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THE SUSY WORLD

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As Fayet reminded us in his introductory talk, realistic models of low-energy supersymmetry have been studied for more than 15 years. At first sight, the absence of direct experimental evidence does not match such an intense theoretical effort, and puts supersymmetry on the same footing as many other extensions of the Standard Model (SM). A number of theoretical and phenomenological reasons, however, make low-energy supersymmetry particularly attractive with respect to its alternatives: the intense activity reported at this workshop is there to prove it! My attempt to summarize its highlights will be organized as follows. In Section 1, I shall review the main motivations that lead most of us to consider the ‘SUSY world’ as a plausible scenario. The simplest realization of low-energy supersymmetry, the Minimal Supersymmetric extension of the Standard Model (MSSM), will be recalled, with some comments on possible non-minimal variations. Section 2 will summarize the phenomenology of supersymmetric particle searches, including Higgs bosons, at present and future accelerators. Finally, Section 3 will review some open theoretical problems connected with spontaneous supersymmetry breaking in supergravity and superstring models, and draw some conclusions. Unavoidably, the selection of topics will depend on space limitations and on my personal view of the subject: I apologize in advance with the participants whose contributions would have deserved a better treatment.

1. The MSSM: a paradigm for low-energy SUSY

1.1. Theoretical motivations

As discussed in the talks by Fayet and Kounnas, there are many good reasons to believe that four-dimensional supersymmetry and its local version, supergravity, could be relevant in a fundamental theory of particle interactions. In particular, superstrings are the present best candidates for a consistent quantum theory unifying gravity with all the other fundamental interactions, and supersymmetry seems to play a very important role for the quantum stability of superstring solutions in flat four-dimensional spacetime. Experimental data, however, tell us that supersymmetry is broken, but strings have not yet given us any insight about the scale of supersymmetry breaking.

The only motivation for low-energy supersymmetry, i.e. supersymmetry effectively broken around the electroweak scale, comes from the naturalness or hierarchy problem of the SM, whose formulation will now be sketched. Despite its remarkable phenomenological success, it is impossible not to regard the SM as an effective low-energy theory, valid up to some energy scale $\Lambda$ at which it is replaced by some more fundamental theory. Certainly $\Lambda$ is less than the Planck scale $M_P \sim 10^{19}$ GeV, since one needs a theory of quantum gravity to describe physics at these energies. However, the study of the Higgs sector of the SM suggests that $\Lambda$ should rather be close to the Fermi scale, $G_F^{-1/2} \sim 300$ GeV. The argument goes as follows. Consistency of the SM requires the SM Higgs mass to be less than $O(1 \text{ TeV})$. If one then tries to extend the validity of the SM to energy scales $\Lambda \gg G_F^{-1/2}$, one is faced with the fact that in the SM there is...
no symmetry to justify the smallness of the Higgs mass with respect to the (physical) cut-off $\Lambda$. This is related to the existence of quadratically divergent loop corrections to the Higgs mass in the SM. Motivated by this problem, much theoretical effort has been devoted to finding descriptions of electroweak symmetry breaking, which modify the SM at scales $\Lambda \sim G_F^{-1/2}$. Here supersymmetry comes into play because of its improved ultraviolet behaviour with respect to ordinary quantum field theories, due to cancellations between bosonic and fermionic loop diagrams. If one wants to have a low-energy effective Lagrangian valid up to scales $\Lambda \gg G_F^{-1/2}$, with one or more elementary scalar fields, kept light without unnatural fine-tuning of parameters, the solution is to introduce supersymmetry, effectively broken in the vicinity of the electroweak scale. This does not yet explain why the scale $M_{\text{SUSY}}$ of supersymmetry breaking is much smaller than $\Lambda$ (further considerations on this problem will be made in the final section), but at least links the Fermi scale $G_F^{-1/2}$ to the supersymmetry-breaking scale $M_{\text{SUSY}}$, and makes the hierarchy $G_F^{-1/2} \sim M_{\text{SUSY}} << \Lambda$ stable against radiative corrections.

1.2. The MSSM

The most economical realization of low-energy supersymmetry is the MSSM, whose defining assumptions were recalled by Fayet\(^1\). The gauge group is $G = SU(3)_C \times SU(2)_L \times U(1)_Y$, and the matter content corresponds to three generations of quarks and leptons, as in the SM. To give masses to all charged fermions and to avoid chiral anomalies, however, one is forced to introduce two complex Higgs doublets, one more than in the SM case. To enforce baryon and lepton number conservation in renormalizable interactions, one imposes a discrete $R$-parity: in practice, $R = +1$ for all ordinary particles (quarks, leptons, gauge and Higgs bosons), $R = -1$ for their superpartners (spin-0 squarks and sleptons, spin-1/2 gauginos and higgsinos). A globally supersymmetric Lagrangian $\mathcal{L}_{\text{SUSY}}$ is then fully determined by the superpotential (in standard notation)

$$f = h^U QU^c H_2 + h^D QD^c H_1 + h^E LE^c H_1 + \mu H_1 H_2.$$ (1)

To proceed towards a realistic model, one has to introduce supersymmetry breaking. In the absence of a fully satisfactory mechanism for spontaneous supersymmetry breaking at a fundamental level, it seems sensible to parametrize supersymmetry breaking at low energy by a collection of soft terms, $\mathcal{L}_{\text{soft}}$, which preserve the good ultraviolet properties of global supersymmetry. This $\mathcal{L}_{\text{soft}}$ contains mass terms for scalar fields and gauginos, as well as a restricted set of scalar interaction terms

$$-\mathcal{L}_{\text{soft}} = \sum_i \tilde{m}_i^2 |\varphi_i|^2 + \frac{1}{2} \sum_A M_A \bar{\lambda}_A \lambda_A + \left( h^U A^U QU^c H_2 + h^D A^D QD^c H_1 + h^E A^E LE^c H_1 + m_3^2 H_1 H_2 + \text{h.c.} \right),$$ (2)

where $\varphi_i$ ($i = H_1, H_2, Q, U^c, D^c, L, E^c$) denotes the generic spin-0 field, and $\lambda_A$ ($A = 1, 2, 3$) the generic gaugino field. Observe that, since $A^U, A^D$ and $A^E$ are matrices in generation space, $\mathcal{L}_{\text{soft}}$ contains in principle a huge number of free parameters. Moreover, for generic values of these parameters one encounters phenomenological problems
with flavour-changing neutral currents, new sources of CP violation, and charge- and colour-breaking vacua. All the above problems can be solved at once if one assumes that the running mass parameters in $\mathcal{L}_{\text{soft}}$, defined at the one-loop level and in a mass-independent renormalization scheme, can be parametrized, at some grand-unification scale $M_U$, by a universal gaugino mass $m_{1/2}$, a universal scalar mass $m_0$, and a universal trilinear scalar coupling $A$, whereas $m_2^3$ remains in general an independent parameter.

1.3. Non-minimal models

The above assumptions, which define the MSSM, are plausible but not compulsory: relaxing them leads to non-minimal supersymmetric extensions of the SM.

For example, as discussed in the talks by Dreiner and Kobayashi, one can consider models in which $R$-parity is explicitly broken by some superpotential couplings. The acceptable ones have either the baryon or the lepton number violated by renormalizable interactions among light particles, and give rise to phenomenological signatures that can be drastically different from the ones of the MSSM. A proof of this is the fact that, with some luck, one could be able to detect signals of supersymmetry even at HERA, which in the case of the MSSM cannot add much to what we already know.

Another possibility is to enlarge the Higgs sector of the model, for example by adding a gauge-singlet Higgs superfield, as discussed in the talk by Kane. In this case the restrictions imposed by supersymmetry on Higgs masses and couplings are much weaker than in the minimal case. On the other hand, requiring perturbative unification of couplings can still lead to interesting constraints, and in particular to an upper bound on the lightest Higgs mass of the order of 150 GeV.

As for the boundary conditions at the unification scale, one can observe with Ibáñez that the simplest unification conditions on the gauge coupling constants are not really compulsory in string unification. In a general four-dimensional string model with gauge group $SU(3) \times SU(2) \times U(1) \times G$, one can have tree-level relations such as $g_1 k_1 = g_2 k_2 = g_3 k_3$, where the integer numbers $k_a$ ($a = 1, 2, 3$) are the so-called Kac-Moody levels. In string unification there is no fundamental reason for the tree-level prediction $\sin^2 \theta_W \equiv (3/5)g_1^2/[(3/5)g_1^2 + g_2^2] = 3/8$, which is so successful when combined with the MSSM quantum corrections. Such an occurrence could be related to the existence of a gauge $U(1)_X$ symmetry of the Peccei-Quinn type, whose anomaly is cancelled by a Green-Schwarz mechanism, but no realistic string model with these properties has been found yet.

Less radically, one can also consider the possibility of non-universal boundary conditions on the soft supersymmetry-breaking terms. Such a possibility, which could be realized in string model-building (an example was given in the talk by Antoniadis, others were recently discussed in Ref. 8), is strongly constrained by the phenomenological limits on flavour-changing neutral currents, but would lead to modified relations among the low-energy parameters with respect to the MSSM case.

All these non-minimal extensions remind us that we should not take the MSSM as the only viable paradigm for low-energy supersymmetry, and that experimental analy-
ses should rather rely on the smallest possible amount of theoretical assumptions. On the other hand, non-minimal models typically increase the number of free parameters without correspondingly increasing the physical motivation, so we shall not discuss them further.

1.4. Phenomenological virtues of the MSSM

It is perhaps useful, at this point in the discussion, to recall some phenomenological virtues of the MSSM (besides the solution of the ‘technical’ part of the hierarchy problem), which were mentioned at this workshop.

As stressed in the talk by Haber, an aspect that became particularly relevant after the recent precision measurements at LEP is the fact that electroweak data put little indirect constraints, via radiative corrections, on the MSSM parameters. In most of its parameter space, the MSSM predictions for electroweak observables coincide in practice with those of the SM for a relatively light Higgs. Deviations comparable to the present experimental accuracy can only occur for a light stop-bottom sector with large mass splittings, or for a chargino with mass just above the production threshold at LEP I. This is not the case, for example, of technicolor and extended technicolor models, which are severely constrained by the recent LEP data.

Another important property of the MSSM, discussed in the talks by Haber and Wagner, is related to the fact that the running top Yukawa coupling \( h_t(Q) \) has an effective infrared fixed point, smaller than in the SM case. Neglecting mixing and the Yukawa couplings of the remaining fermions, \( h_t \) obeys the following one-loop renormalization group equation (RGE)

\[
\frac{dh_t}{dt} = \frac{h_t}{8\pi^2} \left( -\frac{8}{3} g_3^2 - \frac{3}{2} g_2^2 - \frac{13}{18} g'^2 + 3h_t^2 \right) , \quad (t \equiv \log Q) .
\]

A close look at Eq. (3) can give us some important information about the acceptable values for the top-quark mass, \( m_t \), and the ratio of vacuum expectation values, \( \tan \beta \equiv v_2/v_1 \). However high is the value assigned to it at the unification scale, \( h_t \) evaluated at the electroweak scale never exceeds a certain maximum value \( h_t^{\text{max}} \approx 1 \). This implies that, for any given value of \( \tan \beta \), there is a corresponding maximum value for the top-quark mass. A naïve one-loop calculation gives \( m_t^{\text{max}} \sim 200 \text{ GeV} \cdot \sin \beta \), which naturally puts the top-quark mass in the range presently allowed by direct searches and electroweak precision data. The results of a more refined calculation, which includes the effects of all third-generation Yukawa couplings and of the supersymmetric threshold \( M_{\text{SUSY}} \), are shown in Fig. 1.

In the MSSM, \( R \)-parity makes the lightest \( R \)-odd supersymmetric particle (LSP) absolutely stable. In most of the otherwise acceptable parameter space, the LSP is neutral and weakly interacting, rarely a sneutrino, and typically the lightest, \( \tilde{\chi}_1 \), of the neutralinos (the mass eigenstates of the neutral gaugino-higgsino sector). Then the LSP is a natural candidate for cold dark matter, as discussed by Roszkowski, who reported calculations of the neutralino relic density in different regions of the MSSM.
Fig. 1. The region of the \((\tan \beta, m_t)\) parameter space in which all running Yukawa couplings remain finite at energy scales up to \(\Lambda = 10^{16} \text{ GeV}\) (from Ref. 9).

parameter space. His results could be summarized as follows. In most of the otherwise acceptable parameter space, the LSP is cosmologically harmless, in the sense that its relic density is smaller than the closure density of the Universe. Moreover, in a small but non-negligible region of parameter space, the LSP relic density turns out to be large enough to be of cosmological interest in relation with the dark-matter problem. It is then conceivable, even if not very likely, that the first evidence for supersymmetry could come from the dedicated experiments searching for dark-matter signals!

One of the most attractive features of the MSSM is the possibility of describing the spontaneous breaking of the electroweak gauge symmetry as an effect of radiative corrections, as discussed by Kounnas and Lahanas. It is remarkable that, starting from universal boundary conditions at the unification scale, it is possible to explain naturally why fields carrying colour or electric charge do not acquire non-vanishing VEVs, whereas the neutral components of the Higgs doublets do. We give here a simplified description of the mechanism in which the physical content is transparent. The starting point is a set of boundary conditions on the independent model parameters at the unification scale \(Q = M_U\). One then evolves all the running parameters from the grand-unification scale to a low scale \(Q \sim G_F^{-1/2}\), according to the RGEs, and considers
the renormalization-group-improved tree-level potential

\[
V_0(Q) = m_1^2 |H_1|^2 + m_2^2 |H_2|^2 + m_3^2 (H_1 H_2 + \text{h.c.})
+ \frac{g'^2}{8} \left( H_2 \bar{\sigma} H_2 + H_1 \bar{\sigma} H_1 \right)^2
+ \frac{g'^2}{8} \left( |H_2|^2 - |H_1|^2 \right)^2.
\]

(4)

All masses and coupling constants in \(V_0(Q)\) are running parameters, evaluated at the scale \(Q\). The minimization of the potential in Eq. (4) is straightforward. To generate non-vanishing VEVs \(v_1 \equiv \langle H_1^0 \rangle\) and \(v_2 \equiv \langle H_2^0 \rangle\), one needs

\[
\mathcal{B} \equiv m_1^2 m_2^2 - m_3^4 < 0.
\]

(5)

In addition, a certain number of conditions have to be satisfied to have a stable minimum with the correct amount of symmetry breaking and with unbroken colour, electric charge, baryon and lepton number: for example, all the running squark and slepton masses have to be positive. In the whole process, a crucial role is played by the top Yukawa coupling, which strongly influences the RGE for \(m_2^2\). For appropriate boundary conditions, the RGEs drive \(\mathcal{B} < 0\) at scales \(Q \sim G_F^{-1/2}\), whereas all the squark and slepton masses remain positive as desired, to give a phenomenologically acceptable breaking of the electroweak symmetry.

1.5. Supersymmetric grand-unification

The previous list has left out one of the most impressive arguments in favour of low-energy supersymmetry, i.e. the agreement of the generic predictions of supersymmetric grand unification with the extracted values of the gauge coupling constants at the electroweak scale. This topic has been discussed at great length by Zichichi at this workshop, and I will try to give here my personal summary of the subject.

Starting from the boundary condition

\[
g_3(M_U) = g_2(M_U) = g_1(M_U) \equiv g_U,
\]

(6)

where \(g_1 = \sqrt{5/3} \cdot g'\) as in most grand-unified models, one can solve the appropriate RGEs to obtain the running gauge coupling constants \(g_A(Q)\) \((A = 1, 2, 3)\) at scales \(Q << M_U\). At the one-loop level, and assuming that there are no new physics thresholds between \(M_U\) and \(Q\), one finds

\[
\frac{1}{g_A^2(Q)} = \frac{1}{g_U^2} + \frac{b_A}{8\pi^2} \log \frac{M_U}{Q} \quad (A = 1, 2, 3),
\]

(7)

where the one-loop coefficients \(b_A\) depend only on the \(SU(3)_C \times SU(2)_L \times U(1)_Y\) quantum numbers of the light particle spectrum. In the MSSM

\[
b_3 = -3, \quad b_2 = 1, \quad b_1 = \frac{33}{5},
\]

(8)
whereas in the SM
\[ b_0^3 = -7, \quad b_0^2 = -\frac{19}{6}, \quad b_0^1 = \frac{41}{10}. \]  
(9)
Starting from three input quantities at the electroweak scale, for example \( \alpha_3(m_Z) \), \( \alpha_{em}^{-1}(m_Z) \) and \( \sin^2 \theta_W(m_Z) \), one can perform consistency checks of the grand-unification hypothesis in different models.

In the minimal \( SU(5) \) model and indeed in any other model where Eq. (6) holds and the light-particle content is just that of the SM (with no intermediate mass scales between \( m_Z \) and \( M_U \) ), Eqs. (7) and (9) are incompatible with experimental data. This was first realized by noticing that the prediction \( M_U \approx 10^{14-15} \text{ GeV} \) is incompatible with the experimental data on nucleon decay. Subsequently, also the prediction \( \sin^2 \theta_W \approx 0.21 \) was shown to be in conflict with the experimental data, and this conflict became even more significant after the recent LEP precision measurements.

In the MSSM, assuming for simplicity that all supersymmetric particles have masses of order \( m_Z \), one obtains \( M_U \approx 10^{16} \text{ GeV} \) (which increases the proton lifetime for gauge-boson-mediated processes beyond the present experimental limits) and \( \sin^2 \theta_W \approx 0.23 \). At the time of Ref. 16, when data were pointing towards a significantly smaller value of \( \sin^2 \theta_W \), this was considered by some a potential phenomenological shortcoming of the MSSM. The high degree of compatibility between data and supersymmetric grand unification became manifest only later, after improved data on neutrino-nucleon deep inelastic scattering were obtained; it was recently re-emphasized after the LEP precision measurements. One should not forget, however, that unification of the MSSM is not the only solution that can fit the present extracted values of the gauge coupling constants at \( Q = m_Z \); for example, non-supersymmetric models with \textit{ad hoc} light exotic particles or intermediate symmetry-breaking scales could also do the job. The MSSM, however, stands out as the simplest physically motivated solution.

If one wants to make the comparison between low-energy data and the predictions of specific grand-unified models more precise, there are several factors that should be further taken into account. After the inclusion of higher-loop corrections and threshold effects, Eq. (7) is (schematically) modified as follows

\[ g_A^2(Q) = g_U^2 + \frac{b_A}{8\pi^2} \log \frac{M_U}{Q} + \Delta_A^{th} + \Delta_A^{l>1} \quad (A = 1, 2, 3). \]  
(10)
In Eq. (10), \( \Delta_A^{th} \) represents the so-called \textit{threshold effects}, which arise whenever the RGEs are integrated across a particle threshold, and \( \Delta_A^{l>1} \) represents the corrections due to two- and higher-loop contributions to the RGEs. Both \( \Delta_A^{th} \) and \( \Delta_A^{l>1} \) are scheme-dependent, so that one should be careful to compare data and predictions within the same renormalization scheme. The \( \Delta_A^{th} \) receives contributions both from thresholds around the electroweak scale (top quark, Higgs boson, and in SUSY-GUTs also the additional particles of the MSSM spectrum) and from thresholds around the grand-unification scale (superheavy gauge and Higgs bosons, and in SUSY-GUTs also their
superpartners). Needless to say, these last threshold effects can be computed only in the framework of a specific grand-unified model, and typically depend on a number of free parameters. Besides the effects of gauge couplings, $\Delta_{T}^{l=1}$ must include also the effects of Yukawa couplings, since, even in the simplest mass-independent renormalization schemes, gauge and Yukawa couplings mix beyond the one-loop order. In minimal $SU(5)$ grand unification, and for sensible values of the top and Higgs masses, all these corrections are small and do not substantially affect the conclusions derived from the naïve one-loop analysis. This is no longer the case, however, for supersymmetric grand unification. First of all, one should notice that the MSSM by itself does not uniquely define a SUSY-GUT, whereas threshold effects and even the proton lifetime (owing to a new class of diagrams, which can be originated in SUSY-GUTs) become strongly model-dependent. Furthermore, the simplest SUSY-GUT, containing only chiral Higgs superfields in the 24, 5 and $\bar{5}$ representations of $SU(5)$, has a severe problem in accounting for the huge mass splitting between the $SU(2)$ doublets and the $SU(3)$ triplets sitting together in the 5 and $\bar{5}$ Higgs supermultiplets, and cannot reproduce correctly the observed pattern of fermion masses and mixing angles. Threshold effects are typically larger than in ordinary GUTs, because of the much larger number of particles in the spectrum, and in any given model they depend on several unknown parameters. Also two-loop effects of Yukawa couplings can be quantitatively important in SUSY-GUTs, since they depend not only on the top-quark mass, but also on the ratio $\tan \beta = \frac{v_2}{v_1}$ of the VEVs of the two neutral Higgs fields: these effects become large for $m_t > 140$ GeV and $\tan \beta \sim 1$, which correspond to a strongly interacting top Yukawa coupling. All these effects, and others, have been recently re-evaluated. The conclusion is that, even imagining a further reduction in the experimental errors on the three gauge couplings, it is impossible to claim indirect evidence for supersymmetry and to predict the MSSM spectrum with any significant accuracy. The only safe statement is that, at the level of precision corresponding to the naïve one-loop approximation, there is a remarkable consistency between experimental data and the prediction of supersymmetric grand unification, with the MSSM $R$-odd particles roughly at the electroweak scale. These conclusions are summarized in Fig. 2, borrowed from Ref. 22, which compares post-LEP and pre-LEP uncertainties, both theoretical and experimental, in the determination of $\sin^2 \theta_W (m_Z)$.

At this point it is worth mentioning how the unification constraints can be applied to the low-energy effective theories of four-dimensional heterotic string models. The basic fact to be realized is that the only free parameter of these models is the string tension, which fixes the unit of measure of the massive string excitations. All the other scales and parameters are related to VEVs of scalar fields, the so-called moduli, corresponding to flat directions of the scalar potential. In particular, there is a relation between the string mass $M_S \sim \alpha'^{-1/2}$, the Planck mass $M_P \sim G_N^{-1/2}$, and the unified coupling constant $g_U$, which reflects unification with gravity and implies that in any given string vacuum one has one more prediction than in ordinary field-theoretical grand unification. In a large class of string models, one can write down an equation
Fig. 2. Comparison of theory and experiment in the determination of the electroweak mixing angle from the unification hypothesis, now and before LEP (from Ref. 22).
of the same form as (10), and compute \( g_U \), \( M_U \), \( \Delta_A \), \( \ldots \) in terms of the relevant VEVs. In the \( \overline{DR} \) scheme one finds \( M_U \simeq 0.5 \times g_U \times 10^{18} \text{ GeV} \), more than one order of magnitude higher than the naive extrapolations from low-energy data illustrated before. This means that significant threshold effects are needed in order to reconcile string unification with low-energy data: for example, the minimal version of the flipped-\( SU(5) \) model is by now ruled out. To get agreement, one needs some more structure in the spectrum, either at the compactification scale or in the form of light exotics, but even in this case one suffers a loss of predictivity. However, unification constraints now stand as a very important phenomenological test for any realistic string model.

To conclude the discussion of supersymmetric grand unification, a few more words on proton decay seem appropriate. We have already stressed that the RGE of the MSSM imply a unification scale \( M_U \simeq 10^{16} \text{ GeV} \). This means that proton decay mediated by heavy vector bosons, which favours decay modes such as \( p \to e^+ \pi^0 \) and whose rate is proportional to \( M^{-4}_U \), is suppressed to unobservable levels. On the other hand, an entirely new possibility arises in SUSY-GUTs. With the superfield content of the MSSM, one can construct supersymmetric gauge-invariant operators with \( \Delta B = \Delta L = \pm 1 \) and mass dimension \( d = 5 \). These operators can be generated, for example, by the exchange of some heavy Higgs superfields of minimal supersymmetric \( SU(5) \). The nucleon decay amplitudes are obtained by dressing these operators by SUSY particle exchanges, to convert sfermion lines into light fermion lines. The resulting decay rate scales now as \( M^{-2}_U M^{-2}_{\text{SUSY}} \), with the actual numerical value depending both on the details on the SUSY-GUT model and on the details of the low-energy particle spectrum. In view of the first class of ambiguities, to my mind it is not particularly interesting to take minimal SUSY \( SU(5) \) and to look for constraints on the soft breaking terms from the limits on proton decay. On the other hand, an important generic feature emerges from the symmetry structure of the dimension-five operators: the dominant nucleon decay modes should involve strange particles in the final state, as in \( p \to K^+ \pi^- \). Detection of nucleon decay in one of these channels would certainly be a very strong argument in favour of supersymmetric grand unification.

2. Supersymmetry searches

2.1. The particle spectrum of the MSSM

In the \( R \)-even sector, the only new feature of the MSSM with respect to the SM is its extended Higgs sector, with two independent VEVs, \( v_1 \equiv \langle H^0_1 \rangle \) and \( v_2 \equiv \langle H^0_2 \rangle \), which can be taken to be real and positive without loss of generality. Quarks of charge 2/3 have masses proportional to \( v_2 \), quarks of charge 1/3 and charged leptons have masses proportional to \( v_1 \). The W and Z masses are proportional to \( \sqrt{v_1^2 + v_2^2} \), which is therefore fixed by their measured values. The remaining freedom is conveniently parametrized by \( \tan \beta \equiv v_2/v_1 \), whose allowed range of variation in the MSSM is \( 1 \lesssim \tan \beta \lesssim m_t/m_b \). The physical states of the MSSM Higgs sector are three neutral bosons (two CP-even, \( h \) and \( H \), and one CP-odd, \( A \)) and a charged boson, \( H^\pm \). It is important to realize that, at the tree level, all Higgs masses and couplings can be expressed in terms of two
parameters only: a convenient choice is, for example, \((m_A, \tan \beta)\). Radiative corrections to Higgs masses and couplings, however, can be large, as we shall review later, and have to be taken into account in phenomenological analyses.

In the \(R\)-odd sector of the MSSM, the spin-0 fields are the squarks and the sleptons. Neglecting intergenerational mixing, and leaving aside the stop squarks for the moment, their masses can be easily calculated in terms of the fundamental parameters \(m_0, m_{1/2}\) and \(\tan \beta\):

\[
m_f^2 = m_f^2 + \tilde{m}_f^2 + m_D^2(f),
\]

\[
\tilde{m}_f^2 = m_0^2 + C(\tilde{f}) m_{1/2}^2,
\]

\[
m_D^2(\tilde{f}) = m_Z^2 \cos 2\beta (T_{f3}^L - \sin^2 \theta_W Q_f),
\]

where, omitting generation indices, \(f = [q \equiv (u, u^c, d, d^c), \ell \equiv (\nu, e, e^c)]\) and \(C(\tilde{q}) \sim 5 - 8, C(\tilde{l}) \simeq 0.5, C(\tilde{e}^c) \simeq 0.15\).

Among the spin-1/2 particles one finds the strongly interacting gluinos, \(\tilde{g}\), whose mass is directly related to the \(SU(2)\) and \(U(1)\) gaugino masses and to \(m_{1/2}\) by

\[
\frac{m_{\tilde{g}}}{\alpha_S} \approx \frac{M_2}{\alpha_2} \approx \frac{M_1}{\alpha_1} \approx \frac{m_{1/2}}{\alpha_U},
\]

where \(\alpha_1 \equiv (5/3)g'^2/(4\pi)\) and \(\alpha_U\) is the gauge coupling strength at the grand unification scale. The weakly interacting spin-1/2 particles are the \(SU(2) \times U(1)\) gauginos \((\tilde{W}^\pm, \tilde{W}^0, \tilde{B})\) and the higgsinos \((\tilde{H}^\pm, \tilde{H}_1^0, \tilde{H}_2^0)\). These interaction eigenstates mix non-trivially via their mass matrices: the two charged mass eigenstates, called charginos, are denoted by \(\tilde{\chi}^\pm_i\) \((i = 1, 2)\), and the four neutral mass eigenstates, called neutralinos, by \(\tilde{\chi}_k^0\) \((k = 1, 2, 3, 4)\). All masses and couplings in the chargino-neutralino sector can be described in terms of the three parameters \(m_{1/2}\) [which determines the \(SU(2) \times U(1)\) gaugino masses via eq. (14)], \(\mu\) (the supersymmetric Higgs-Higgsino mass term) and \(\tan \beta\). It should be noted that \(\tilde{\chi}_1^0\), often denoted simply as \(\tilde{\chi}\), is the favourite candidate for being the LSP. An alternative candidate is \(\tilde{\nu}_\tau\), but this is actually the LSP for a much smaller range of parameter space. Notice also that there is no particular reason to assume that \(\tilde{\chi}\) is a pure photino, as is often done in phenomenological analyses.

In summary, the particle spectrum of the MSSM can be approximately described in terms of five basic parameters:

- The mass \(m_A\) of the CP-odd neutral Higgs boson (or any other SUSY Higgs mass)
- The ratio of VEVs \(\tan \beta \equiv v_2/v_1\)
- The universal gaugino mass \(m_{1/2}\), or equivalently the gluino mass \(m_{\tilde{g}}\)
- The universal scalar mass \(m_0\)
- The supersymmetric Higgs-Higgsino mass \(\mu\)
Of course, the top-quark mass is undetermined, as in the SM. Also, as we shall see later, more subtleties have to be introduced for the description of the stop system.

2.2. Searches for SUSY Higgs bosons

We have already mentioned the fact that, at the classical level, the Higgs sector of the MSSM is very tightly constrained. However, as extensively discussed at this workshop [20-22], Higgs boson masses and couplings are subject to large, finite radiative corrections, dominated by loops involving the top quark and its supersymmetric partners.

To illustrate the case with a simple example, we can assume a universal soft SUSY-breaking squark mass, $m_{\tilde{q}}$, and negligible mixing in the stop mass matrix. The leading correction to the neutral CP-even mass matrix is then

$$\left( \Delta M_R^2 \right)_{22} = \frac{3}{8\pi^2} \frac{g^2 m_t^4}{m_W^2 \sin^2 \beta} \log \left( 1 + \frac{m_{\tilde{q}}^2}{m_t^2} \right). \tag{15}$$

The most striking fact in Eq. (15) is that the correction $(\Delta M_R^2)_{22}$ is proportional to $(m_t^4/m_W^2)$ for fixed $(m_{\tilde{q}}/m_t)$. This implies that, for $m_t$ in the presently allowed range, the tree-level predictions for $m_h$ and $m_H$ can be badly violated, as for the related inequalities. The other free parameter in Eq. (15) is $m_{\tilde{q}}$, but the dependence on it is much milder. The result of Eq. (15) has been generalized to arbitrary values of the parameters in the stop mass matrix, and the effects of other virtual particles in the loops have been included. Renormalization-group methods have been used to resum the large logarithms that arise when the typical scale of supersymmetric particle masses, $M_{\text{SUSY}}$, is much larger than $m_Z$. Two-loop corrections have been computed in the leading logarithmic approximation, and found to be small. After all these refinements, Eq. (15) still gives the most important mass correction in the most plausible region of parameter space.

The computation of radiative corrections has been extended to the other parameters of the MSSM Higgs sector. One-loop corrections to the charged Higgs mass have been computed and found to be small, at most a few GeV, for generic values of the parameters. In the case of Higgs boson self-couplings, which control decays such as $H \rightarrow hh$, $H \rightarrow AA$ and $h \rightarrow AA$ when they are kinematically allowed, radiative corrections can be numerically large, being formally proportional to $(m_t/m_W)^4$. Higgs couplings to vector boson and fermions can be more easily handled: in most phenomenological studies, radiative corrections to these couplings need to be taken into account only approximately, by improving the tree-level formulae with one-loop-corrected values of the $H-h$ mixing angle, $\alpha$, and with running fermion masses, evaluated at the typical scale $Q$ of the process under consideration. Residual corrections have been computed and found to be numerically small in the experimentally interesting regions of parameter space.

We now move to the discussion of SUSY Higgs searches at present and future accelerators. For definiteness, when making numerical examples we shall make the
same assumptions as for Eq. (15) and choose the numerical values $m_t = 140 \text{ GeV}, m_{\tilde{q}} = 1 \text{ TeV}$: for this parameter choice, the maximum value of $m_h$, reached for $m_A \gg m_Z$ and $\tan \beta \gg 1$, is approximately 110 GeV, $O(20 \text{ GeV})$ larger than the tree-level upper bound. For given $m_A$ and $\tan \beta$, the shift in $m_h$ can be as large as $O(50 \text{ GeV})$, when $\tan \beta \sim 1$. In particular, after radiative corrections one can have not only $m_h > m_Z$, but also $m_h > m_A$.

As discussed by Clare\cite{Clare} and Fisher\cite{Fisher}, the relevant processes for MSSM Higgs boson searches at LEP I are $Z \to hZ^*$ and $Z \to hA$, which play a complementary role, since their rates are proportional to $\sin^2(\beta - \alpha)$ and $\cos^2(\beta - \alpha)$, respectively. An important effect of radiative corrections is to render possible, for some values of the parameters, the decay $h \to AA$, which would be kinematically forbidden according to tree-level formulae. Experimental limits that take radiative corrections into account have by now been obtained by the four LEP collaborations, using different methods to present and analyse the data, and different ranges of parameters in the evaluation of radiative corrections. An example is given in Fig. 3, where the cross-hatched area corresponds to the presently excluded region for the parameter choice $m_t = 140 \text{ GeV}, m_{\tilde{q}} = 1 \text{ TeV}$.

The situation in which the impact of radiative corrections is most dramatic is the search for MSSM Higgs bosons at LEP II, as discussed at this workshop by Katsanevakis\cite{Katsanevakis}. At the time when only tree-level formulae were available, there was hope that LEP could completely test the MSSM Higgs sector. According to tree-level for-
mulae, in fact, there should always be a CP-even Higgs boson with mass smaller than \( m_Z \) (\( h \)) or very close to it (\( H \)), and significantly coupled to the \( Z \) boson. However, this result can be completely upset by radiative corrections. A detailed evaluation of the LEP II discovery potential can be made only if crucial theoretical parameters (such as the top-quark mass and the various soft supersymmetry-breaking masses) and experimental parameters (such as the centre-of-mass energy, the luminosity, and the \( b \)-tagging efficiency) are specified. Taking for example \( \sqrt{s} = 190 \) GeV, and the parameter choice of Fig. 3, there is a region of parameter space where the associated production of a \( Z \) and a CP-even Higgs can be pushed beyond the kinematical limit. Associated \( hA \) production could be a useful complementary signal, but obviously only for \( m_h + m_A < \sqrt{s} \). Associated \( HA \) production is typically negligible at these energies. The hatched area of Fig. 3 shows the domain accessible to LEP II for the mentioned parameter choice, for an integrated luminosity of \( 500 \text{ pb}^{-1} \) and for a detector similar to the ALEPH detector at LEP: one can see that the theoretically allowed parameter space cannot be fully tested.

Of course, one should keep in mind that there is, at least in principle, the possibility of further extending the maximum LEP energy up to values as high as \( \sqrt{s} \approx 230–240 \) GeV, at the price of introducing more (and more performing) superconducting cavities into the LEP tunnel. More boldly, one can consider the possibility of an \( e^+e^- \) linear collider with \( \sqrt{s} \sim 500 \) GeV and a luminosity of order \( 10^{33} \text{ cm}^{-2}\text{sec}^{-1} \): a detailed study of the discovery potential of such a collider has been presented at this workshop by Grivaz. Among the relevant production mechanisms there are those already mentioned for LEP II: (a) \( e^+e^- \rightarrow hZ \), (b) \( e^+e^- \rightarrow HZ \), (c) \( e^+e^- \rightarrow hA \), (d) \( e^+e^- \rightarrow HA \); in addition, one can consider \( WW \) and \( ZZ \) fusion: (e) \( e^+e^- \rightarrow h\nu\bar{\nu} \) or \( he^+e^- \), (f) \( e^+e^- \rightarrow H\nu\bar{\nu} \) or \( He^+e^- \). Considering the domain that will remain unexplored if the centre-of-mass energy at LEP II is limited to 190 GeV, there are four main configurations (see Fig. 3): in (1) \( h \) is SM-like and accessible via (a) and (e); in (2) one has in addition the possibility of detecting (d); in (3) the observable processes are (b), (c) and (f); in (4) all processes are kinematically allowed and only moderately suppressed with respect to the SM case. Also, in regions (2), (3) and (4) one can observe pair production of charged Higgses: (g) \( e^+e^- \rightarrow H^+H^- \). In summary, such a linear \( e^+e^- \) collider would allow for a complete exploration of the MSSM parameter space: if the MSSM is indeed correct, one could expect the guaranteed detection of at least one neutral Higgs state and the concrete possibility of a detailed spectroscopy of the Higgs sector.

Another interesting possibility offered by a linear \( e^+e^- \) collider is the study of \( \gamma\gamma \) collisions at very high energy and luminosity, thanks to a back-scattered laser beam facility. The physics impact of such a machine on the SUSY-Higgs sector has been discussed by Gunion at this meeting, who emphasized its complementarity with respect to the \( e^+e^- \) mode.

The next question, which was discussed by Kunszt and also by Gunion, is whether the LHC/SSC can explore the full parameter space of the MSSM Higgs bosons.
The analysis is complicated by the fact that the $R$-odd particles could play a role both in the production (via loop diagrams) and in the decay (via loop diagrams and as final states) of the MSSM Higgs bosons. For simplicity, one can concentrate on the most conservative case in which all $R$-odd particles are heavy enough not to play any significant role. Still, one has to perform a separate analysis for each $(m_A, \tan \beta)$ point, to include radiative corrections (depending on additional parameters such as $m_t$ and $m_{\tilde{q}}$), and to consider Higgs boson decays involving other Higgs bosons.

The most promising signals at the LHC/SSC are $h, H \rightarrow \gamma\gamma$ (inclusive or in association with a $W$ boson or with a $t\bar{t}$ pair, giving an extra isolated lepton in the final state), $H \rightarrow ZZ \rightarrow 4l^\pm$, $A, H \rightarrow \tau^+\tau^-$ and $t \rightarrow bH^+ \rightarrow b\tau^+\nu_\tau$. A pictorial representation of the LHC/SSC discovery potential corresponding to the different processes is given in Fig. 4, which also shows, as dashed lines, contours associated with a ‘pessimistic’ and an ‘optimistic’ estimate of the LEP II sensitivity. In summary, a global look at Fig. 4 shows that there is a high degree of complementarity between the regions of parameter space accessible to LEP II and to the LHC/SSC. However, for our representative choice of parameters, there is a non-negligible region of the $(m_A, \tan \beta)$ plane that is presumably beyond the reach of LEP II and of the LHC/SSC. This potential problem could be solved, as we said before, by a further increase of the LEP II energy beyond the reference value of $\sqrt{s} \lesssim 190$ GeV or by a high-energy linear $e^+e^-$ collider. One should also keep in mind that indirect information on the particle spectrum of the
MSSM, including its extended Higgs sector, could come from lower-energy precision data. An interesting effect, mentioned at this workshop by Haber, is the contribution of the charged-Higgs loop to the rare decay $b \to s\gamma$, which in the SM proceeds via a $W$-boson loop. The theoretical and experimental errors on the inclusive radiative $B$-decay could already be small enough to put non-trivial constraints on the particle spectrum of the MSSM. In particular, in the limit of very heavy $R$-odd particles one could identify an excluded region in the $(m_A, \tan \beta)$ plane, corresponding to low values of $m_{H^\pm}$: the precise form of such a region strongly depends on the assumed theoretical uncertainties.

One should also mention that a complete study of the SUSY-Higgs phenomenology cannot neglect the possibility of a relatively light spectrum of $R$-odd particles: as reported by Baer and Tata, in this case one expects a worsening of the standard signals, which could be compensated, however, by the appearance of new signals related to Higgs decays into pairs of $R$-odd particles.

2.3. Searches for $R$-odd particles

The phenomenology of $R$-odd SUSY particles has been discussed in many contributions to this workshop.

The present status of experimental searches for supersymmetry is a collection of negative results, which translate into limits on the MSSM parameters. LEP experiments have searched for a variety of possible supersymmetric $Z$ decays

$$Z \rightarrow \tilde{t}^\pm \tilde{t}^\mp, \tilde{\nu} \tilde{\nu}, \tilde{\chi}^\pm \tilde{\chi}^\mp, \tilde{\chi}_1^0 \tilde{\chi}_2^0, \tilde{q} \tilde{g},$$

both directly and indirectly (via measurements of the $Z$ line shape). A crude summary is that $R$-odd particles weighing much less than $m_Z/2$ are in general excluded, with some possible exceptions, corresponding to particles with strongly suppressed couplings to the $Z$. The first exception are the lightest neutralinos $\tilde{\chi}_1^0$ and $\tilde{\chi}_2^0$: for small values of $m_{1/2}$ and $\tan \beta$, one can have the charginos and the heavier neutralinos beyond the kinematical limit of LEP I, and the two lightest neutralinos with dominant gaugino components and therefore decoupled from the $Z$ boson. The LEP lower bounds on the masses of $\tilde{\chi}_1^0$ and $\tilde{\chi}_2^0$ are of the order of 20 and 45 GeV, respectively, for large values of $\tan \beta$, but evaporate for values of $\tan \beta$ sufficiently close to unity. The second exception is the lighter stop state. In the stop mass matrix, the mixing term between $\tilde{t}_L$ and $\tilde{t}_R$ is proportional to the top-quark mass, so it cannot always be neglected. It is quite possible that the lighter stop state, $\tilde{t}_1 \equiv \cos \theta_t \tilde{t}_L + \sin \theta_t \tilde{t}_R$, is significantly lighter than the top quark and the remaining squarks. In this case, its coupling to the $Z$ boson is proportional to $[\cos^2 \theta_t - (4/3) \sin^2 \theta_W]$, so that for certain values of the mixing angle one can still have $m_{\tilde{t}_1} \simeq 25$ GeV. The third and final exception is the gluino, which does not have any tree-level coupling to the $Z$ boson: no limit on the gluino mass has been extracted from LEP data so far.

In the case of squarks and gluinos, additional information comes from the experiments at $p\overline{p}$ colliders. In this case, limits are considerably more model-dependent than
at LEP, due to the complicated pattern of cascade decays that can arise\[1,3,5\]. Figure 5 shows the squark and gluino mass limits assuming a light LSP ($m_\tilde{\chi} < 15$ GeV), six mass-degenerate squarks, and no cascade decays. However, once the MSSM squark and gluino branching ratios are introduced into the analysis, the above limits can be considerably degraded, as can be seen in the example of Fig. 6, corresponding to a representative parameter choice in the chargino-neutralino sector. One should also keep in mind that $p\bar{p}$-collider searches are not sensitive to very light gluino masses, which must be excluded by different methods. At present, it seems very difficult to exclude gluinos weighing between 3 and 4 GeV with lifetimes around $10^{-13}$ s, and gluinos weighing 3 GeV or more and having lifetimes between about $10^{-8}$ and $10^{-10}$ s.

Future accelerators should allow for great progress in the search for $R$-odd supersymmetric particles. At LEP II, one should be sensitive to pair production of sleptons and charginos almost up to the kinematical threshold\[30,31\]. At the LHC and the SSC, one should be able to search for squarks and gluinos up to masses of the order of 1 TeV or more\[34,38\]. This should allow the coverage of most of the theoretically plausible region of parameter space.

3. Theoretical outlook and conclusions

Among the open theoretical problems, the most important and challenging one is to understand the mechanism of spontaneous supersymmetry breaking and the origin of the hierarchy. These problems cannot be addressed, by definition, within the MSSM, where supersymmetry is broken explicitly and the soft breaking terms are controlled by arbitrary input parameters. Present theoretical ideas\[2\] and phenomenological requirements favour the possibility that supersymmetry is spontaneously broken in the hidden
sector of some underlying supergravity (or superstring) model, communicating with the observable sector (the one containing the states of the MSSM) only via gravitational interactions. As for the precise mechanism of spontaneous breaking of local supersymmetry, there are several suggestions, among which non-perturbative phenomena (such as gaugino condensation) and string constructions (such as coordinate-dependent compactifications), but none of them has yet reached a fully satisfactory formulation. If one day a convincing mechanism will be found, by taking the low-energy limit it will be possible to predict the mass parameters of the MSSM, with enormous enhancement in predictive power.

Leaving aside this open theoretical problem, one can say that low-energy supersymmetry, incarnated in the MSSM, stands out as a theoretically motivated, phenomenologically acceptable and calculationally well-defined extension of the SM. To the eyes of many observers, its plausibility has increased over the years. On the one hand, experimental searches have excluded until now only a relatively small part of its natural parameter space. On the other hand, electroweak precision measurements nicely fit an elementary Higgs sector and supersymmetric grand unification.

The phenomenological properties of the MSSM are by now well understood. Accurate computations of cross-sections and branching ratios for the MSSM particles are available, as functions of the model parameters. A lot of simulation work has been and is being performed for present and future colliders. At the level of the MSSM, the present challenge is mainly an experimental one. LEP I, which has impressively improved the previous limits on the weakly interacting MSSM particles, has almost
saturated its discovery potential. The next big step will occur at LEP II, which should allow the search for charginos and charged sleptons up to masses of order 80–90 GeV, and significant progress in the search for supersymmetric Higgs bosons. The present and future runs of the Tevatron collider should significantly push the sensitivity to squarks and gluinos, up to masses of order 200–300 GeV. If the idea of low-energy supersymmetry is correct, a discovery is not unlikely already at this stage. A negative result, however, would not yet discourage the SUSY advocates. Only the LHC and the SSC, and perhaps a linear $e^+e^-$ collider with $\sqrt{s} \gtrsim 500$ GeV, will be able to perform a decisive test: the former should be able to probe squark and gluino masses up to 1–2 TeV, the latter should be sensitive to sleptons and charginos up to 200 GeV and more, and to SUSY Higgs bosons in any plausible model. A positive signal would open up an exciting new generation of experiments, a negative one would presumably push low-energy supersymmetry into oblivion.

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