP) is usually the neutralino, a certain mixture of superpartners of the photon, \(Z\)-boson and neutral Higgs bosons. And it is stable! However, in large amount of the parameter space (left upper corner of the plane where \(m_0 < m_{1/2}\)) the superpartner of \(\tau\)-lepton becomes even lighter and hence have to be stable. But we would have registered a stable charged particle if it exists even if it is heavy. Therefore the requirement of neutrality of LSP have to be fulfilled. The region which is excluded by this requirement also depends on \(\tan \beta\): the larger is \(\tan \beta\) the larger is the excluded region.

The allowed regions of the parameter space in the plane \((m_0, m_{1/2})\) which are left after taking into account all mentioned constraints are shown in Fig. 3 for two values of \(\tan \beta = 35\) and \(\tan \beta = 50\).

Figure 5: The light (blue) band is the region allowed by the WMAP data for \(\tan \beta = 51, \mu > 0\) and \(A_0 = 0.5m_0\). The excluded regions where the LSP is stau (red up left corner), where the radiative electroweak symmetry breaking mechanism does not work (red low right corner), and where the Higgs boson is too light (yellow lower left corner) are shown with dots. The numbers denote: 1 — the main annihilation region, 2 — the co-annihilation region, 3 — the focus point region, 4 — the funnel region, 5 — the EGRET region.
Constraint on the amount of the Dark matter in the interval $\Omega h^2 = 0.1 \div 0.3$, in its turn, cuts out narrow bands in the $(m_0, m_{1/2})$ plane as is shown in Fig. 4 for two values of $\tan \beta$.

With account of recent precise data from WMAP collaboration (Wilkinson Microwave Anisotropy Probe) [6] one has much more severe constraints. As a result the allowed regions in the $(m_0, m_{1/2})$ plane are along the narrow band shown in Fig. 5 [16]. Note that to satisfy WMAP requirement one prefers large values of $\tan \beta \approx 50$. Another comment concerns the importance of the upper limit coming from WMAP data, the lower limit may be influenced by the other unknown particles and invisible macro objects.

The constrained MSSM possesses already high predictive power. In the regions of parameter space where there is no contradiction with experimental data or theoretical requirements one can get the mass spectra of superpartners and the Higgs bosons and to indicate the possibilities if their experimental search.

5 Possible scenaria

The relation between $m_0$ and $m_{1/2}$ inside varies along the the narrow allowed band. Respectively vary the mass spectrum of superpartners, the dominance of different creation and decay processes, the values of the cross sections, the methods of analysis of experimental data at LHC. Remind that part of parameters are practically fixed in a sense that one can choose their values to satisfy the imposed constraints with maximal probability.

Consider several cosmologically acceptable and phenomenologically different regions along the WMAP band. They have a certain mass spectrum typical to the each region that defines the main production and annihilation and/or co-annihilation channels for the neutralino [17].

- The first, mostly studied region is the bulk annihilation region, this is the region of relatively small $m_0$ and $m_{1/2}$ ($m_0 \approx 50 \div 150$ GeV, $m_{1/2} \approx 50 \div 350$ GeV). It is bounded from below by the non-observation of the Higgs boson and the absence of radiative electroweak symmetry breaking as well as by consistency with the $b \to s\gamma$ decay rate. From the left there is a forbidden region where stau is the LSP.

One of the main processes in this region is the annihilation of pair of neutralinos into quarks through the exchange of a squark in the $t$-channel $\tilde{\chi}_1^0\tilde{\chi}_1^0 \to q\bar{q}$. The parameters can be adjusted in a way to give the right amount of the Dark matter.

The size of the region depends on $\tan \beta$ and for low values of $\tan \beta$ it practically disappears due to the non-observation of the Higgs boson.

- The other interesting region is the so-called stau co-annihilation region. Here typically one has small values of $m_0$ and much bigger values of $m_{1/2}$. It is located along the border line between the regions where $\tilde{\tau}_1$-slepton is the LSP
and neutralino $\tilde{\chi}_1^0$ is the LSP. Evidently this corresponds to the case when the particles are almost degenerate in masses $m_{\tilde{\chi}_1^0} \approx m_{\tilde{\tau}_1}$ and in the early Universe there were co-annihilation processes $\tilde{\chi}_1^0\tilde{\tau}_1 (\tilde{\chi}_1^0\tilde{\tau}_1 \to \tau^+ \to \tau \gamma)$ as well as co-annihilation $\tilde{\tau}_1\tilde{\tau}_1$. Neutralino in this case is mostly higgsino and its mass may be large up to 500 GeV without violating the WMAP bound.

Co-annihilation region is interesting from the point of view of existence of long-lived charged sleptons. Their life-time may be large enough to be produced in proton-proton collisions and to fly away from the detector area or to decay inside the detector at a considerable distance from the collision point. Clearly that such an event can not be unnoticed. However, to realize this possibility one need a fine-tuning of the parameters of the model [18].

- As has been already mentioned for large $m_0$ small values of $m_{1/2}$ are forbidden due to the absence of radiative electroweak symmetry breaking. However, along the border of the forbidden region, the WMAP allowed band may stay long enough leading to the masses of squarks and sleptons up to a few TeV. This region is called the focus point region since the values of the Higgs mass parameters here tend to the focus point when running the renormalization group equations. In this region the Higgs mixing parameter $\mu$ happens to be small $|\mu| \sim M_Z$. Then it is possible that two light neutralinos and the light chargino are practically degenerate $m_{\chi_1^0} \sim m_{\chi_2^0} \sim m_{\chi_1^\pm} \sim \mu$. The lightest neutralino in this case is mostly higgsino. The main annihilation channel is the one into the pair of gauge bosons $\chi_1^0\chi_1^0 \to ZZ$ or $\chi_1^0\chi_1^0 \to W^+W^-$ but due to degeneracy of masses of neutralino and chargino there are also possible the co-annihilation processes $\chi_1^0\chi_1^\pm$, $\chi_1^0\chi_2$, $\chi_1^\pm\chi_1^-$ and $\chi_2^0\chi_1^\pm$.

Despite the large values of $m_{1/2}$ up to 1 TeV, $\mu$ remains small and leads to chargino and neutralino masses of the order of a few hundreds GeV. This tells us that the focus point region is accessible by LHC. Even the cross section of gluino pair production is big enough to observe this process.

- For large values of $\tan \beta$ it is possible that $m_A \approx 2m_{\chi_1^0}$. There is no need for precise equality because the width of the $CP$-odd Higgs boson $A$ is about tens of GeV. In this region of $(m_0, m_{1/2})$ plane the allowed WMAP band has a sharp turn in the form of a funnel: $A$-funnel region. The main channel of annihilation in this region is $\chi_1^0\chi_1^0 \to A \to b\bar{b}$ or $\tau\bar{\tau}$. The reason for such a behaviour is that for increasing $\tan \beta$ the mass of a pseudoscalar Higgs boson $A$ decreases while the mass of neutralino practically does not change. Then inevitably the resonance situation when $m_A = 2m_{\chi_1^0}$ occurs. And despite that fact that neutralino in this case is almost photino and does not interact with the Higgs boson $A$, the tiny admixture of higgsino leads to considerable effect due to relatively big coupling of the $A$-boson to quarks and leptons $Ab\bar{b}$ $A\tau\bar{\tau}$. For the same reason the exchange of the heavy Higgs boson $H$ might give an essential contribution.
Besides, in this region the cross section of neutralino $\chi^0_1$ scattering on the nucleus is of the order of $10^{-8} \div 10^{-9}$ pb which is close to the values corresponding to the sensitivity of the modern and the nearest future experiments on the direct Dark matter searches.

In addition to the above mentioned regions there are some small exotic ones. For example, for a specific choice of parameters (very big $A_0$, moderate or big $m_0$ and small $m_{1/2}$) as a result of mixing one of the $t$-squarks becomes practically degenerate with the lightest neutralino $\chi^0_1$. In this case the process of $\tilde{\chi}^0_1 \tilde{t}_1$ co-annihilation is possible. For small values of $m_{1/2}$ (and for appropriate choice of the other parameters) there is a possibility of neutralino annihilation due to light Higgs boson exchange in the $s$-channel. This situation is analogous to that of annihilation through $A$ or $H$.

- One should mention the other interesting constraint on the parameter space of the MSSM related to the supersymmetric interpretation of the excess of the diffuse gamma ray flux in our Galaxy compared to the background calculations. This is the data presented by EGRET collaboration (Energetic Gamma Ray Experiment Telescope) \[^{19}\]. Omitting the details we notice, that it is enough to assume the existence of a neutral stable weakly interacting particle (WIMP) of a certain mass to explain this excess. The fit gives the value of this mass in the region $m_X \approx 50 - 100$ GeV \[^{20}\]. If one takes this WIMP to be the lightest neutralino this will strongly constrain the value of parameter $m_{1/2}$. Moreover, this constraint is compatible with WMAP. The allowed area is in the region of $m_0 \approx 1400$ GeV and $m_{1/2} \approx 180$ GeV \[^{21}\], i.e. practically between the bulk annihilation region and then focus point region.

6 Search for supersymmetry at LHC

The strategy of SUSY searches at LHC is based on the assumption that the masses of superpartners indeed are in the region of 1 TeV so that they might be created at the mass shell with the cross section big enough to distinguish them from the background of the ordinary particles. Calculation of the background in the framework of the Standard Model thus becomes essential since the secondary particles in all the cases will be the same.

There are many possibilities to create superpartners at hadron colliders. Besides the usual annihilation channel there are numerous processes of gluon fusion, quark-antiquark and quark-gluon scattering. The maximal cross sections of the order of a few picobarn can be achieved in the process of gluon fusion.

As a rule all superpartners are short lived and decay into the ordinary particles and the lightest superparticle. The main decay modes of superpartners, i.e. experimental manifestation of SUSY at LHC are presented in Table 2.

Notice the typical events with missing energy and transverse momentum that is the main difference from the background processes of the Standard Model.
The main decay modes

| Creation | The main decay modes | Signature |
|----------|----------------------|-----------|
| \( \tilde{g}, \tilde{g}, \tilde{g} \) | \( g \rightarrow q\bar{q}' \chi_1^0 \) \( m_\tilde{g} > m_\tilde{q} \) | \( E_T \) + multijets (+leptons) |
| \( \tilde{\chi}_1^+ \tilde{\chi}_2^- \) | \( \tilde{\chi}_1^+ \rightarrow \chi_1^0 \ell^+ \nu, \tilde{\chi}_2^- \rightarrow \chi_2^0 \ell \) | trilepton + \( E_T \) |
| \( \tilde{\chi}_1^0 \tilde{\chi}_1^0 \) | \( \tilde{\chi}_1^0 \rightarrow \chi_1^0 X, \tilde{\chi}_1^0 \rightarrow \chi_1^0 X' \) | dilepton + jet + \( E_T \) |
| \( \tilde{t}_1 \tilde{t}_1 \) | \( \tilde{t}_1 \rightarrow c \chi_1^0 \) | 2 noncollinear jets + \( E_T \) |
| \( \tilde{t}_1, \tilde{\nu}, \tilde{\ell} \) | \( \tilde{t}_1 \rightarrow b \tilde{\chi}_1^+, \tilde{\chi}_1^- \rightarrow \chi_1^0 q \bar{q}' \) | single lepton + \( E_T \) + b's |
| | \( \tilde{t}_1 \rightarrow b \tilde{\chi}_1^+, \tilde{\chi}_1^- \rightarrow \chi_1^0 \ell^+ \nu \) | dilepton + \( E_T \) + b's |

missing energy is carried away by the heavy particle with the mass of the order of 100 GeV that is essentially different from the processes with neutrino in the final state. In hadron collisions the superpartners are always created in pairs and then further quickly decay creating a cascade with the ordinary quarks (i.e. hadron jets) or leptons at the end plus the missing energy. For the case of gluon fusion with creation of gluino it is presented in Table 4.

Chargino and neutralino can also be produced in pairs through the Drell-Yang mechanism \( pp \rightarrow \tilde{\chi}_1^\pm \tilde{\chi}_2^0 \) and can be detected via their lepton decays \( \chi_1^\pm \chi_2^0 \rightarrow \ell \ell + E_T \). Hence the main signal of their creation is the isolated leptons and missing energy (Table 4). The main background in trilepton channel comes from creation of the standard particles \( WZ/Z\bar{Z}, tt, Zbb \) and \( bb \). There might be also the supersymmetric background from the cascade decays of squarks and gluino into multilepton modes.

The cross sections for various superpartners creation at LHC are shown in Fig. 6. One can see that in some regions they may reach a few pb that is for a planned luminosity of LHC allows one to provide reliable detection. In the case of light neutralino and chargino the cross sections of their pair production can reach those of the strongly interacting particles [22]. To present the region of reach for the LHC in different channels of sparticle production one uses the same plane of soft SUSY breaking parameters \( m_0 \) and \( m_{1/2} \). In this case one usually assumes certain luminosity which will be presumably achieved during the accelerator operation.
Table 3: Creation of the pair of gluino with further cascade decay

| Process | final states |
|---------|-------------|
| $g \rightarrow \tilde{g} \tilde{g}$ | $2\ell$, $6j$, $\not{E}_T$ |

Thus, for instance, in Fig. 7 it is shown the regions of reach in different channels. The lines of a constant squark mass form the arch curves, and those for gluino are almost horizontal. The curved lines show the reach bounds in different channel of creation of secondary particles. The theoretical curves are obtained within the MSSM for a certain choice of the other soft SUSY breaking parameters. In the left plot the calculations are performed for the luminosity equal to $10^5 \text{ pb}^{-1}$ and in the right one for the luminosity $10^2 \text{ pb}^{-1}$ which will be presumably reached at the first stage. As one can see, for the fortunate circumstances the wide range of the parameter space up to the masses of the order of 2 Tev will be examined.

The other example is shown in Fig. 8 where the regions of reach for squarks and gluino are shown for various luminosities. One can see that for the maximal luminosity the discovery range for squarks and gluino reaches 3 TeV for the center of mass energy of 14 TeV and even higher for the double energy.

The same is true for the sleptons as shown in Fig. 9. The slepton pairs can be
Table 4: Creation of the lightest chargino and the second neutralino with further cascade decay.

| Process | final states |
|---------|--------------|
| $p (q)$ | $\ell \nu \ell \nu$ |
| $p (\bar{q})$ | $\ell \chi^0_1$ |

created via the Drell-Yang mechanism $pp \to \gamma^*/Z^* \to \ell^+\ell^-$ and can be detected through the slepton decays $\ell \to \ell + \chi^0_1$. The typical signal used for slepton detection is the dilepton pair with the missing energy without hadron jets. For the luminosity of $L_{tot} = 100$ fb$^{-1}$ the LHC will be able to discover sleptons with the masses up to 400 GeV \cite{24}.

We do not discuss here the different possibilities of detection of long lived supersymmetric particles, staus or supersymmetric hadrons. The very existence of these particles requires the fine tuning of parameters. However, if these particles exist, their decay inside the detector would give a characteristic signal with creation of a jet or a charged lepton at a point distinct from the collision point which might be detected.
Figure 6: The cross sections of superpartners creation as functions of $m_{1/2}$ and $m_0$ for $\tan \beta = 51$, $A_0 = 0$ and positive sign of $\mu$.

Figure 7: Expected range of reach for superpartners in various channels at LHC [23].
The CMS q, g mass reach in $E_T^{\text{miss}} + \text{jets}$ inclusive channel for various integrated luminosities at $\sqrt{s} = 28$ TeV.

Figure 8: Expected range of reach for squarks and gluino for different luminosities at LHC [23].

The CMS q, g mass reach in $E_T^{\text{miss}} + \text{jets}$ inclusive channel for various integrated luminosities at $\sqrt{s} = 28$ TeV.

Figure 9: Expected range of reach for sleptons at LHC [23].
SUSY Higgs boson

The presence of an extra Higgs doublet in SUSY model is a novel feature of the theory. In the MSSM one has two doublets with the quantum numbers (1,2,-1) and (1,2,1), respectively:

\[
H_1 = \left( \frac{H_1^0}{H_1^-} \right) = \left( v_1 + \frac{S_1 + iP_1}{\sqrt{2} v_1} \right), \quad H_2 = \left( \frac{H_2^+}{H_2^0} \right) = \left( v_2 + \frac{S_2 + iP_2}{\sqrt{2} v_2} \right),
\]

(11)

where \(v_i\) are the vacuum expectation values of the neutral components.

Thus, in the MSSM, as actually in any two Higgs doublet model, one has \(8=4+4=5+3\) degrees of freedom. As in the case of the SM, \(3\) degrees of freedom can be gauged away, and one is left with five physical Higgs bosons: two CP-even neutral, one CP-odd neutral and two charged ones.

The Higgs potential in the MSSM is totally defined by superpotential \(W_R\) and the soft terms (8). Due to the structure of \(W_R\) it contributes only to the mass matrix while the Higgs self-interaction comes from the interaction with the gauge fields. The tree level potential is

\[
V_{\text{tree}}(H_1, H_2) = m_1^2 |H_1|^2 + m_2^2 |H_2|^2 - m_3^2 (H_1 H_2 + h.c.) + \frac{g^2 + g'^2}{8} (|H_1|^2 - |H_2|^2)^2 + \frac{g^2}{2} |H_1 H_2|^2,
\]

(12)

where \(m_1^2 = m_{H_1}^2 + \mu^2, m_2^2 = m_{H_2}^2 + \mu^2\). At the GUT scale \(m_1^2 = m_2^2 = m_0^2 + \mu_0^2, m_3^2 = -B\mu_0\). Notice that the Higgs self-coupling in eq.(12) is fixed and defined by the gauge interactions as opposed to the SM.

The mass eigenstates are [2]:

\[
\begin{align*}
C^0 &= -\cos \beta P_1 + \sin \beta P_2, & \text{Goldstone boson} \rightarrow Z_0, \\
A &= \sin \beta P_1 + \cos \beta P_2, & \text{Neutral CP} = -1 \text{ Higgs},
\end{align*}
\]

\[
\begin{align*}
G^+ &= -\cos \beta (H_1^-)^* + \sin \beta H_2^+, & \text{Goldstone boson} \rightarrow W^+, \\
H^+ &= \sin \beta (H_1^-)^* + \cos \beta H_2^+, & \text{Charged Higgs}, \\
h &= -\sin \alpha S_1 + \cos \alpha S_2, & \text{SM Higgs boson CP} = 1, \\
H &= \cos \alpha S_1 + \sin \alpha S_2, & \text{Extra heavy Higgs boson},
\end{align*}
\]

where the mixing angle \(\alpha\) is given by

\[
\tan 2\alpha = \tan 2\beta \left( \frac{m_A^2 + M_Z^2}{m_A^2 - M_Z^2} \right).
\]

The physical Higgs bosons acquire the following masses [2]:

\[
\begin{align*}
\text{CP-odd neutral Higgs} \quad A : & \quad m_A^2 = m_1^2 + m_2^2, \\
\text{Charge Higgses} \quad H^\pm : & \quad m_{H^\pm}^2 = m_A^2 + M_W^2,
\end{align*}
\]

(13)
CP-even neutral Higgses $H, h$:

$$m_{H,h}^2 = \frac{1}{2} \left[ m_A^2 + M_Z^2 \pm \sqrt{(m_A^2 + M_Z^2)^2 - 4m_A^2M_Z^2\cos^22\beta} \right], \quad (14)$$

where, as usual,

$$M_W^2 = \frac{g^2}{2}v^2, \quad M_Z^2 = \frac{g^2 + g'^2}{2}v^2.$$ 

This leads to the once celebrated SUSY mass relations

$$m_{H^\pm} \geq M_W,$$
$$m_H \leq m_A \leq M_H,$$
$$m_h \leq M_Z|\cos 2\beta| \leq M_Z,$$
$$m_A^2 + m_H^2 = m_A^2 + M_Z^2.$$

Thus, the lightest neutral Higgs boson happens to be lighter than the $Z$-boson, which clearly distinguishes it from the SM one since not knowing the mass of the Higgs boson in the SM one has several indirect constraints leading to the lower boundary of $m_{h_{\text{SM}}} \geq 135 \text{ GeV}$ \cite{25}. However, after including the radiative corrections, the mass of the lightest Higgs boson in the MSSM, $m_h$, increases.

These radiative corrections vanish when supersymmetry is not broken and depend on the values of the soft breaking parameters. The main contribution comes from top (stop) quarks. Contributions from the other particles are much smaller \cite{26}. In the one loop order one has the following modification of the tree-level relation for the lightest Higgs mass

$$m_h^2 \approx M_Z^2\cos^22\beta + \frac{3g^2m_t^4}{16\pi^2M_W^2} \log \frac{\tilde{m}_t^2}{\tilde{m}_t^2}, \quad (16)$$

One finds that the one-loop correction is positive and increases the mass value. Two loop corrections have the opposite effect but are smaller and result in slightly lower value of the Higgs mass \cite{27}. To find out numerical values of these corrections, one has to determine the masses of all superpartners.

Within the Constrained MSSM, imposing various constraints, one can define the allowed region in the parameter space and calculate the spectrum of superpartners and, hence, the radiative corrections to the Higgs boson mass. The Higgs mass depends mainly on the following parameters: the top mass, the squark masses, the mixing in the stop sector, the pseudoscalar Higgs mass and $\tan\beta$. The maximum of the Higgs mass is obtained for large $\tan\beta$, for the maximal value of the top and squark masses and the minimal value of the stop mixing.

We present the value of the lightest Higgs mass in the whole $m_0, m_{1/2}$ plane for the high $\tan\beta$ solutions in Fig.10 \cite{28}. One can see that it is practically constant in the whole plane and is saturated for high values of $m_0$ and $m_{1/2}$.

The lightest Higgs boson mass $m_h$ is shown as a function of $\tan\beta$ in Fig.11 \cite{28}. The shaded band corresponds to the uncertainty from the stop mass and stop
$\tan \beta = 35$, $\mu > 0$

Figure 10: The value of the Higgs mass $m_0, m_{1/2}$ plane for the high tan $\beta$ solution $\tan \beta = 35$.

Figure 11: The mass of the lightest Higgs boson as a function of tan $\beta$ mixing for $m_t = 175$ GeV. The upper and lower lines correspond to $m_t = 170$ and 180 GeV, respectively.

The parameters used for the calculation of the upper limit are: $m_t = 180$ GeV, $A_0 = -3m_0$ and $m_0 = m_{1/2} = 1000$ GeV. The lowest line of the same figure gives the minimal values of $m_h$. For high tan $\beta$ the values of $m_h$ range from 105 GeV 125 GeV. At present, there is no preference for any of the values in this range but it can be seen that the 95% C.L. lower limit on the Higgs mass [29] of 113.3 GeV excludes $\tan \beta < 3.3$.

So combining all the uncertainties discussed before the results for the Higgs mass in the CMSSM can be summarized as follows:

- The low tan $\beta$ scenario ($\tan \beta < 3.3$) of the CMSSM is excluded by the lower
limit on the Higgs mass of 113.3 GeV.

- For the high tan β scenario the Higgs mass is found to be in the range from 110 to 120 GeV for $m_t = 175$ GeV. The central value is found to be [28]:

$$m_h = 115 \pm 3 \text{ (stop mass)} \pm 1.5 \text{ (stop mixing)} \pm 2 \text{ (theory)} \pm 5 \text{ (top mass)} \text{ GeV},$$

where the errors are the estimated standard deviations. This prediction is independent of tan β for tan β > 20 and decreases for lower tan β.

However, these SUSY limits on the Higgs mass may not be so restricting if non-minimal SUSY models are considered. In a SUSY model extended by a singlet, the so-called Next-to-Minimal model, eq.(15) is modified and at the tree level the upper bound looks like [30]

$$m_h^2 \simeq M_Z^2 \cos^2 2\beta + \lambda^2 v^2 \sin^2 2\beta,$$

(17)

where λ is an additional singlet Yukawa coupling. This coupling being unknown brings us back to the SM situation, though its influence is reduced by sin 2β. As a result, for low tan β the upper bound on the Higgs mass is slightly modified (see Fig[12]).

Even more dramatic changes are possible in models containing non-standard fields at intermediate scales. These fields appear in scenarios with gauge mediated supersymmetry breaking. In this case, the upper bound on the Higgs mass may increase up to 155 GeV [30] (the upper curve in Fig[12]), though it is not necessarily saturated. One should notice, however, that these more sophisticated models do not change the generic feature of SUSY theories, the presence of the light Higgs boson.

![Figure 12: Dependence of the upper bound on the lightest Higgs boson mass on tan β in MSSM (lower curve), NMSSM (middle curve) and extended SSM (upper curve)](image-url)
8 Observation of the Higgs boson at LHC

In principle, LHC will be able to cover the whole interval of SUSY and Higgs masses up to a few TeV. However, due to severe background, specially for the Higgs mass around the $Z$-boson mass one needs large integrated luminosity. From the point of view of observation there is no much difference between the SM Higgs boson and the lightest Higgs of the MSSM. The main production processes at hadron colliders are shown in Fig. 13. The cross section and the role of different channels depend on the mass of the Higgs boson as shown in Fig. 14.

Figure 13: The main Higgs boson production processes at hadron colliders

Figure 14: The cross sections of various Higgs boson production processes at LHC

Being created the Higgs boson will decay. The signatures of the Higgs boson are related to the dominant decay modes which again depend on the mass of the Higgs
boson. The main decay channels are listed below

- \( H, h \to \gamma\gamma, b\bar{b} \) \((H \to b\bar{b} \text{ in } WH, t\bar{t}H)\)
- \( h \to \gamma\gamma \text{ in } WH, t\bar{t}h \to \ell\gamma\gamma \)
- \( h, H \to ZZ^*, ZZ \to 4\ell \)
- \( h, H, A^+ \to \tau^+\tau^- \to (e/\mu)^+ + H^- + E_T^{miss} \to e^+ + \mu^- + E_T^{miss} \to H^+ + H^- + E_T^{miss} \) \(\text{inclusively in } b\bar{b}H_{SUSSY}\)
- \( H^+ \to \tau^+\nu \text{ from } t\bar{t} \)
- \( H^+ \to \tau^+\nu \text{ and } H^+ \to t\bar{b} \text{ for } M_H > M_{top} \)
- \( A \to Zh \) with \( h^- > b\bar{b} \); \( A \to \gamma\gamma \)
- \( H, A \to \tilde{\chi}_2^0\tilde{\chi}_2^0, \tilde{\chi}_1^+\tilde{\chi}_1^- \) \(\text{promising}\)
- \( H^+ \to \tilde{\chi}_2^+\tilde{\chi}_2^0 \)
- \( H \to \tau^+\tau^- \text{ in } WH, t\bar{t}H \)

The LHC will either discover the SM or the MSSM Higgs boson, or prove their absence. In terms of exclusion plots shown in Fig. 15, the LHC collider will cover the whole region of SUSY parameter space. Various decay modes allow one to probe different areas, as shown in Fig. 15, though the background will be very essential.

Figure 15: Exclusion plots for LHC hadron collider for different Higgs decay modes

9 Conclusion

The LHC hadron collider will have all the possibilities for important discoveries already in the first year of its operation (one day of LHC with the luminosity of
10^{33} \text{cm}^{-2}\text{s}^{-1} \text{ is equivalent to 10 years of work of the previous accelerator). Supersymmetry, if the scenarios described above are realized, might be discovered almost immediately. Slightly more complicated is the situation with the Higgs boson. Therefore the stable functioning of the accelerator with high luminosity is crucial. However, to get the desired result one needs enormous efforts on data processing and calculation of the background processes within the Standard Model at the center of mass energy of 14 TeV.}

Financial support from Russian Foundation of Basic Research (grant 05–02–17603) and the grant of the President of Russian Federation for support of leading scientific schools (NSh–5362.2006.2) is kindly acknowledged. We are grateful to the organizers of the ITEP Winter school for fruitful atmosphere and pleasant conditions.

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