M_W AND SUPERSYMMETRY*

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Various aspects of the link between the weak gauge boson masses and supersymmetry are discussed.

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The masses of the weak gauge bosons, $W^{\pm}$ and Z^0, originate from the Higgs mechanism. In the Standard Model, viewed as an effective low energy theory, the Higgs potential looks very unnatural and the theory faces the well known hierarchy problem. Supersymmetry offers an interesting solution to the hierarchy problem and it is here where the strong link exists between the gauge boson masses and supersymmetry.

The content of this lecture is the following:

1) The Higgs potential in the Standard Model and need for new physics.
2) Is the Higgs boson light or heavy — precision tests at the Z^0 pole in the Standard Model.
3) The hierarchy problem and supersymmetry.
4) Supersymmetric Higgs sector and its naturalness.
5) Minimal Supersymmetric Standard Model and electroweak measurements ($M_W etc$).
6) Have superpartners been already discovered?

As we shall see the considerations in Section 4 and 5 provide interesting upper and lower bounds, respectively, on the supersymmetric spectrum.

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1. The Higgs potential in the Standard Model and the need for new physics

Spontaneous breaking of the electroweak gauge symmetry SU(2)× U(1) is now confirmed experimentally with one per mille accuracy (for instance, the measured value $M_W = (80.356\pm0.125)$ GeV is predicted in terms of G_μ, α_{EM}, M_Z and m_t) [1]. However, the actual mechanism of this symmetry breaking still remains unknown and waits for experimental discovery. This is, by far, the most central question to particle physics and, in particular, to the experimental programs at LEP 2 and the LHC. It is very likely that the understanding of the mechanism of the electroweak symmetry breaking will not only provide us with the missing link in the Standard Model but, also, will be an important bridge to physics beyond it.

The minimal model for spontaneous electroweak gauge symmetry breaking is the Higgs mechanism, whose minimal version (the minimal Standard Model) requires one scalar SU(2) doublet (Higgs doublet). The scalar potential at some scale Λ is:

$$V(\Lambda) = m^2(\Lambda)|H(\Lambda)|^2 + \tfrac{1}{2}\lambda(\Lambda)|H(\Lambda)|^4 \tag{1}$$

with the dependence on Λ controlled by the renormalization group evolution (RGE). The mass of the physical scalar (Higgs boson) is $M^2_{\phi^0} = \lambda(M_Z)v^2(M_Z)$ where $v(M_Z) = \sqrt{4M^2_W/g^2_2(M_Z)} \approx 246$ GeV. There exist the well known theoretical bounds on the Higgs boson mass which follow from certain constraints on the behaviour of the self-coupling $\lambda(\Lambda)$. One can distinguish two types of bounds. The most general upper bound on M_{ϕ^0} follows from the requirement that the Standard Model is a unitary and weakly interacting theory at the energy scale $\mathcal{O}(M_Z)$. We get then $M_{\phi^0} \lesssim \mathcal{O}(1\ \text{TeV})$. Stronger bounds are Λ-dependent and are known under the names of the triviality (upper) bound and the vacuum stability (lower) bound. They follow respectively from the requirements that the theory remains perturbative ($\lambda(\Lambda) < 16\pi^2$) and the vacuum remains stable ($\lambda(\Lambda) > 0$) up to a certain scale Λ. Those bounds are particularly interesting in the presence of the heavy top quark, $m_t = (175 \pm 6)$ GeV [2]. They are shown in Fig. 1 and lead to the striking conclusions: We see that the discovery of a light Higgs boson ($M_{\phi^0} \lesssim 80$ GeV) or a heavy one ($M_{\phi^0} \gtrsim 500$ GeV) would be a direct information about the existence of new physics below the scale $\Lambda \sim \mathcal{O}(1$ TeV) (or at least of a strongly interacting Higgs sector). On the other hand, if the SM in its perturbative regime is to be valid up to very large scales Λ, of the order of the GUT scale $\Lambda \approx 10^{16}$ GeV, one gets strong bounds 140 GeV $\lesssim M_{\phi^0} \lesssim$ 180 GeV. In this case we face the well known hierarchy problem in the SM: M_{ϕ^0}, $v \ll \Lambda$ and it is difficult to understand how the

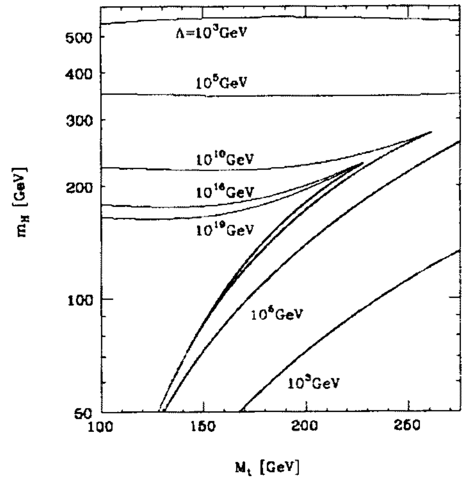

Fig. 1. Triviality (upper) and stability (lower) bounds on the SM Higgs boson mass as a function of the top-quark mass for different cut-off scales Λ. The figure has been taken from Ref. [3].

scalar potential remains stable under radiative corrections of the full theory. One way or another, the bounds on M_{ϕ^0} in the SM are strongly suggestive that the mechanism of spontaneous electroweak gauge symmetry breaking is directly related to the existence of a new scale (not much above the electroweak scale) in fundamental interactions. The central question can then be phrased as this: discover and investigate the next scale in fundamental interactions. Is it the scale of new strong interactions (strongly interacting Higgs sector or techicolour interactions or compositness scale) or the scale of soft supersymmetry breaking? We would like to stress the basic difference between these two lines of approach. In the first one, the new scale is also a cut-off scale for the perturbative validity of the electroweak theory. Supersymmetry offers a solution to the hierarchy problem while maintaining the perturbative nature of the theory up to the GUT or even Planck scale. This is a welcome feature if such facts as the gauge coupling unification are not to be considered as purely accidental.

Another difference is in the expectations for the Higgs boson mass: in the strong interaction scenarios it is naturally heavy, with its mass close to the new scale Λ. In supersymmetric extensions of the SM the lightest Higgs scalar h^0 generically remains light, $M_{h^0} \sim \mathcal{O}(100 \text{ GeV})$ and only logarithmically correlated with the scale of the soft supersymmetry breaking.

2. Is the Higgs boson light or heavy-precision tests at the Z^0 pole in the Standard Model

The bulk of the electroweak precision measurements (M_W, Z^0-pole observables, νe, lp, ...) shows that the global comparison of the SM predictions with the data shows impressive agreement. Both, the experiment and the theory have at present similar accuracy, typically $\mathcal{O}(1‰)$! The dramatic change in the data, reported recently, are the new values for R_b and R_c. Both are now in agreement with the SM. The predictions of the SM are usually given in terms of the very precisely known parameters G_μ, α_{EM}, M_Z and the other three parameters $\alpha_s(M_Z)$, m_t, M_h. The top quark mass and the strong coupling constant are now also reported from independent experiments with considerable precision: $m_t = (175 \pm 6)$ GeV [2] and $\alpha_s(M_Z) = 0.118 \pm 0.003$ [4], but those measurements are difficult and it is safer to take α_s, m_t, M_h as parameters of an overall fit. Such fits give values of m_t and α_s very well consistent with the above values [1].

The theoretical uncertainties in the SM predictions (for fixed m_t, M_h, α_s) come mainly from the RG evolution of $\alpha_{EM} \equiv \alpha(0) \to \alpha(M_Z)$ (to the scale M_Z) which depends on the hadronic contribution to the photon vacuum polarization $\alpha(s) = \alpha(0)/(1 - \Delta\alpha(s))$ where $\Delta\alpha(s) = \Delta\alpha_{\text{hadr}} + \ldots$ and $\Delta\alpha_{\text{hadr}} = 0.0280 \pm 0.0007$ [5]. The present error in the hadronic vacuum polarization propagates as $\mathcal{O}(1‰)$ errors in the final predictions. The other uncertainties come from the neglected higher order corrections and manifest themselves as renormalization scheme dependence, higher order arbitrariness in resummation formulae *etc.* Those effects have been estimated to be smaller than $\mathcal{O}(1‰)$, hence the conclusion is that the theory and experiment agree with each other at the level of $\mathcal{O}(1‰)$ accuracy. In particular, the genuine weak loop corrections are now tested at $\mathcal{O}(5\sigma)$ level and the precision is already high enough to see some sensitivity to the Higgs boson mass.

The electroweak observables depend only logarithmically on the Higgs boson mass (whereas the dependence on the top quark mass is quadratic). Global fits to the present data give $M_h \approx 145^{+160}_{-80}$ GeV and the 95% C.L. upper bound is around 600 GeV [1,6]. Thus, the data give some indication for a light Higgs boson. (It is worth noting that $M_h = 1$ TeV is $\gtrsim 3\sigma$ away from the best fit). The direct experimental lower limit on M_h is ~ 65 GeV. These results should be placed in the context of the theoretical lower and upper bounds for the SM Higgs boson mass (discussed in the previous section). We conclude that the precision data give some, although not very strong, indication for a light Higgs boson. This gives some support to the supersymmetric solution to the hierarchy problem.

3. The hierarchy problem and supersymmetry

The hierarchy [7] problem can be illustrated with a simple model of two scalar fields φ and Φ, such that their respective masses satisfy $m \ll M$. The Lagrangian density reads (we impose the symmetry under $\varphi \to -\varphi$, $\Phi \to -\Phi$):

$$\begin{aligned} \mathcal{L} = & \frac{1}{2}\left[(\partial\varphi)^2 - m^2\varphi^2 + (\partial\Phi)^2 - M^2\Phi^2\right] \\ & - \frac{\lambda_1}{4!}\varphi^4 - \frac{\lambda_2^2}{4}\varphi^2\Phi^2 - \frac{\lambda_2}{4!}\Phi^4 . \end{aligned} \tag{2}$$

For $m^2 < 0$, at the tree level there exist a ground state such that $\langle\varphi\rangle = \sqrt{-6m^2/\lambda_1}$, $\langle\Phi\rangle = 0$ and the discrete symmetry is spontaneously broken. However, quantum corrections ΔV to the effective potential destablilize this result. At 1-loop level one obtains:

$$\Delta V = \frac{1}{2}\delta m^2\varphi^2 + \frac{1}{4}\delta\lambda_1\varphi^4 + \mathcal{O}\left(\frac{1}{M^2}\right) . \tag{3}$$

where

$$\delta m^2 = \frac{1}{(4\pi)^2}\lambda_2 M^2\left(-1 + \log\frac{M^2}{\mu^2}\right) + \mathcal{O}\left(m^2\right) \tag{4}$$

$$\delta\lambda_1 = \frac{1}{(4\pi)^2}\lambda_2^2\log\frac{M^2}{\mu^2} + \mathcal{O}\left(\log m^2\right) \tag{5}$$

(μ is the renormalization scale). Since δm^2 depends quadratically on M^2, those new contributions destabilize the tree level result $\langle\varphi\rangle \ll M$, unless the parameters m^2, M^2 and λ_2 are tuned very precisely! However, the effect of fine tuning remains hidden for us if only low energy sector is observed, which can be described by the effective theory with only one field φ (Appelquist–Carrazone decoupling). Then, the relevant parameter is $m^2_{\text{eff}} = m^2 + \delta m^2$ and, as every renormalizable parameter, it is taken from experiment. Thus we face a severe hierarchy problem once we suppose that there exists a heavy field coupled to a light field in a scalar field theory.

A solution to this problem offered by supersymmetry is based on the fact that the quadratic dependence on M in the diagrams of the scalar field theory generated by the $\varphi^2\Phi^2$ interaction can be cancelled by the diagrams which couple the field φ with a Weyl spinor field (of mass M) provided the Yukawa coupling, $\lambda_y\varphi\lambda\bar{\lambda}$, satisfies the relation $\lambda_1 = \lambda_y^2$. This cancellation occurs provided the number of scalar and spin 1/2 degrees of freedom strictly match each other *i.e.* for instance a supersymmetric theory requires a complex scalar and a Weyl spinor as superpartners.

Supersymmetry is of interest for a number of reasons. It is likely that it is linked to the electroweak symmetry breaking (hierarchy problem). It is at present the only theoretical framework which allows to extrapolate to very short distances (Planck length). It is an appealing mathematical structure. Supersymmetric field theories have several interesting properties which make them more predictive that non-supersymmetric theories. And finally, on the purely pragmatic level, the Minimal Supersymmetric Standard Model is so far the only framework beyond the SM which addresses the phenomenology of elementary interactions at the electroweak scale and just above it in a complete and quantitative way. As such, it plays an important stimulating role in experimental search for physics beyond the SM.

The minimal model is based on the three main assumptions [8]:

a) minimal particle content consistent with the known spectrum and supersymmetry

b) most general soft supersymmetry breaking terms which are $SU_c(3)\times SU_L(2)\times U_Y(1)$ symmetric

c) R-parity conservation

Two "mild" extensions of the minimal model include the models with R-parity explicitly broken and those with additional full SU(5) matter multiplets, at "low" scale. New R-parity violating couplings in the superpotential must be small enough to avoid problems with the baryon and lepton number violating processes. Additional complete SU(5) multiplets do not destroy the unification of couplings.

4. Supersymmetric Higgs sector and its naturalness

In the Minimal Supersymmetric Standard Model the Higgs sector is particularly simple and predictive. Supersymmetry and the minimal particle content imply that it consists of two Higgs doublets, each coupled to only one type (H_1 (H_2) couples to the down (up)) of fermions. The scalar Higgs potential reads:

$$V = m_1^2\overline{H}_1H_1 + m_2^2\overline{H}_2H_2 + m_3^2\left(\varepsilon_{ab}H_1^aH_2^b + c.c\right) + \tfrac{1}{8}(g_1^2+g_2^2)(\overline{H}_1H_1 - \overline{H}_2H_2)^2 + \tfrac{1}{2}g_2^2|\overline{H}_1H_2|^2\,, \qquad (6)$$

where $\varepsilon_{12} = -1$ and m_1^2, m_2^2 and m_3^2 are the soft supersymmetry breaking mass parameters. The crucial point about the potential (6) is that its quartic couplings are the electroweak gauge couplings (*i.e.* there is no F-term contribution to the scalar Higgs potential). The only free parameters are the three mass parameters. The tree level mass eigenstates of the Higgs bosons are: two CP-even (h^0, H^0), one CP-odd (A^0) and 2 charged ($H^\pm$) physical particles and three Goldstone bosons "eaten up" by the gauge bosons. An important parameter is $\tan\beta \equiv v_2/v_1$ where v_i minimize the tree level potential (6) and are given by $v_1 = v\cos\beta$, $v_2 = v\sin\beta$ with

$$v^2 = \frac{8}{g_1^2+g_2^2}\frac{m_1^2 - m_2^2\tan^2\beta}{\tan^2\beta - 1} \tag{7}$$

$$\sin 2\beta = \frac{2m_3^2}{m_1^2+m_2^2}\,. \tag{8}$$

Since v is fixed by the Z^0 mass, all physical Higgs boson masses are expressed in terms of only two free parameters. They can be taken *e.g.* as $\tan\beta$ and the mass $M_{A^0}^2$ of the CP–odd Higgs scalar A^0 given by $M_{A^0}^2 = m_1^2 + m_2^2$. The CP–even Higgs boson masses then read:

$$M_{h^0,H^0}^2 = \frac{1}{2}\left(M_{A^0}^2 + M_{Z^0}^2 + \sqrt{\left(M_{A^0}^2 - M_{Z^0}^2\right)^2 - 4M_{A^0}^2M_{Z^0}^2\cos^2 2\beta}\right) \tag{9}$$

leading to the bound $M_{h^0} < M_{Z^0}$ and to the "natural" (*i.e.* independent of any other parameters) relation $M_{h^0}^2 + M_{H^0}^2 = M_{A^0}^2 + M_{Z^0}^2$. The other relation is $M_{H^\pm}^2 = M_{W^\pm}^2 + M_{A^0}^2$ [9].

The origin and the magnitude of radiative corrections to the Higgs boson masses can be easily understood. Let M be the scale of the soft supersymmetry breaking sfermion masses. Neglecting terms suppressed by inverse powers of M, the dominant one-loop corrections to the effective potential $V_{\mathbf{eff}}$, due to the top and stop loops, can be absorbed into renormalization of the parameters in the Higgs potential. One gets:

$$\begin{aligned} V &= \tilde{m}_1^2\overline{H}_1H_1 + \tilde{m}_2^2\overline{H}_2H_2 + \tilde{m}_3^2\left(\varepsilon_{ab}H_1^aH_2^b + c.c\right) \\ &\quad +\lambda_1|H_1|^4 + \lambda_2|H_2|^4 + \lambda_3|H_1|^2|H_2|^2 + \lambda_4|\overline{H}_1H_2|^2\,. \end{aligned} \tag{10}$$

The appearance of other couplings is protected by the symmetries of the model. It is clear on the dimensional grounds that

$$\delta m_i^2 = \tilde{m}_i^2 - m_i^2 \sim \mathcal{O}(M^2)\,. \tag{11}$$

They are logarithmically divergent but can be absorbed into the free parameters of the model. The corrections $\delta\lambda_i$ defined by

$$\begin{aligned} \delta\lambda_1 &= \lambda_1 - \tfrac{1}{8}(g_1^2+g_2^2), \qquad & \delta\lambda_2 &= \lambda_2 - \tfrac{1}{8}(g_1^2+g_2^2) \\ \delta\lambda_3 &= \lambda_3 + \tfrac{1}{4}(g_1^2+g_2^2), \qquad & \delta\lambda_4 &= \lambda_4 - \tfrac{1}{2}g_2^2\,. \end{aligned} \tag{12}$$

are all $\mathcal{O}(\log M)$. Moreover, from the non-renormalization theorem, the corrections $\delta\lambda_i$ are calculable (finite) in terms of the remaining parameters of the model. From the top and stop loops with attached four Higgs boson legs one gets

$$\delta\lambda_i \sim \frac{12}{16\pi^2} h_t^4 \log\left(\frac{M_{\tilde{t}}^2}{m_t^2}\right) . \tag{13}$$

where h_t is the top quark Yukawa coupling (factor 12 comes from 4 top degrees of freedom multiplied by the color factor of 3) and $M_{\tilde{t}}$ denotes the scale of the stop masses. Thus, the correction to the h^0 mass is [10]

$$\delta M_{h^0}^2 \sim \mathcal{O}\left(\frac{6g_2^2}{16\pi^2}\frac{m_t^4}{M_W^2}\log\left(\frac{M_{\tilde{t}}^2}{m_t^2}\right)\right) . \tag{14}$$

In general, taking into account the full structure of the stop mass matrix, the lightest Higgs boson mass in the MSSM is parametrized by

$$M_{h^0} = M_{h^0}\left(M_{A^0}, \tan\beta, m_t, M_{\tilde{t}_1}, M_{\tilde{t}_2}, A_t, \mu, \ldots\right) , \tag{15}$$

where $M_{\tilde{t}_i}$ are the physical stop masses, A_t and μ determine their mixing angle (as well as some of their trilinear couplings to the Higgs bosons) and ellipsis stand for other parameters whose effects are not dominant (*e.g.* the sbottom sector parameters).

In Fig. 2 we show M_{h^0} as a function of M_{A^0} for two values of $\tan\beta$ and $M_{\tilde{t}_1} = M_{\tilde{t}_2} = 1$ TeV, $\mu = 0$ and two values of the A_t parameter. We see that maximal M_{h^0} is always obtained for $M_{A^0} \gg M_{Z^0}$ (in practice, the bound is saturated for $M_{A^0} \gtrsim 250$ GeV).

We conclude that in the MSSM there exists a strong upper bound on the lightest Higgs boson mass which depends only logarithmically on the (unknown but expected to be at most $\mathcal{O}(1$ TeV) — see below) scale of soft masses of the third generation squarks.

Finally, we are going to discuss several constraints on the range of the soft supersymmetry breaking masses m_1^2, m_2^2, m_3^2, the top squark masses and the mixing parameters A_t, μ (*i.e.* the remaining parameters relevant for M_h) which follow from the extrapolation of the MSSM to high energies.

There has been often addressed the question of fine-tuning (large cancellations) in the Higgs potential in models with the soft terms generated at large scales [11–13]. Indeed, if supersymmetry is to be the solution to the hierarchy problem in the SM, it should not introduce another fine-tuning in the Higgs potential. The origin of the problem is easy to see. Using the renormalization group equations we can express M_{Z^0} for a given $\tan\beta$ in terms

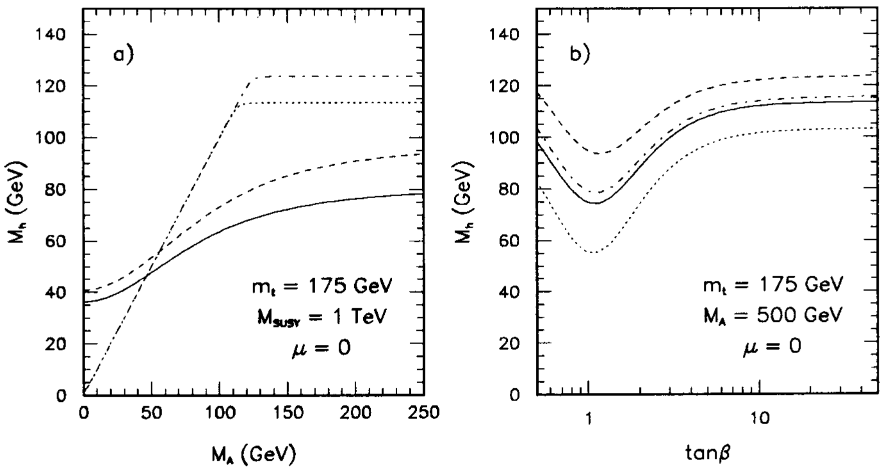

Fig. 2. Radiatively corrected M_h in the MSSM (1- and 2-loops). (a) As a function of the CP-odd Higgs mass for $M_{\rm SUSY} = 1$ TeV and for $\tan\beta = 1.5$ (solid and dashed lines) and $\tan\beta = 50$ (dotted and dash-dotted lines). Lower (upper) lines correspond to $A_t = 0$ $(2.5 M_{\rm SUSY})$ (b) As a function of $\tan\beta$ for $m_Q = m_U =$ 1 TeV, $A_t = 0$ (2.5 TeV) solid (dashed) line and for $m_Q = 500$, $m_U = 100$ GeV, $A_t = 0$ (1 TeV) dotted (dash-dotted) line.

of the initial values $m_K^2(0)$ of the scalar masses ($K = H_1,\ H_2,\ Q,\ U,\ D$ *etc.*), $M_{1/2}$ (universal gaugino mass) and the μ parameter:

$$\begin{aligned} M_{Z^0}^2 &= -2\mu^2(t) + a_{H_1} m_{H_1}^2(0) + a_{H_2} m_{H_2}^2(0) + a_{QU}\left(m_Q^2(0) + m_U^2(0)\right) \\ &\quad + a_{AA} A_t^2(0) + a_{AM} A_t(0) M_{1/2} + a_M M_{1/2}^2 . \end{aligned} \tag{16}$$

For $m_t = 175$ GeV the generic values of the coefficients in Eq. (16) in the supergravity scenario *e.g.* for $\tan\beta \approx 1.65(2.2)$ are $a_{H_1} \approx 1.1(0.5)$, $a_{H_2} \approx -1.7(-1.5)$, $a_{QU} \approx 1.5(1.1)$, $a_{AA} \approx 0.1(0.2)$, $a_{AM} \approx -0.3(-0.7)$, $a_M \approx 11.1(8.3)$. Eq. (16) shows that for values of μ, $M_{1/2}$ and/or $m_K^2(0)$ much larger than M_{Z^0} one needs large cancellations. Asymptotically, we are back to the hierarchy problem in the SM. Although the idea of "naturalness" is only qualitative, one can at least correlate the magnitude of the necessary cancellations with the values of the parameters μ, $m_K^2(0)$ and $M_{1/2}$ and, in consequence, with the low energy mass parameters. One notes, in particular, that the smallness of a_{QU} (compared to a_M) puts weaker constraints on the "natural" values of $m_K^2(0)$ than large a_M does on $M_{1/2}$. However, in the physical spectrum this effect is partially counterbalanced by the fact that the stop soft masses tend to be suppressed compared to $m_{Q,U}(0)$ by the running with large top quark Yukawa coupling. This effect is stronger for the right handed stop than for the left handed one and gives the hierarchy $M_{\tilde{t}_R} < M_{\tilde{t}_L}$. Important source of fine-tuning can also be the relation (8) which correlates the values of $\tan\beta$ and the B_0 parameter. Finally it is important to observe that, only the third generation squark masses enter into the Eq. (16). Thus, squarks of the first two generations can be very heavy without facing the problem of naturalness in the Higgs potential.

5. MSSM and electroweak measurements

The simplest interpretation of the success of the SM is that the superpartners are heavy enough to decouple from the electroweak observables. Explicit calculations (with the same precision as in the SM) show that this happens if the common supersymmetry breaking scale is $\geq \mathcal{O}(300-400)$GeV. This is very important as such a scale of supersymmetry breaking is still low enough for supersymmetry to cure the hierarchy problem. However, in this case the only supersymmetric signature at the electroweak scale and just above it is the Higgs sector withe a light, $M_h \sim \mathcal{O}(100$ GeV), Higgs boson. This prediction is consistent with the SM fits discussed earlier. We can, therefore, conclude at this point that the supersymmetric extension of the SM, with all superpartners $\geq \mathcal{O}(300)$ GeV, is phenomenologically as successful as the SM itself and has the virtue of solving the hierarchy problem. Discovery of a light Higgs boson is the crucial test for such an extension.

The relatively heavy superpartners discussed in the previous paragraph are sufficient for explaining the success of the SM. But is it necessary that all of them are that heavy? Is there a room for some light superpartners with masses $\mathcal{O}(M_Z)$ or even below? Since the Higgs boson mass in the MSSM is very strongly constrained, we may hope that the precision electroweak data show some sensitivity to the superpartner spectrum. This question is of great importance for LEP2. Indeed, a closer look at the electroweak observables shows that the answer to this question is positive. The dominant quantum corrections to the electroweak observables are the so-called "oblique" corrections to the gauge boson self-energies. They are economically summarized in terms of the S, T, U parameters [14]

$$S \sim \Pi'_{3Y}(0) = \Pi'_{L3,R3} + \Pi'_{L3,B-L}\,, \tag{17}$$

(the last decomposition is labelled by the $SU_L(2)\times SU_R(2)\times U_{B-L}(1)$ quantum numbers)

$$\alpha T \equiv \Delta\rho \sim \Pi_{11}(0) - \Pi_{33}(0)\,, \tag{18}$$

$$U \sim \Pi'_{11}(0) - \Pi'_{33}(0)\,, \tag{19}$$

where $\Pi_{ij}(0)$ ($\Pi'_{ij}(0)$) are the (i,j) left -handed gauge boson self-energies at the zero momentum (their derivatives) and the self- energy correction to the S parameter mixes $W^{\pm}_{\mu}$ and B_{μ} gauge bosons. It is clear from their definitions that the parameters S, T, U have important symmetry properties [15]: T and U vanish in the limit when quantum corrections to the left-handed gauge boson self-energies have unbroken "custodial" $\mathrm{SU}_V(2)$ symmetry. The parameter S vanishes if $\mathrm{SU}_L(2)$ is an exact symmetry (notice that, since $\mathbf{3_L} \times \mathbf{3_R} = \mathbf{1} \oplus \mathbf{5}$ under $\mathrm{SU}_V(2)$, exact $\mathrm{SU}_V(2)$ is not sufficient for the vanishing of S) [15]. The success of the SM means that it has just the right amount of the $\mathrm{SU}_V(2)$ breaking (and of the $\mathrm{SU}_L(2)$ breaking), encoded mainly in the top quark-bottom quark mass splitting. Any extension of the SM, to be consistent with the precision data, should not introduce additional sources of large $\mathrm{SU}_V(2)$ breaking in sectors which couple to the left-handed gauge bosons. In the MSSM, the main potential origin of new $\mathrm{SU}_V(2)$ breaking effects in the left-handed sector is the splitting between the left-handed stop and sbottom masses:

$$\begin{aligned} M^2_{\tilde{t}_L} &= m^2_Q + m^2_t - \cos 2\beta (M^2_Z - 4M^2_W)\,, \\ M^2_{\tilde{b}_L} &= m^2_Q + m^2_b - \cos 2\beta (M^2_Z + 2M^2_W)\,. \end{aligned} \tag{20}$$

The $\mathrm{SU}_V(2)$ breaking is small if the common soft mass m^2_Q is large enough. So, from the bulk of the precision data one gets a lower bound on the masses of the left-handed squarks of the third generation [1]. However, the right-handed squarks can be very light, at their experimental lower bound ~ 45 GeV. Another interesting observation is that in the low $\tan\beta$ region the top squark masses are strongly constrained also by the present experimental lower bound on the lightest supersymmetric Higgs boson mass, $M_h \geq 60$ GeV. For low $\tan\beta$, the tree level Higgs boson mass is close to zero and radiative corrections are very important. They depend logarithmically on the product $M_{\tilde{t}_1} M_{\tilde{t}_2}$.

In Fig. 3 we show the lower bound on the mass of the heavier top squark as a function of the mass of the lighter stop, which follows from the requirement that a fit in the MSSM is at most by $\Delta\chi^2 = 2$ worse than the analogous fit in the SM and from the lower bound on M_h. The limits on the stop masses obtained from the bound on M_h are of similar strength as the χ^2 limits but apply only for low $\tan\beta$. The important role played in the fit by the precise result for $\sin^2\theta^{\mathrm{lept}}_{\mathrm{eff}}$ is illustrated in Fig. 4(a). The world average value is obtained in the SM model with $m_t = (175 \pm 6)$ GeV and $M_h \sim$

[1] Additional source of the $\mathrm{SU}_V(2)$ breaking is in the A-terms. In principle, there can be cancellations between the soft mass terms and the A-terms, such that another solution with small SU $_V(2)$ breaking exists with a large inverse hierarchy $m^2_U \gg m^2_Q$. This is very unnatural from the point of view of the GUT boundary conditions and here we assume $m^2_Q > m^2_U$.

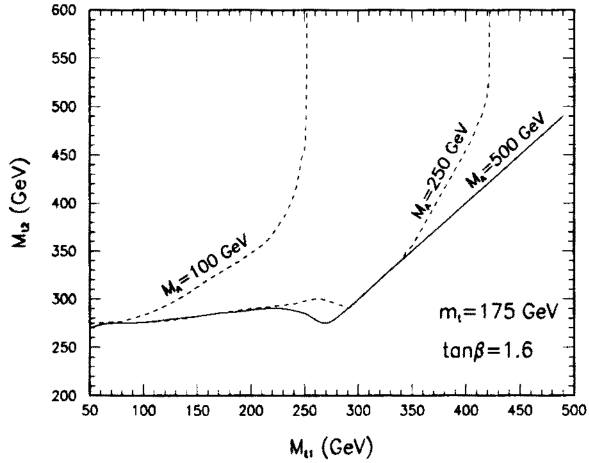

Fig. 3. Lower bounds on the heavier stop, $\tilde{t}_2$, as a function of $M_{\tilde{t}_1}$. The solid line shows the bound from $M_h > 60$ GeV and the requirement of a good χ^2 fit to the electroweak observables. The dashed lines show the bounds obtained from the $b \to s\gamma$ constraint for two values of the CP-odd Higgs boson mass M_A. The latter bound is obtained under the assumption that the chargino masses $m_{C_1^\pm} = 90$ GeV.

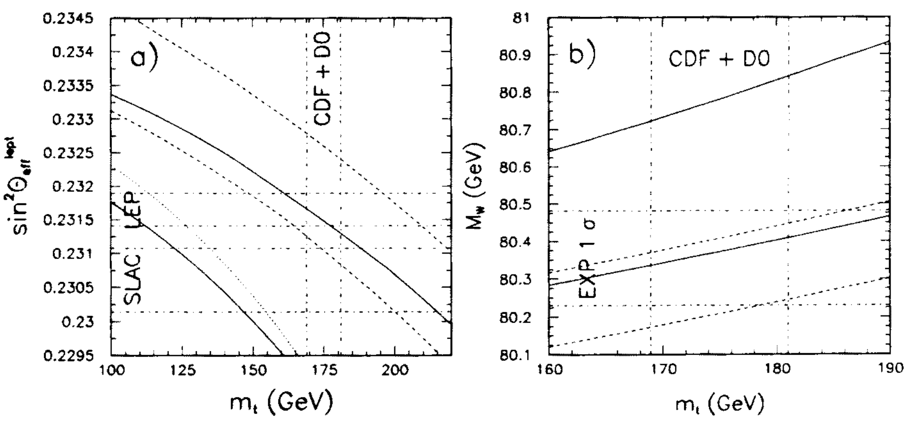

Fig. 4. Predictions for $\sin^2\theta_{\rm eff}^{\rm lept}$ (a) and M_W (b) in the SM (the band bounded by the dashed lines) and in the MSSM (solid lines) as functions of the top-quark mass. The SLC and the (average) LEP measurements for $\sin^2\theta_{\rm eff}^{\rm lept}$ and 1σ experimental range for M_W are marked by horizontal dash-dotted lines. Dotted line in (a) shows the lower limit for $\sin^2\theta_{\rm eff}^{\rm lept}$ in the MSSM if all sparticles are heavier than Z^0.

(120–150) GeV, with little room for additional supersymmetric contribution. Hence, the relevant superpartners ($\tilde{t}_L$ and $\tilde{b}_L$) have to be heavy. With lighter superpartners, one obtains the band (solid lines) shown in Fig. 4(a). We see that the SLD result for $\sin^2\theta_{\rm eff}^{\rm lept}$ leaves much more room for light $\tilde{t}_L$ and $\tilde{b}_L$. Thus, settling the SLD/LEP dispute is very relevant for new physics. Similar dependence for M_W is shown in Fig. 4(b).

All squarks of the first two generations as well as sleptons are almost unconstrained by the precision electroweak data. The same applies to the gaugino/higgsino sectors, since they do not give any strong $SU_V(2)$ breaking effects. In conclusion, most of the superpartners decouple from most of the electroweak observables, even if very light, $\leq \mathcal{O}(M_Z)$. This high degree of screening follows from the basic structure of the model.

The remarkable exception is the famous R_b [16]. Additional supersymmetric contributions to the $Z^0\bar{b}b$ vertex, precisely from the chargino-right-handed stop loop, can be non-negligible when both are light (and from the

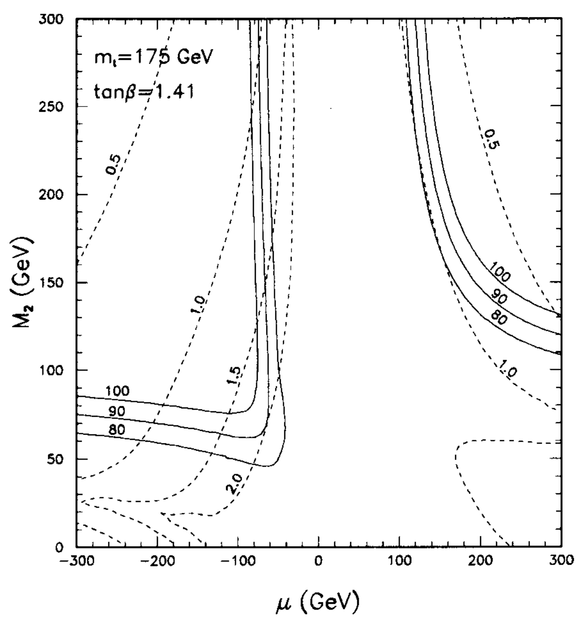

Fig. 5. Contours of constant lighter chargino masses $m_{C_1^\pm}$ =80, 90, 100 GeV (solid lines) and of $\delta R_b \times 10^3$ =2.0, 1.5, 1.0, 0.5 (dashed lines) in the (μ, M_2) plane for m_t = 175 GeV and $\tan\beta$ = 1.41. The region below the lines $m_{C_1^\pm}$ =80 GeV is excluded after the LEP run at $s^{1/2}$ = 161 GeV.

CP-odd Higgs loop in the large $\tan\beta$ region). (Note that those contributions do not change the value of $\mathcal{A}_b$ as they dominantly modify the left-handed effective coupling.) However, even with the chargino and stop very light, at their present experimental mass limit, in the MSSM the prediction for R_b depends on the chargino composition and on the stop mixing angle. The values ranging from 0.2158 (the SM prediction) up to 0.218 (0.219) for small (large) $\tan\beta$ can be realistically obtained (given all the experimental constraints) [17].

No significant modification of the SM result for R_c is possible, though. This predictions hold with or without R-parity conservation and with or without the GUT relation for the gaugino masses. The upper bound is reachable for chargino masses up to $\mathcal{O}(90$ GeV) provided they are mixed gaugino-higgsino states ($M_2/|\mu| \sim 1$). In the same chargino mass range $\delta R_b \to 0$ in the deep higgsino and gaugino regions. Clearly, the new values of R_b and R_c are good news for supersymmetry! At the same time, one should face the fact that, unfortunately, in the MSSM

$$\delta R_b^{\max} \sim \mathcal{O}(1\ \sigma^{\exp})$$

so much better experimental precision is needed for a meaningful discussion. The contours of δR_b in the (M_2, μ) plane are shown in Fig. 5.

6. Have light superpartners already been discovered?

The constraints on the superpartner spectrum which we have discussed so far apply to the MSSM and to its "mild" extensions such as the so-called gauge mediated models with the lightest sparticle decaying into gravitino and the models with broken R-parity. However, those models often have different from the SM signatures for the direct search for superpartners and therefore such direct limits as $m_{\tilde{X}}$ >85 GeV (based on the missing energy signature) do not apply. A considerable attention has been recently paid to several exotic pieces of experimental information: a single event $e^+e^-\gamma\gamma + missing\ E_T$ has been reported by the CDF, the results from the LEP 1.5 run at $\sqrt{s} = 136$ GeV include peculiar four jet events reported by ALEPH and the results from Hera may suggest a production of a new particle. Those findings should be taken with extreme caution and are likely to be a statistical fluctuation. Nevertheless, they generate some speculations on being a possible manifestation of supersymmetry. If so they would require precisely those previously mentioned mild extension of the MSSM. The CDF event can be interpreted as a selectron pair production with a subsequent chain decay:

$$p\bar{p} \to \tilde{e}^+\tilde{e}^- \to (e^+X_2)(e^-X_2) \to \\ \to (e^+X_1\gamma)(e^-X_1\gamma)\,, \qquad (21)$$

where X_1 is the LSP which carries the missing energy and X_2 is next-to-the LSP particle. The event can be interpreted in two possible ways [18]:

a) X_2 - neutralino (gaugino)
   X_1 - neutralino (higgsino)

The signatures of the event are reproduced for $m_{X_2} - m_{X_1} \geq 30$ GeV, $\tan\beta \sim 1$ and "nonunified" gaugino masses $M_1 \approx M_2$. This interpretation is consistent with a light supersymmetric spectrum of the type discussed earlier (in particular, the one which may give some enhancement in R_b).

b) X_2 - neutralino
   X_1 - gravitino ($\tilde{G}$) with BR($X_2 \to \tilde{G}\gamma$) ~ 1,
   $M_{X_2} \leq 100$ GeV, $m_{\tilde{G}} \leq 250$ eV (for X_2 to decay in the detector)

This second interpretation fits nicely into the ideas of the so-called gauge mediated low energy supersymmetry breaking [19].

The supersymmetric interpretation of the ALEPH four-jet events is also possible, though not strikingly "natural". Their main signatures (the absence of missing energy and of b-quark jets) can be consistent with the so-called light gluino scenario [20] or otherwise needs broken R-parity. In the latter case the events can be interpreted as a production of a pair of sparticles [22,23] (sneutrinos or right-handed stops or charginos) with subsequent R-parity violating decay into a pair of quarks via baryon or lepton number violating couplings.

Hera excess [21] in the structure functions can be interpreted as a leptoquark production or a squark production with R-parity non-conservation. It is fair to wait for experimental clarification before speculating further.

7. Conclusions

Although the Standard Model is strikingly successful in its description of the electroweak data, we need new physics to "explain" M_W. The discovery of the Higgs boson and the measurement of its mass are important clues to physics beyond the SM. Light versus heavy Higgs boson has its correspondence in supersymmetry versus dynamical electroweak symmetry breaking.

Supersymmetric extension of the SM is not only theoretically motivated but naturally accommodates the success of the SM. There are important constraints on supersymmetric spectrum from naturalness and from precision tests. Large m_t is crucial for those constraints. However some of the superpartners maybe very light $m \leq M_Z$. This is a consequence of the structure of the theory and not of fine-tuning of its parameters.

Very light superpartners (*e.g.* $C^{\pm}$, $\tilde{t}_R$, N^0, ...) may have important effects on few selected observables such as R_b, $b \to s\gamma$, B^0-$\bar{B}^0$, ... which, however, require still better experimental accuracy to be confirmed.

We have a couple of exotic experimental observations which could find their interpretation in the supersymmetric framework with mild extension of the minimal supersymmetric model. Extreme caution with any firm conclusion is, however, advised before further experimental clarification.

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