NEW PHYSICS AND THE LHC

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In these lectures I start by briefly reviewing the status of the electroweak theory, in the Standard Model and beyond. I then discuss the motivation and the possible avenues for new physics, on the brink of the LHC start.

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1. The programme of LHC physics

The first collisions at the LHC are expected in '08 and the physics run at 14 TeV will start soon after. The particle physics community eagerly waits for the answers that one expects from the LHC to a number of big questions. The main physics issues at the LHC, addressed by the ATLAS and CMS collaborations, will be: 1) the experimental clarification of the Higgs sector of the electroweak (EW) theory, 2) the search for new physics at the weak scale that, on conceptual grounds, one predicts should be in the LHC discovery range, and 3) the identification of the particle(s) that make the dark matter in the Universe. In addition the LHCb detector will be devoted to the study of precision B physics, with the aim of going deeper in the knowledge of the Cabibbo-Kobayashi-Maskawa (CKM) matrix and of CP violation. The LHC will also devote a number of runs to accelerate heavy ions and the ALICE collaboration will study their collisions for an experimental exploration of the QCD phase diagram.

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2. The Higgs problem

The Higgs problem is really central in particle physics today. On the one hand, the experimental verification of the Standard Model (SM) cannot be considered complete until the physics of the Higgs sector is not established by experiment. On the other hand, the Higgs is directly related to most of the major open problems of particle physics, like the flavour problem or the hierarchy problem, the latter strongly suggesting the need for new physics near the weak scale (which could possibly clarify the dark matter identity). It is clear that the fact that some sort of Higgs mechanism is at work has already been established. The W or the Z with longitudinal polarization that we observe are not present in an unbroken gauge theory (massless spin-1 particles, like the photon, are transversely polarized). The longitudinal degree of freedom for the W or the Z is borrowed from the Higgs sector and is an evidence for it. Also, the couplings of quarks and leptons to the weak gauge bosons $W^{\pm}$ and Z are indeed precisely those prescribed by the gauge symmetry. To a lesser accuracy the triple gauge vertices γWW and ZWW have also been found in agreement with the specific predictions of the $SU(2) \bigotimes U(1)$ gauge theory. This means that it has been verified that the gauge symmetry is unbroken in the vertices of the theory: all currents and charges are indeed symmetric. Yet there is obvious evidence that the symmetry is instead badly broken in the masses. Not only the W and the Z have large masses, but the large splitting of, for example, the t-b doublet shows that even a global weak SU(2) is not at all respected by the fermion spectrum. Symmetric coupling and completely non symmetric spectrum are a clear signal of spontaneous symmetry breaking which, in a gauge theory, is implemented via the Higgs mechanism. The big remaining questions are about the nature and the properties of the Higgs particle(s).

The present experimental information on the Higgs sector, mainly obtained from LEP as described in section 4, is surprisingly limited. It can be summarized in a few lines, as follows. First, the relation $M_W^2 = M_Z^2 \cos^2\theta_W$, modified by small, computable radiative corrections, has been experimentally proven. This relation means that the effective Higgs (be it fundamental or composite) is indeed a weak isospin doublet. The Higgs particle has not been found but, in the SM, its mass can well be larger than the present direct lower limit $m_H > 114.4$ GeV (at 95% c.l.) obtained from searches at LEP-2. As we shall see, the radiative corrections computed in the SM when compared to the data on precision electroweak tests lead to a clear indication for a light Higgs, not too far from the present lower bound.

The experimental upper limit on m_H, obtained from fitting the data in the SM, depends on the value of the top quark mass m_t (the one-loop radiative corrections are quadratic in m_t and logarithmic in m_H). The CDF and D0 combined value after Run II is at present[1] $m_t = 172.6 \pm 1.4$ GeV (it went down with respect to the value $m_t = 178 \pm 4.3$ GeV from Run I and also the experimental error is now sizably reduced). As a consequence the present limit on m_H is more stringent: $m_H < 190$ GeV (at 95% c.l., after including the information from the 114.4 GeV direct bound). On the Higgs the LHC will address the following questions: do the Higgs particles actually exist? How many: one doublet, several doublets, additional singlets? SM Higgs or SUSY Higgses? Fundamental or composite (of fermions, of WW ...)? Pseudo-Goldstone boson of an enlarged symmetry? A manifestation of large extra dimensions (5th component of a gauge boson, an effect of orbifolding or of boundary conditions...)? Or some combination of the above or something so far unthought of?

3. Theoretical bounds on the SM Higgs

The LHC has been designed to solve the Higgs puzzle. In the SM lower and upper limits on the Higgs mass can be derived from theoretical considerations. It is well known[2,3,4] that in the SM with only one Higgs doublet a lower limit on m_H can be derived from the requirement of vacuum stability (or, in milder form, from a moderate instability, compatible with the lifetime of the Universe[5]). The limit is a function of m_t and of the energy scale Λ where the SM model breaks down and new physics appears. The Higgs mass enters because it fixes the initial value of the quartic Higgs coupling λ for its running up to the large scale Λ. Similarly an upper bound on m_H (with mild dependence on m_t) is obtained[6] from the requirement that in λ, up to the scale Λ, no Landau pole appears, or in more explicit terms, that the perturbative description of the theory remains valid. The upper limit on the Higgs mass in the SM is clearly important for assessing the chances of success of the LHC as an accelerator designed to solve the Higgs problem. Even if Λ is as small as a few TeV the limit is $m_H < 600 - 800$ GeV and becomes $m_H < 180$ GeV for $\Lambda \sim M_{Pl}$. We now briefly recall the derivation of these limits.

The possible instability of the Higgs potential $V[\phi]$ is generated by the quantum loop corrections to the classical expression of $V[\phi]$. At large ϕ the derivative $V'[\phi]$ could become negative and the potential would become unbound from below. The one-loop corrections to $V[\phi]$ in the SM are well known and change the dominant term at large ϕ according to $\lambda\phi^4 \to (\lambda + \gamma \log \phi^2/\Lambda^2)\,\phi^4$. The one-loop approximation is not enough in this case, because it fails at large enough ϕ, when $\gamma \log \phi^2/\Lambda^2$ becomes of order 1. The renormalization group improved version of the corrected potential leads to the replacement $\lambda\phi^4 \to \lambda(\Lambda)\phi'^4(\Lambda)$ where $\lambda(\Lambda)$ is the running coupling and $\phi'(\mu) = \phi \exp \int^t \gamma(t')dt'$, with $\gamma(t)$ being an anomalous dimension function and $t = \log \Lambda/v$ (v is the vacuum expectation value $v = (2\sqrt{2}G_F)^{-1/2}$). As a result, the positivity condition for the potential amounts to the requirement that the running coupling $\lambda(\Lambda)$ never becomes negative. A more precise calculation, which also takes into account the quadratic term in the potential, confirms that the requirements of positive $\lambda(\Lambda)$ leads to the correct bound down to scales Λ as low as ~ 1 TeV. The running of $\lambda(\Lambda)$ at one loop is given by:

$$\frac{d\lambda}{dt} = \frac{3}{4\pi^2}\left[\lambda^2 + 3\lambda h_t^2 - 9h_t^4 + \text{small gauge and Yukawa terms}\right], \qquad (1)$$

with the normalization such that at $t = 0$, $\lambda = \lambda_0 = m_H^2/2v^2$ and, for the top Yukawa coupling, $h_t^0 = m_t/v$. We see that, for m_H small and m_t fixed at its measured value, λ decreases with t and can become negative. If one requires that λ remains positive up to $\Lambda = 10^{15}$–10^{19} GeV, then the resulting bound on m_H in the SM with only one Higgs doublet is given by[4]:

$$m_H(\mathrm{GeV}) > 132 + 2.1\,[m_t - 172.6] - 4.5\,\frac{\alpha_s(m_Z) - 0.118}{0.006}. \qquad (2)$$

Note that this limit is evaded in models with more Higgs doublets. In this case the limit applies to some average mass but the lightest Higgs particle can well be below, as it is the case in the minimal SUSY extension of the SM (MSSM).

The upper limit on the Higgs mass in the SM is clearly important for assessing the chances of success of the LHC as an accelerator designed to solve the Higgs problem. The upper limit[6] arises from the requirement that the Landau pole associated with the non asymptotically free behaviour of the $\lambda\phi^4$ theory does not occur below the scale Λ. The initial value of λ at the weak scale increases with m_H and the derivative is positive at large λ (because of the positive λ^2 term in eq.(1) – the $\lambda\varphi^4$ theory is not asymptotically free – which overwhelms the negative top-Yukawa term). Thus if m_H is too large the point where λ, computed from the perturbative beta function, becomes infinite (the Landau pole) occurs at too low an energy. Of course in the vicinity of the Landau pole the 2-loop evaluation of the beta function is not reliable. Indeed the limit indicates the frontier of the domain where the theory is well described by the perturbative expansion. Thus the quantitative evaluation of the limit is only indicative, although it has been to some extent supported by simulations of the Higgs sector of the EW theory on the lattice. For the upper limit on m_H one nds[6] m_H < 180 GeV for M_{GUT} M_{Pl} and m_H < 0:5 0:8 TeV for 1 TeV. Actually, for m_t 172 GeV, only a small range of values for m_H is allowed, 130 < m_H < 200 GeV, if the SM holds up to M_{GUT} or M_{Pl}. An additional argument indicating that the solution of the Higgs problem cannot be too far away is the fact that, in the absence of a Higgs particle or of an alternative mechanism, violations of unitarity appear in scattering amplitudes involving longitudinal gauge bosons (those most directly related to the Higgs sector) at energies in the few TeV range[7]. In conclusion, it is very unlikely that the solution of the Higgs problem can be missed at the LHC which has a good sensitivity up to m_H 1 TeV.

4. Precision tests of the standard electroweak theory

The most precise tests of the electroweak theory apply to the QED sector. The anomalous magnetic moments of the electron and of the muon are among the most precise measurements in the whole of physics. Recently there have been new precise measurements of a_e and a for the electron[8] and the muon[9] ($a = (g \quad 2)/2$). On the theory side, the QED part has been computed analytically for i = 1;2;3, while for i = 4 there is a numerical calculation (see, for example, ref.[11]). Some terms for i = 5 have also been estimated for the muon case. The weak contribution is from W or Z exchange. The hadronic contribution is from vacuum polarization insertions and from light by light scattering diagrams. For the electron case the weak contribution is essentially negligible and the hadronic term does not introduce an important uncertainty. As a result the a_e measurement can be used to obtain the most precise determination of the ne structure constant[10]. In the muon case the experimental precision is less by about 3 orders of magnitude, but the sensitivity to new physics e ects is typically increased by a factor $(m \; / m_e)^2$ $4 \cdot 10^4$. The dominant theoretical ambiguities arise from the hadronic terms in vacuum polarization and in light by light scattering. If the vacuum polarization terms are evaluated from the $e^+ e$ data a discrepancy of 3 is obtained (the data would indicate better agreement, but the connection to a is less direct and recent new data have added solidity to the $e^+ e$ route)[12]. Finally, we note that, given the great accuracy of the a measurement and the estimated size of the new physics contributions, for example from SUSY, it is not unreasonable that a first signal of new physics would appear in this quantity.

The results of the electroweak precision tests as well as of the searches for the Higgs boson and for new particles performed at LEP and SLC are now available in final form [1]. Taken together with the measurements of m_t, m_W and the searches for new physics at the Tevatron, and with some other data from low energy experiments, they form a very stringent set of precise constraints to be compared with the SM or with any of its conceivable extensions[13]. All high energy precision tests of the SM are summarized in fig. 1[1]. For the analysis of electroweak data in the SM one starts from the input parameters: as in any renormalizable theory masses and couplings have to be specified from outside. One can trade one parameter for another and this freedom is used to select the best measured ones as input parameters. Some of them, α, G_F and m_Z, are very precisely known, some other ones, $m_{f_{light}}$, m_t and $\alpha_s(m_Z)$ are far less well determined while m_H is largely unknown. Among the light fermions, the quark masses are badly known, but fortunately, for the calculation of radiative corrections, they can be replaced by $\alpha(m_Z)$, the value of the QED running coupling at the Z mass scale. The value of the hadronic contribution to the running, $\Delta\alpha^{(5)}_{had}(m_Z)$, reported in Fig. 1, is obtained through dispersion relations from the data on $e^+e^- \to$ hadrons at low centre-of-mass energies [1]. From the input parameters one computes the radiative corrections to a sufficient precision to match the experimental accuracy. Then one compares the theoretical predictions with the data for the numerous observables which have been measured, checks the consistency of the theory and derives constraints on m_t, $\alpha_s(m_Z)$ and m_H.

The computed radiative corrections include the complete set of one-loop diagrams, plus some selected large subsets of two-loop diagrams and some sequences of resummed large terms of all orders (large logarithms and Dyson resummations). In particular large logarithms, e.g., terms of the form $(\alpha/\pi \ln(m_Z/m_{f'}))^n$ where f' is a light fermion, are resummed by well-known and consolidated techniques based on the renormalisation group. For example, large logarithms dominate the running of α from m_e, the electron mass, up to m_Z, which is a 6% effect, much larger than the few per mil contributions of purely weak loops. Also, large logs from initial state radiation dramatically distort the line shape of the Z resonance observed at LEP-1 and SLC and have been accurately taken into account in the measurement of the Z mass and total width.

Among the one loop EW radiative corrections a remarkable class of con-

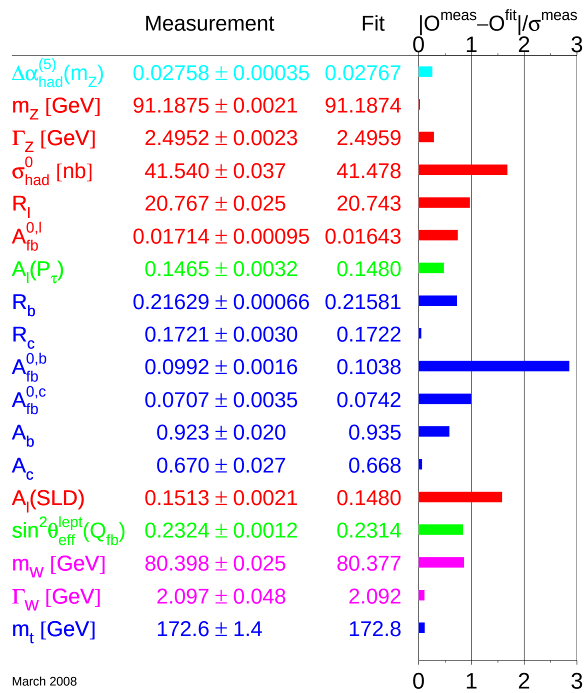

Figure 1. Precision tests of the Standard EW theory from LEP, SLC and the TeVatron (March '08).

tributions are those terms that increase quadratically with the top mass. The large sensitivity of radiative corrections to m_t arises from the existence of these terms. The quadratic dependence on m_t (and possibly on other widely broken isospin multiplets from new physics) arises because, in spontaneously broken gauge theories, heavy loops do not decouple. On the contrary, in QED or QCD, the running of α and α_s at a scale Q is not affected by heavy quarks with mass $M \gg Q$. According to an intuitive decoupling theorem [14], diagrams with heavy virtual particles of mass M can be ignored for $Q \ll M$ provided that the couplings do not grow with M and that the theory with no heavy particles is still renormalizable. In the spontaneously broken EW gauge theories both requirements are violated. First, one important difference with respect to unbroken gauge theories is in the longitudinal modes of weak gauge bosons. These modes are generated by the Higgs mechanism, and their couplings grow with masses (as is also the case for the physical Higgs couplings). Second, the theory without the top quark is no more renormalizable because the gauge symmetry is broken if the b quark is left with no partner (while its couplings show that the weak isospin is 1/2). Because of non decoupling precision tests of the electroweak theory may be sensitive to new physics even if the new particles are too heavy for their direct production.

While radiative corrections are quite sensitive to the top mass, they are unfortunately much less dependent on the Higgs mass. If they were sufficiently sensitive, by now we would precisely know the mass of the SM Higgs. In fact, the dependence of one loop diagrams on m_H is only logarithmic: $\sim G_F m_W^2 \log(m_H^2/m_W^2)$. Quadratic terms $\sim G_F^2 m_H^2$ only appear at two loops and are too small to be important. The difference with the top case is that $m_t^2 - m_b^2$ is a direct breaking of the gauge symmetry that already affects the relevant one loop diagrams, while the Higgs couplings to gauge bosons are "custodial-SU(2)" symmetric in lowest order.

The various asymmetries determine the effective electroweak mixing angle for leptons with highest sensitivity. The weighted average of all results, including small correlations, is:

$$\sin^2\theta_{eff} = 0.23153 \pm 0.00016 . \qquad (3)$$

Note, however, that this average has a χ^2 of 11.8 for 5 degrees of freedom, corresponding to a probability of 3.7%. The χ^2 is pushed up by the two most precise measurements of $\sin^2\theta_{eff}$, namely those derived from the measurements of A_l by SLD, dominated by the left-right asymmetry A_{LR}, and of the forward-backward asymmetry measured in $b\bar{b}$ production at LEP, A^b_{FB}, which differ by about 3.2 σ's. In general, there appears to be a discrepancy between $\sin^2\theta_{eff}$ measured from leptonic asymmetries ($(\sin^2\theta_e)_l$) and from hadronic asymmetries ($(\sin^2\theta_e)_h$), as seen from Figure 2. In fact, the result from A_{LR} is in good agreement with the leptonic asymmetries measured at LEP, while all hadronic asymmetries, though their errors are large, are better compatible with the result of A^b_{FB}. This very unfortunate fact makes the interpretation of precision tests less sharp and some perplexity remains: is it an experimental error or a signal of some new physics?

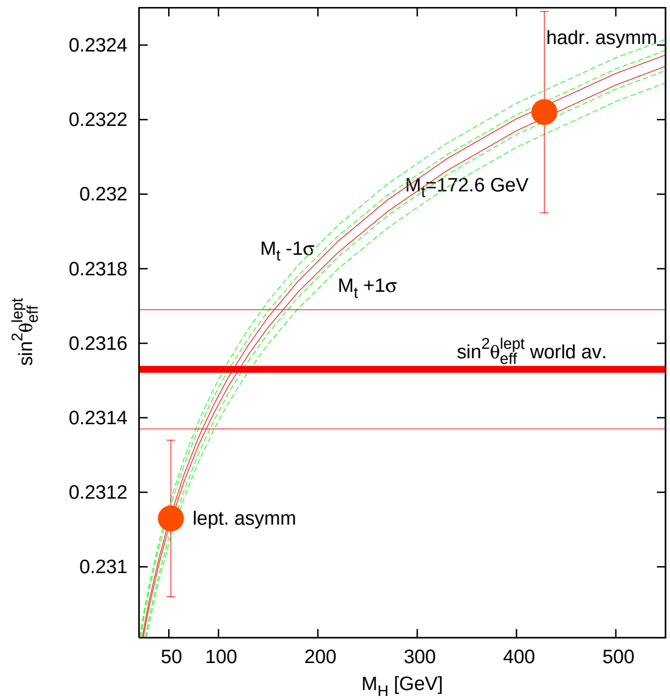

Figure 2. The data for $\sin^2\theta_{e}^{lept}$ are plotted vs m_H. For presentation purposes the measured points are shown each at the m_H value that would ideally correspond to it given the central value of m_t.

The situation is shown in Figure 2 [15]. The values of $(\sin^2\theta_e)_l$, $(\sin^2\theta_e)_h$ and their formal combination are shown each at the m_H value that would correspond to it given the central value of m_t. Of course, the value for m_H indicated by each $\sin^2\theta_{eff}$ has an horizontal ambiguity determined by the measurement error and the width of the $\pm 1\sigma$ band for m_t. Even taking this spread into account it is clear that the implications on m_H are sizably different.

One might imagine that some new physics effect could be hidden in the Zbb vertex. Like for the top quark mass there could be other non decoupling effects from new heavy states or a mixing of the b quark with some other heavy quark. However, it is well known that this discrepancy is not easily explained in terms of some new physics effect in the Zbb vertex. In fact, A^b_{FB} is the product of lepton- and b-asymmetry factors: $A^b_{FB} = (3/4)A_eA_b$. The sensitivity of A^b_{FB} to A_b is limited, because the A_e

factor is small, so that a rather large change of the b-quark couplings with respect to the SM is needed in order to reproduce the measured discrepancy (precisely a $\sim 30\%$ change in the right-handed coupling g_R^b, an effect too large to be a loop effect but which could be produced at the tree level, e.g., by mixing of the b quark with a new heavy vectorlike quark[16] or of the Z with an heavier Z'[17]). But this effect is not confirmed by the direct measurement of A_b performed at SLD using the left-right polarized b asymmetry, which agrees with the precision within the moderate precision of this result. Also, no deviation is manifest in the accurate measurement of $R_b \propto g_{Rb}^2 + g_{Lb}^2$ (but there g_R^b is not dominant). Thus, even introducing an ad hoc mixing the overall fit of A_{FB}^b, A_b and R_b is not terribly good, but we cannot exclude the possibility of new physics completely. Alternatively, the observed discrepancy could be due to a large statistical fluctuation or an unknown experimental problem. In any case the effective ambiguity in the measured value of $\sin^2\theta_{eff}$ is actually larger than the nominal error, reported in Eq. 3, obtained from averaging all the existing determinations.

We now discuss fitting the data in the SM. One can think of different types of fit, depending on which experimental results are included or which answers one wants to obtain. For example[1], in Table 1 we present in column 1 a fit of all Z pole data plus m_W, Γ_W (this is interesting as it shows the value of m_t obtained indirectly from radiative corrections, to be compared with the value of m_t measured in production experiments), in column 2 a fit of all Z pole data plus m_t (here it is m_W which is indirectly determined), and, finally, in column 3 a fit of all the data listed in Fig. 1 (which is the most relevant fit for constraining m_H). From the fit in column 1 of Table 1 we see that the extracted value of m_t is in good agreement with the direct measurement (see the value reported in Fig. 1). Similarly we see that the direct determination of m_W reported in Fig. 1 is still a bit larger with respect to the value from the fit in column 2 (although the direct value of m_W went down recently). We have seen that quantum corrections depend only logarithmically on m_H. In spite of this small sensitivity, the measurements are precise enough that one still obtains a quantitative indication of the Higgs mass range in the SM. From the fit in column 3 we obtain: $\log_{10} m_H(\mathrm{GeV}) = 1.94 \pm 0.16$ (or $m_H = 87^{+36}_{-27}$ GeV). We see that the central value of m_H from the fit is below the lower limit on the SM Higgs mass from direct searches $m_H > 114$ GeV, but within 1σ from this bound. If we had reasons to remove the result on A_{FB}^b from the fit, the fitted value of m_H would move down to something like: $m_H = 55^{+30}_{-20}$ GeV, further away from the lower limit.

Fit	1	2	3
Measurements	m_W	m_t	m_t, m_W
m_t (GeV)	178.7^{+12}_{-9}	172.6 ± 1.4	172.8 ± 1.4
m_H (GeV)	143^{+236}_{-80}	111^{+56}_{-39}	87^{+36}_{-27}
$\log[m_H(\mathrm{GeV})]$	2.16 ± 0.39	2.05 ± 0.18	1.94 ± 0.16
$\alpha_s(m_Z)$	0.1190 ± 0.0028	0.1190 ± 0.0027	0.1185 ± 0.0026
m_W (MeV)	80385 ± 21	80363 ± 20	80377 ± 15

We have already observed that the experimental value of m_W (with good agreement between LEP and the Tevatron) is a bit high compared to the SM prediction (see Figure 3,[15]). The value of m_H indicated by m_W is on the low side, just in the same interval as for $\sin^2\theta_e^{lept}$ measured from leptonic asymmetries. The recent decrease of the experimental value of m_t maintains the tension between the experimental values of m_W and $\sin^2\theta_e^{lept}$ measured from leptonic asymmetries on the one side and the lower limit on m_H from direct searches on the other side [18,19].

With all these words of caution in mind it remains true that on the whole the SM performs rather well, so that it is fair to say that no clear indication for new physics emerges from the data. Actually the result of precision tests on the Higgs mass is particularly remarkable. The value of $\log_{10}[m_H(\mathrm{GeV})]$ is, within errors, inside the small window between ~ 2 and ~ 3 which is allowed, on the one side, by the direct search limit ($m_H > 114.4$ GeV from LEP-2 [1]), and, on the other side, by the theoretical upper limit on the Higgs mass in the minimal SM [6], $m_H < 600 - 800$ GeV.

Thus the whole picture of a perturbative theory with a fundamental Higgs is well supported by the data on radiative corrections. It is important that there is a clear indication for a particularly light Higgs: at 95% c.l. $m_H < 190$ GeV. This is quite encouraging for the ongoing search for the Higgs particle. More in general, if the Higgs couplings are removed from the Lagrangian the resulting theory is non renormalizable. A cutoff Λ must be introduced. In the quantum corrections $\log m_H$ is then replaced by $\log\Lambda$ plus a constant. The precise determination of the associated finite terms would be lost (that is, the value of the mass in the denominator in the argument of the logarithm). A heavy Higgs would need some conspiracy or some dynamical reason[20]: the finite terms, different in the new theory from

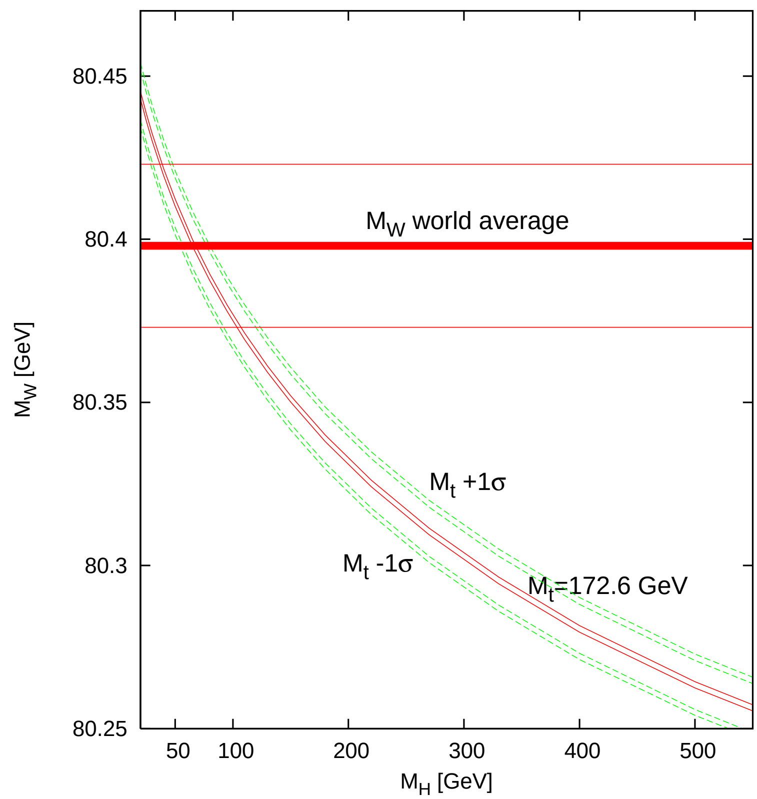

Figure 3. The world average for m_W is plotted vs m_H.

those of the SM, should accidentally or dynamically compensate for the heavy Higgs in a few key parameters of the radiative corrections (mainly ϵ_1 and ϵ_3, see, for example, [21]). Alternatively, additional new physics, for example in the form of effective contact terms added to the minimal SM lagrangian, should do the compensation, which again needs some sort of conspiracy or some special dynamics, although this possibility is not so unlikely to be apriori discarded.

5. The physics of flavour

In the last decade great progress in different areas of flavour physics has been achieved. In the quark sector, the amazing results of a generation of frontier experiments, obtained at B factories and at accelerators, have become available[22]. QCD has been playing a crucial role in the interpretation of experiments by a combination of effective theory methods (heavy quark effective theory, NRQCD, SCET), lattice simulations and perturbative calculations. A great achievement obtained by many theorists over the last years is the calculation at NNLO of the branching ratio for $B \to X_s \gamma$ with B a beauty meson[24]. The effect of the photon energy cut, $E_\gamma > E_0$, necessary in practice, has been evaluated at NNLO [25]. The central value of the theoretical prediction is now slightly below the data: for $\mathcal{B}[B \to X_s \gamma; E_0 = 1.6$ GeV$](10^{-4})$ the experimental value is 3.55(26)[23] and the theoretical value is 3.15(23)[24] or 2.98(26)[25], which to me is good agreement. The hope of the B-decay experiments was to detect departures from the CKM picture of mixing and of CP violation as signals of new physics. Finally, in quantitative terms, all measurements are in agreement with the CKM description of mixing and CP violation as shown in Fig. 4[26]. The recent measurement of Δm_s by CDF and D0, in fair agreement with the SM expectation, has closed another door for new physics. But in some channels, especially those which occur through penguin loops, it is well possible that substantial deviations could be hidden (possible hints are reported in $B \to K\pi$ decays[27] and in $b \to s$ transitions[28]). But certainly the amazing performance of the SM in flavour changing and/or CP violating transitions in K and B decays poses very strong constraints on all proposed models of new physics[29].

In the leptonic sector the study of neutrino oscillations has led to the discovery that at least two neutrinos are not massless and to the determination of the mixing matrix[30]. Neutrinos are not all massless but their masses are very small (at most a fraction of eV). Probably masses are small because ν's are Majorana fermions, and, by the see-saw mechanism, their masses are inversely proportional to the large scale M where lepton number (L) non conservation occurs (as expected in GUT's). Indeed the value of $M \sim m_{\nu_R}$ from experiment is compatible with being close to $M_{GUT} \sim 10^{14} - 10^{15}$ GeV, so that neutrino masses fit well in the GUT picture and actually support it. The interpretation of neutrinos as Majorana particles enhances the importance of experiments aimed at the detection of neutrinoless double beta decay and a huge effort in this direction is underway[31]. It was realized that decays of heavy ν_R with CP and L non conservation can produce a B-L asymmetry. The range of neutrino masses indicated by neutrino phenomenology turns out to be perfectly compatible with the idea of baryogenesis via leptogenesis[32]. This elegant model for baryogenesis has by now replaced the idea of baryogenesis near the weak

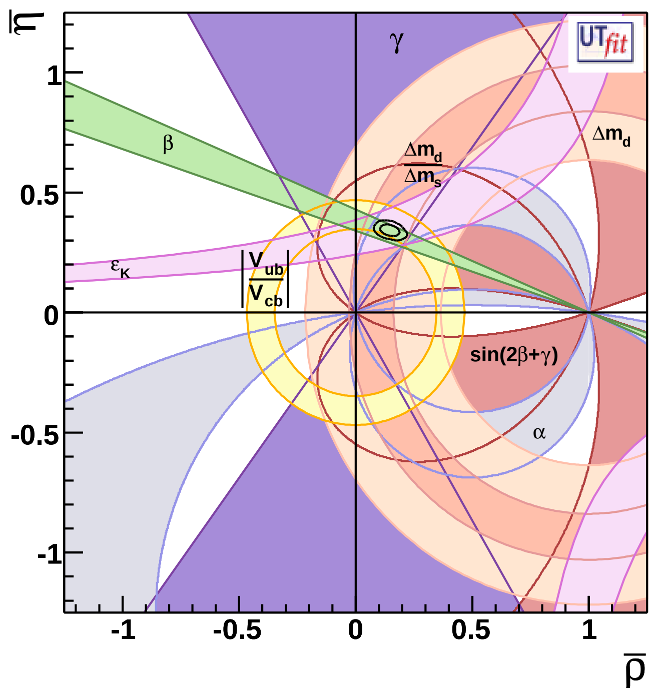

Figure 4. Constraints in the ; plane including the most recent , and M $_s$ inputs in the global CKM t.

scale, which has been strongly disfavoured by LEP.

It is remarkable that we now know the neutrino mixing matrix with good accuracy. Two mixing angles are large and one is small. The atmospheric angle $_{23}$ is large, actually compatible with maximal but not necessarily so: at 3 [33]: 0:34 sin^2 $_{23}$ 0:68 with central value around 0:5. The solar angle $_{12}$ (the best measured) is large, sin^2 $_{12}$ 0:3, but certainly not maximal (by more than 5). The third angle $_{13}$, strongly limited mainly by the CHOOZ experiment, has at present a 3 upper limit given by about sin^2 $_{13}$ 0:04. The non conservation of the three separate lepton numbers and the large leptonic mixing angles make it possible that processes like ! e or ! might be observable, not in the SM but in extensions of it like the MSSM. Thus, for example, the outcome of the now running experiment MEG at PSI [34], aiming at improving the limit on ! e by 1 or 2 orders of magnitude, is of great interest.

6. Problems of the Standard Model

No signals of new physics were found neither in electroweak precision tests nor in flavour physics. Given the success of the SM why are we not satisfied with that theory? Why not just find the Higgs particle, for completeness, and declare that particle physics is closed? The reason is that there are both conceptual problems and phenomenological indications for physics beyond the SM. On the conceptual side the most obvious problems are the proliferation of parameters, the puzzles of family replication and of flavour hierarchies, the fact that quantum gravity is not included in the SM and the related hierarchy problem. Some of these problems could be postponed to the more fundamental theory at the Planck mass. For example, the explanation of the three generations of fermions and the understanding of fermion masses and mixing angles can be postponed. But other problems, like the hierarchy problem, must find their solution in the low energy theory. Among the main phenomenological hints for new physics we can list the quest for Grand Unification and coupling constant merging, dark matter, neutrino masses (explained in terms of L non conservation), baryogenesis and the cosmological vacuum energy (a gigantic naturalness problem).

6.1. Dark matter and dark energy

We know by now[35] that the Universe is flat and most of it is not made up of known forms of matter: while $\Omega_{tot} \sim 1$ and $\Omega_{matter} \sim 0.3$, the normal baryonic matter is only $\Omega_{baryonic} \sim 0.044$, where Ω is the ratio of the density to the critical density. Most of the energy in the Universe is Dark Matter (DM) and Dark Energy (DE) with $\Omega_\Lambda \sim 0.7$. We also know that most of DM must be cold (non relativistic at freeze-out) and that significant fractions of hot DM are excluded. Neutrinos are hot DM (because they are ultrarelativistic at freeze-out) and indeed are not much cosmo-relevant: $\Omega_\nu \lesssim 0.015$. The identification of DM is a task of enormous importance for both particle physics and cosmology. The LHC has good chances to solve this problem in that it is sensitive to a large variety of WIMP's (Weekly Interacting Massive Particles). WIMP's with masses in the 10 GeV–1TeV range with typical EW cross-sections turn out to contribute terms of o(1) to Ω. This is a formidable hint in favour of WIMP's as DM candidates. By comparison, axions are also DM candidates but their mass and couplings must be tuned for this purpose. If really some sort of WIMP's are a main component of DM they could be discovered at the LHC and this will be a great service of particle physics to cosmology. Also, we have seen that vacuum energy accounts for about 2/3 of the critical density: ~ 0.7[36]. Translated into familiar units this means for the energy density $\sim (2\ 10^{-3}\ \mathrm{eV})^4$ or $(0.1\ \mathrm{mm})^{-4}$. It is really interesting (and not at all understood) that $^{1/4}$ $\sim$ $^2_{EW}/M_{Pl}$ (close to the range of neutrino masses). It is well known that in eld theory we expect $\sim$ $^4_{cutoff}$. If the cut o is set at M_{Pl} or even at 0(1 TeV) there would be an enormous mismatch. In exact SUSY = 0, but SUSY is broken and in presence of breaking $^{1/4}$ is in general not smaller than the typical SUSY multiplet splitting. Another closely related problem is "why now?": the time evolution of the matter or radiation density is quite rapid, while the density for a cosmological constant term would be at in time. If so, then how comes that precisely now the two density sources are comparable? This suggests that the vacuum energy is not a cosmological constant term, but rather the vacuum expectation value of some eld (quintessence) and that the "why now?" problem is solved by some dynamical coupling of the quintessence eld with gauge singlet elds (perhaps RH neutrinos)[37].

6.2. The hierarchy problem

The computed evolution with energy of the e ective gauge couplings clearly points towards the uni cation of the EW and strong forces (Grand Uni ed Theories: GUT's) at scales of energy $M_{GUT} \sim 10^{15} - 10^{16}$ GeV which are close to the scale of quantum gravity, $M_{Pl} \sim 10^{19}$ GeV. GUT's are so attractive that are by now part of our culture: they provide coupling uni cation, an explanation of the quantum numbers in each generation of fermions (e.g. one generation exactly lls the 16 dimensional representation of SO(10)), transformation of quarks into leptons and proton decay etc. One step further and one is led to imagine a uni ed theory of all interactions also including gravity (at present superstrings provide the best attempt at such a theory). Thus GUT's and the realm of quantum gravity set a very distant energy horizon that modern particle theory cannot ignore. Can the SM without new physics be valid up to such large energies? The answer is presumably not: the structure of the SM could not naturally explain the relative smallness of the weak scale of mass, set by the Higgs mechanism at $\sim 1/\sqrt{G_F} \sim 250$ GeV with G_F being the Fermi coupling constant, with respect to M_{GUT} or M_{Pl}. This so-called hierarchy problem is due to the instability of the SM with respect to quantum corrections. This is related to the presence of fundamental scalar elds in the theory with quadratic mass divergences and no protective extra symmetry at = 0. For fermion masses, first, the divergences are logarithmic and, second, they are forbidden by the $SU(2) \bigotimes U(1)$ gauge symmetry plus the fact that at $m = 0$ an additional symmetry, i.e. chiral symmetry, is restored. Here, when talking of divergences, we are not worried of actual infinities. The theory is renormalizable and finite once the dependence on the cut off Λ is absorbed in a redefinition of masses and couplings. Rather the hierarchy problem is one of naturalness. We can look at the cut off as a parameterization of our ignorance on the new physics that will modify the theory at large energy scales. Then it is relevant to look at the dependence of physical quantities on the cut off and to demand that no unexplained enormously accurate cancellations arise.

In the past in many cases naturalness has been a good guide in particle physics. For example, without charm and the GIM mechanism the short distance contribution to the $(K - \bar{K})$ mass difference would be of order $G_F^2 f_K^2 m_W^2 m_K$, while the correct result is of order $G_F^2 f_K^2 m_c^2 m_K$ and, without GIM, an unnatural cancellation between long and short distance contributions would be needed. Also note that $\Lambda_{QCD} << M_{GUT}$ is natural because, due to the logarithmic running of α_s, dimensional transmutation brings in exponential suppression.

The hierarchy problem can be put in less abstract terms (the "little hierarchy problem"): loop corrections to the higgs mass squared are quadratic in the cut off Λ. The most pressing problem is from the top loop. With $m_H^2 = m_{bare}^2 + \delta m_H^2$ the top loop gives

$$\delta m^2_{H|top} \sim -\frac{3G_F}{2\sqrt{2}\pi^2} m_t^2 \Lambda^2 \sim -(0.2\Lambda)^2 \qquad (4)$$

If we demand that the correction does not exceed the light Higgs mass indicated by the precision tests, Λ must be close, $\Lambda \sim o(1\ \text{TeV})$. Similar constraints arise from the quadratic Λ dependence of loops with gauge bosons and scalars, which, however, lead to less pressing bounds. So the hierarchy problem demands new physics to be very close (in particular the mechanism that quenches the top loop). Actually, this new physics must be rather special, because it must be very close, yet its effects are not clearly visible in the EW precision tests (the "LEP Paradox"[38]) now also accompanied by a similar "flavour paradox"[29] arising from the recent precise experimental results in B and K decays. The main avenues open for new physics are discussed in the following sections[39].

7. Supersymmetry: the standard way beyond the SM

Models based on supersymmetry (SUSY)[40] are the most developed and widely known. In the limit of exact boson-fermion symmetry the quadratic divergences of bosons cancel, so that only logarithmic divergences remain. However, exact SUSY is clearly unrealistic. For approximate SUSY (with soft breaking terms), which is the basis for all practical models, Λ in eq.(4) is essentially replaced by the splitting of SUSY multiplets. In particular, the top loop is quenched by partial cancellation with s-top exchange, so the s-top cannot be too heavy.

The Minimal SUSY Model (MSSM) is the extension of the SM with minimal particle content. To each ordinary particle a s-particle is associated with 1/2 spin difference: to each helicity state of a spin 1/2 fermion of the SM a scalar is associated (for example, the electron states e_L and e_R correspond to 2 scalar s-electron states). Similarly to each ordinary boson a s-fermion is associated: for example to each gluon a gluino (a Majorana spin 1/2 fermion) is related. Why not even one s-particle was seen so far? A clue: observed particles are those whose mass is forbidden by $SU(2) \bigotimes U(1)$. When SUSY is broken but $SU(2) \bigotimes U(1)$ is unbroken s-particles get a mass but particles remain massless. Thus if SUSY breaking is large we understand that no s-particles have been observed yet. It is an important fact that two Higgs doublets, H_u and H_d, are needed in the MSSM with their corresponding spin 1/2 s-partners, to give mass to the up-type and to the down-type fermions, respectively. This duplication is needed for cancellation of the chiral anomaly and also because the SUSY rules forbid that $H_d = H_u^\dagger$ as is the case in the the SM. The ratio of their two vacuum expectation values $\tan\beta = v_u/v_d$ (with the SM vev v being given by $v = \sqrt{v_u^2 + v_d^2}$) plays an important role for phenomenology.

The most general MSSM symmetric renormalizable lagrangian would contain terms that violate baryon B and lepton L number conservation (which in the SM, without ν_R, are preserved at the renormalizable level, so that they are "accidental" symmetries). To eliminate those terms it is sufficient to invoke a discrete parity, R-parity, whose origin is assumed to be at a more fundamental level, which is $+1$ for ordinary particles and -1 for s-partners. The consequences of R-parity are that s-particles are produced in pairs at colliders, the lightest s-particle is absolutely stable (it is called the Lightest SUSY Particle, LSP, and is a good candidate for dark matter) and s-particles decay into a final state with an odd number of s-particles (and, ultimately, in the decay chain there will be the LSP).

The necessary SUSY breaking, whose origin is not clear, can be phenomenologically introduced through soft terms (i. e. with operator dimension < 4) that do not spoil the good convergence properties of the theory (renormalizability and non renormalization theorems are maintained). We denote by m_{soft} the mass scale of the soft SUSY breaking terms. The most general soft terms compatible with the SM gauge symmetry and with R-parity conservation introduce more than one hundred new parameters. In general new sources of flavour changing neutral currents (FCNC) and of CP violation are introduced e.g. from s-quark mass matrices. Universality (proportionality of the mass matrix to the identity matrix for each charge sector) and/or alignment (near diagonal mass matrices) must be assumed at a large scale, but renormalization group running can still produce large effects. The MSSM does provide a viable flavour framework in the assumption of R-parity conservation, universality of soft masses and proportionality of trilinear soft terms to the SM Yukawas (still broken by renormalization group running). As already mentioned, observable effects in the lepton sector are still possible (e.g. $\mu \to e\gamma$ or $\tau \to \mu\gamma$). This is made even more plausible by large neutrino mixings.

How can SUSY breaking be generated? Conventional spontaneous symmetry breaking cannot occur within the MSSM and also in simple extensions of it. Probably the soft terms of the MSSM arise indirectly or radiatively (loops) rather than from tree level renormalizable couplings. The prevailing idea is that it happens in a "hidden sector" through non renormalizable interactions and is communicated to the visible sector by some interactions. Gravity is a plausible candidate for the hidden sector. Many theorists consider SUSY as established at the Planck scale M_{Pl}. So why not to use it also at low energy to fix the hierarchy problem, if at all possible? It is interesting that viable models exist. Suitable soft terms indeed arise from supergravity when it is spontaneoulsly broken. Supergravity is a non renormalizable SUSY theory of quantum gravity[40]. The SUSY partner of the spin-2 graviton $g_{\mu\nu}$ is the spin-3/2 gravitino ψ_{μ}^{i} (i: spinor index, μ: Lorentz index). The gravitino is the gauge field associated to the SUSY generator. When SUSY is broken the gravitino takes mass by absorbing the 2 goldstino components (super-Higgs mechanism). In gravity mediated SUSY breaking typically the gravitino mass $m_{3/2}$ is of order m_{soft} (the scale of mass of the soft breaking terms) and, on dimensional ground, both are given by $m_{3/2} \sim m_{soft} \sim \langle F\rangle/M_{Pl}$, where F is the dimension 2 auxiliary field that takes a vacuum expectation value $\langle F\rangle$ in the hidden sector (the denominator M_{Pl} arises from the gravitational coupling that transmits the breaking down to the visible sector). For $m_{soft} \sim 1$ TeV, the scale of SUSY breaking is very large of order $\sqrt{\langle F\rangle} \sim \sqrt{m_{soft} M_{Pl}} \sim 10^{11}$ GeV. With TeV mass and gravitational coupling the gravitino is not relevant for LHC physics but perhaps for cosmology (it could be the LSP and a dark matter candidate). In gravity mediation the neutralino is the typical LSP and an excellent dark matter candidate. A lot of missing energy is a signature for gravity mediation.

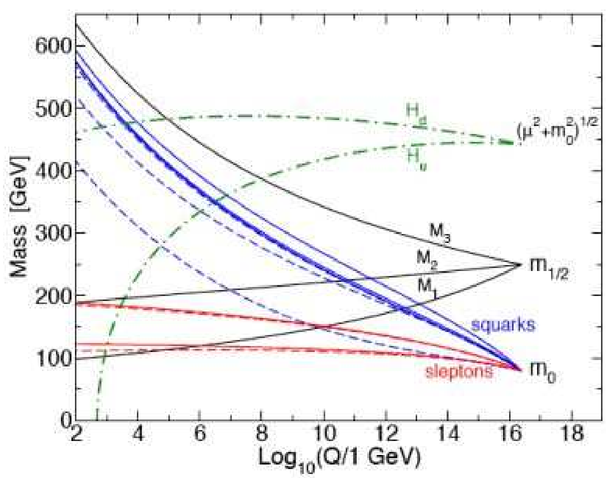

Figure 5. A SUSY spectrum generated by universal boundary conditions at the GUT scale

Different mechanisms of SUSY breaking are also being considered. In one alternative scenario[43] the (not so much) hidden sector is connected to the visible one by messenger heavy fields, with mass M_{mess}, which share ordinary gauge interactions and thus, in amplitudes involving only external light particles, appear in loops so that $m_{soft} \sim \frac{\alpha_i}{4\pi}\frac{\langle F\rangle}{M_{mess}}$. Both gaugino and s-fermion masses are of order m_{soft}. Messengers can be taken in complete SU(5) representations, like $5+\bar{5}$, so that coupling unification is not spoiled. As gauge interactions are much stronger than gravitational interactions, the SUSY breaking scale can be much smaller, as low as $\sqrt{\langle F\rangle} \sim M_{mess} \sim 10-100$ TeV. It follows that the gravitino is very light (with mass of order or below 1 eV typically) and, in these models, always is the LSP. Its couplings are observably large because the gravitino couples to

SUSY particle multiplets through its spin 1/2 goldstino components. Any SUSY particle will eventually decay into the gravitino. But the decay of the next-to-the lightest SUSY particle (NLSP) could be extremely slow, with a travel path at the LHC from microscopic to astronomical distances. The main appeal of gauge mediated models is a better protection against FCNC: if one starts at M_{mess} with sufficient universality/alignment then the very limited interval for renormalization group running down to the EW scale does not spoil it. Indeed at M_{mess} there is approximate alignment because the mixing parameters $A_{u,d,l}$ in the soft breaking lagrangian are of dimension of mass and arise at two loops, so that they are suppressed.

What is unique to SUSY with respect to most other extensions of the SM is that SUSY models are well defined and computable up to M_{Pl} and, moreover, are not only compatible but actually quantitatively supported by coupling unification and GUT's. At present the most direct phenomenological evidence in favour of SUSY is obtained from the unification of couplings in GUT's. Precise LEP data on $\alpha_s(m_Z)$ and $\sin^2\theta_W$ show that standard one-scale GUT's fail in predicting $\alpha_s(m_Z)$ given $\sin^2\theta_W$ and $\alpha(m_Z)$ while SUSY GUT's are compatible with the present, very precise, experimental results (of course, the ambiguities in the MSSM prediction are larger than for the SM case because of our ignorance of the SUSY spectrum). If one starts from the known values of $\sin^2\theta_W$ and $\alpha(m_Z)$, one finds[44] for $\alpha_s(m_Z)$ the results: $\alpha_s(m_Z) = 0.073 \pm 0.002$ for Standard GUT's and $\alpha_s(m_Z) = 0.129 \pm 0.010$ for SUSY GUT's to be compared with the world average experimental value $\alpha_s(m_Z) = 0.118 \pm 0.002$[45]. Another great asset of SUSY GUT's is that proton decay is much slowed down with respect to the non SUSY case. First, the unification mass $M_{GUT} \sim$ few 10^{16} GeV, in typical SUSY GUT's, is about 20 times larger than for ordinary GUT's. This makes p decay via gauge boson exchange negligible and the main decay amplitude arises from dim-5 operators with higgsino exchange, leading to a rate close but still compatible with existing bounds (see, for example,[46]).

By imposing on the MSSM model universality constraints at M_{GUT} one obtains a drastic reduction in the number of parameters at the price of more rigidity and model dependence (see Figure 5[40]). This is the SUGRA or CMSSM (C for "constrained") limit[40]. An interesting exercise is to repeat the fit of precision tests in the CMSSM, also including the additional data on the muon $(g-2)$, the dark matter relic density and the $b \to s\gamma$ rate. The result[47] is that the central value of the lightest Higgs mass m_h goes up (in better harmony with the bound from direct searches) with moderately large $\tan\beta$ and relatively light SUSY spectrum.

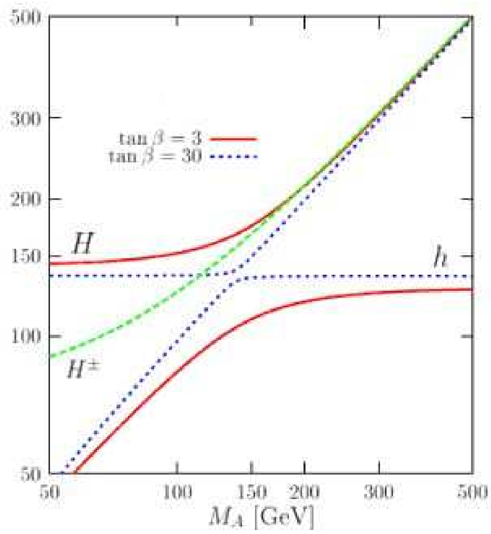

Figure 6. The MSSM Higgs spectrum as function of m_A: h is the lightest Higgs, H and A are the heavier neutral scalar and pseudoscalar Higgs, respectively, and $H^{\pm}$ are the charged Higgs bosons. The curves refer to $m_t = 178$ GeV and large top mixing A_t

In spite of all these virtues it is true that the lack of SUSY signals at LEP and the lower limit on m_H pose problems for the MSSM. The predicted spectrum of Higgs particles in the MSSM is shown in Figure 6[48]. As apparent from the figure the lightest Higgs particle is predicted in the MSSM to be below $m_h \lesssim 130$ GeV (with the esperimental value of m_t going down the upper limit is slightly decreased). In fact, at tree level $m_h^2 = m_Z^2 \cos^2 2\beta$ and it is only through radiative corrections that m_h can increase beyond m_Z:

$$m_h^2 \sim \frac{3G_F}{\sqrt{2}\pi^2} m_t^4 \log \frac{m_{\tilde{t}_1} m_{\tilde{t}_2}}{m_t^2} \tag{5}$$

Here $\tilde{t}_{1,2}$ are the s-top mass eigenstates. The direct limit on m_h from the Higgs search at LEP, shown in Figure 7[49], considerably restricts the available parameter space of the MSSM requiring relatively large $\tan\beta$ and heavy s-top quarks. Stringent naturality constraints also follow from imposing that the EW breaking occurs at the right energy scale: in SUSY models the breaking is induced by the running of the H_u mass starting from a common scalar mass m_0 at M_{GUT} (see Figure 5). The squared Z mass m_Z^2 can be expressed as a linear combination of the SUSY param-

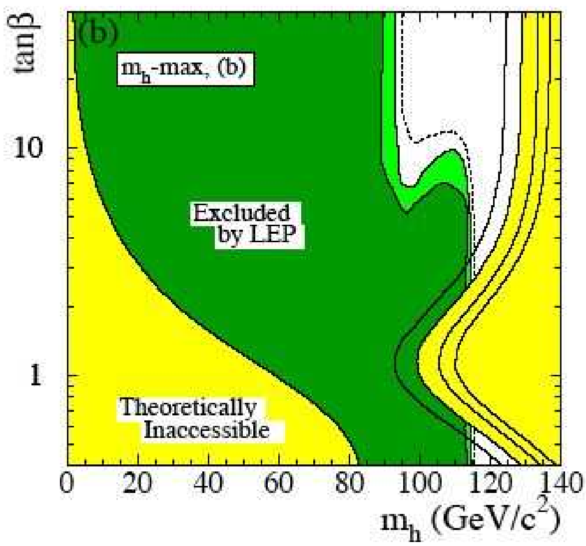

Figure 7. Experimental limits in the $\tan\beta - m_h$ plane from LEP. With h one denotes the lightest MSSM Higgs boson.

eters $m_0^2, m_{1/2}^2, A_t^2, \mu^2$,... with known coefficients. Barring cancellations that need fine tuning, the SUSY parameters, hence the SUSY s-partners, cannot be too heavy. The LEP limits, in particular the chargino lower bound $m_{\chi^+} > 100$ GeV, are sufficient to eliminate an important region of the parameter space, depending on the amount of allowed fine tuning. For example, models based on gaugino universality at the GUT scale, like the CMSSM, need a fine tuning by at least a factor of 20. Without gaugino universality[51] the strongest limit remains on the gluino mass: the relation reads $m_Z^2 \sim 0.7\, m_{gluino}^2 + \ldots$ and is still compatible with the present limit $m_{gluino} > 250 - 300$ GeV from the TeVatron (see Figure 8[50])

This is the case of the MSSM with minimal particle content. Of course, minimality is only a simplicity assumption that could possibly be relaxed. For example, adding an additional Higgs singlet S considerably helps in addressing naturalness constraints (Next-to Minimal SUSY SM: NMSSM)[41,42]. An additional singlet can also help solving the "μ-problem"[40]. In the exact SUSY and gauge symmetric limit there is a single parameter with dimension of mass in the superpotential. The μ term in the superpotential is of the form $W_{\mu\, term} = \mu H_u H_d$. The mass μ, which

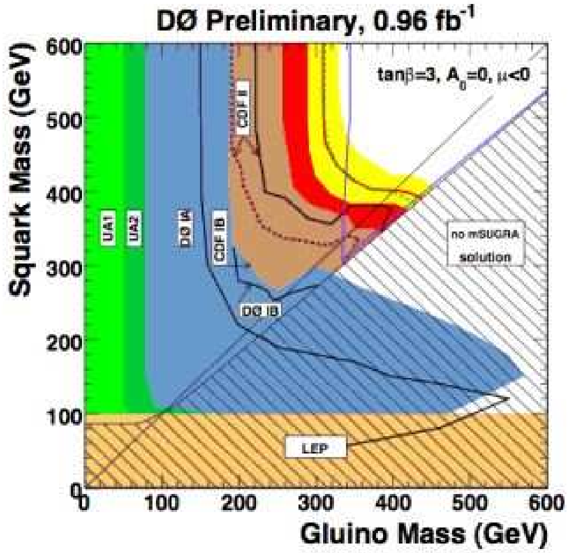

Figure 8. Present experimental limits on s-quarks and gluinos

contributes to the Higgs sector masses, must be of order m_{soft} for phenomenological reasons. The problem is to justify this coincidence, because could in principle be much larger given that it already appears at the symmetric level. A possibility is to forbid the term by a suitable symmetry in the SUSY unbroken limit and then generate it together with the SUSY breaking terms. For example, one can introduce a discrete parity that forbids the term. Then Giudice and Masiero[52] have observed that in general, the low energy limit of supergravity, also induces a SUSY conserving term together with the soft SUSY breaking terms and of the same order. A di erent phenomenologically appealing possibility is to replace with the vev of a new singlet scalar eld S, thus enlarging the Higgs sector as in the NMSSM.

In conclusion the main SUSY virtues are that the hierarchy problem is drastically reduced, the model agrees with the EW data, is consistent and computable up to M_{Pl}, is well compatible and indeed supported by GUT's, has good dark matter candidates and, last not least, is testable at the LHC. The delicate points for SUSY are the origin of SUSY breaking and of R-parity, the -problem, the avour problem and the need of sizable ne tuning.

8. Little Higgs models

The non discovery of SUSY at LEP has given further impulse to the quest for new ideas on physics beyond the SM. In "little Higgs" models the symmetry of the SM is extended to a suitable global group G that also contains some gauge enlargement of $SU(2) \otimes U(1)$, for example $G \quad [SU(2) \otimes U(1)]^2 \quad SU(2) \otimes U(1)$. The Higgs particle is a pseudo-Goldstone boson of G that only takes mass at 2-loop level, because two distinct symmetries must be simultaneously broken for it to take mass, which requires the action of two di erent couplings in the same diagram. Then in the relation eq.(4) between m_h^2 and ^{2} there is an additional coupling and an additional loop factor that allow for a bigger separation between the Higgs mass and the cut-o . Typically, in these models one has one or more Higgs doublets at $m_h \quad 0.2$ TeV, and a cut-o at 10 TeV. The top loop quadratic cut-o dependence is partially canceled, in a natural way guaranteed by the symmetries of the model, by a new coloured, charge 2/3, vectorlike quark of mass around 1 TeV (a fermion not a scalar like the s-top of SUSY models). Certainly these models involve a remarkable level of group theoretic virtuosity. However, in the simplest versions one is faced with problems with precision tests of the SM [71]. These problems can be xed by complicating the model[54]: one can introduce a parity symmetry, T-parity, and additional "mirror" fermions. T-parity interchanges the two $SU(2) \otimes U(1)$ groups: standard gauge bosons are T even while heavy ones are T odd. As a consequence no tree level contributions from heavy W and Z appear in processes with external SM particles. Therefore all corrections to EW observables only arise at loop level. A good feature of T-parity is that, like for R-parity in the MSSM, the lightest T-odd particle is stable (usually a B') and can be a candidate for Dark Matter (missing energy would here too be a signal) and T-odd particles are produced in pairs (unless T-parity is not broken by anomalies[55]). Thus the model could work but, in my opinion, the real limit of this approach is that it only o ers a postponement of the main problem by a few TeV, paid by a complete loss of predictivity at higher energies. In particular all connections to GUT's are lost. Still it is very useful to o er to experiment a di erent example of possible new physics.

9. Extra dimensions

Extra dimensions models are among the most interesting new directions in model building. Early formulations were based on "large" extra dimensions [56,57]. These are models with factorized metric: $ds^2 = \eta_{\mu\nu} dx^\mu dx^\nu + h_{ij}(y)dy^i dy^j$, where $y^{i,j}$ denote the extra dimension coordinates and indices. Large extra dimension models propose to solve the hierarchy problem by bringing gravity down from M_{Pl} to $m \sim o(1\ \mathrm{TeV})$ where m is the string scale. Inspired by string theory one assumes that some compactified extra dimensions are sufficiently large and that the SM fields are confined to a 4-dimensional brane immersed in a d-dimensional bulk while gravity, which feels the whole geometry, propagates in the bulk. We know that the Planck mass is large just because gravity is weak: in fact $G_N \sim 1/M_{Pl}^2$, where G_N is Newton constant. The new idea is that gravity appears so weak because a lot of lines of force escape in extra dimensions. Assume you have $n = d - 4$ extra dimensions with compactification radius R. For large distances, $r >> R$, the ordinary Newton law applies for gravity: in natural units, the force between two units of mass is $F \sim G_N/r^2 \sim 1/(M_{Pl}^2 r^2)$. At short distances, $r \lesssim R$, the flow of lines of force in extra dimensions modifies Gauss law and $F^{-1} \sim m^2 (mr)^{d-4} r^2$. By matching the two formulas at $r = R$ one obtains $(M_{Pl}/m)^2 = (Rm)^{d-4}$. For $m \sim 1$ TeV and $n = d - 4$ one finds that $n = 1$ is excluded ($R \sim 10^{15}$ cm), for $n = 2$ R is very marginal and also at the edge of present bounds $R \sim 1$ mm on departures from Newton law [58], while for $n = 4, 6$, $R \sim 10^{-9}, 10^{-12}$ cm and these cases are not excluded.

A generic feature of extra dimensional models is the occurrence of Kaluza-Klein (KK) modes. Compactified dimensions with periodic boundary conditions, like the case of quantization in a box, imply a discrete spectrum with momentum $p = n/R$ and mass squared $m^2 = n^2/R^2$. In any case there are the towers of KK recurrences of the graviton. They are gravitationally coupled but there are a lot of them that sizably couple, so that the net result is a modification of cross-sections and the presence of missing energy. There are many versions of these models. The SM brane can itself have a thickness r with $r \lesssim 10^{-17}$ cm or $1/r \gtrsim 1$ TeV, because we know that quarks and leptons are pointlike down to these distances, while for gravity in the bulk there is no experimental counter-evidence down to $R \lesssim 0.1$ mm or $1/R \gtrsim 10^{-3}$ eV. In case of a thickness for the SM brane there would be KK recurrences for SM fields, like W_n, Z_n and so on in the TeV region and above. Large extra dimensions provide an exciting scenario.

Already it is remarkable that this possibility is compatible with experiment. However, there are a number of criticisms that can be brought up. First, the hierarchy problem is more translated in new terms rather than solved. In fact the basic relation $Rm = (M_{Pl}/m)^{2/n}$ shows that Rm, which one would apriori expect to be 0(1), is instead ad hoc related to the large ratio M_{Pl}/m. Also it is not clear how extra dimensions can by themselves solve the LEP paradox (the large top loop corrections should be controlled by the opening of the new dimensions and the onset of gravity): since m_H is light $\Lambda \sim 1/R$ must be relatively close. But precision tests put very strong limits on Λ. In fact in typical models of this class there is no mechanism to sufficiently quench the corrections.

More recently models based on the Randall-Sundrum (RS) solution for the metric have attracted most of the model builders attention[59,60]. In these models the metric is not factorized and an exponential "warp" factor multiplies the ordinary 4-dimensional coordinates in the metric: $ds^2 = e^{-2kR\phi}\eta_{\mu\nu}dx^\mu dx^\nu - R^2\phi^2$ where ϕ is the extra coordinate. This non-factorizable metric is a solution of Einstein equations with specified 5-dimensional cosmological term. Two 4-dimensional branes are often localized at $\phi = 0$ (the Planck or ultraviolet brane) and at $\phi = \pi$ (the infrared brane). In the simplest models all SM fields are located on the infrared brane. All 4-dim masses m_4 are scaled down with respect to 5-dimensional masses $m_5 \sim k \sim M_{Pl}$ by the warp factor: $m_4 = M_{Pl}e^{-kR\pi}$. In other words mass and energies on the infrared brane are redshifted by the $\sqrt{g_{00}}$ factor. The hierarchy suppression m_W/M_{Pl} could arise from the warping exponential $e^{-kR\pi}$, for not too large values of the warp factor exponent: $kR \sim 12$ (extra dimension are not "large" in this case). The question of whether these values of kR can be stabilized has been discussed in ref.[61]. It was shown that the determination of kR at a compatible value can be assured by a scalar field in the bulk ("radion") with a suitable potential which offer the best support to the solution of the hierarchy problem in this context. In the original RS models where the SM fields are on the brane and gravity is in the bulk there is a tower of spin-2 KK graviton resonances. Their couplings to ordinary particles are of EW order (because their propagator masses are red shifted on the infrared brane) and universal for all particles. These resonances could be visible at the LHC. Their signature is spin-2 angular distributions and universality of couplings. The RS original formulation is very elegant but when going to a realistic formulation it has problems, for example with EW precision tests. Also, In a description of physics from m_W to M_{Pl} there should be place for GUTs. But, if all SM particles are on the TeV brane the effective theory cut-off is low and no way to M_{GUT} is open. Inspired by RS different realizations of warped geometry were tried: gauge fields in the bulk and/or all SM fields (except the Higgs) on the bulk. The hierarchy of fermion masses can be seen as the result of the different profiles of the corresponding distributions in the bulk: the heaviest fermions are those closest to the brane where the Higgs is located. While no simple, realistic model has yet emerged as a benchmark, it is attractive to imagine that ED could be a part of the truth, perhaps coupled with some additional symmetry or even SUSY.

Extra dimensions offer new possibilities for SUSY breaking. In fact, ED can realize a geometric separation between the hidden (on the Planck brane) and the visible sector (on the TeV brane), with gravity mediation in the bulk. In anomaly mediated SUSY breaking[62] 5-dim quantum gravity effects act as messengers. The name comes because L_{soft} can be understood in terms of the anomalous violation of a local superconformal invariance. In a particular formulation of 5 dimensional supergravity, at the classical level, the soft term are exponentially suppressed on the MSSM brane. SUSY breaking effects only arise at quantum level through beta functions and anomalous dimensions of the brane couplings and fields. In this case gaugino masses are proportional to gauge coupling beta functions, so that the gluino is much heavier than the electroweak gauginos.

In the general context of extra dimensions an interesting direction of development is the study of symmetry breaking by orbifolding and/or boundary conditions. Orbifolding means that we have a 5 (or more) dimensional theory where the extra dimension $x_5 = y$ is compactified. Along y one or more Z_2 reflections are defined, for example $P = y \leftrightarrow -y$ (a reflection around the horizontal diameter) and $P^0 = y \leftrightarrow -y - \pi R$ (a reflection around the vertical diameter). A field $\phi(x;y)$ with definite P and P^0 parities can be Fourier expanded along y. Then $\phi_{++};\phi_{+-};\phi_{-+};\phi_{--}$ have the n-th Fourier components proportional to $\cos\frac{2ny}{R};\cos\frac{(2n+1)y}{R};\sin\frac{(2n+1)y}{R};\sin\frac{(2n+2)y}{R}$, respectively. On the branes located at the fixed points of P and P^0, $y = 0$ and $y = -\pi R/2$, the symmetry is reduced: indeed at $y = 0$ only ϕ_{++} and ϕ_{+-} are non vanishing and only ϕ_{++} is massless.

For example, at the GUT scale, symmetry breaking by orbifolding can be applied to obtain a reformulation of SUSY GUT's where many problematic features of ordinary GUT's (e.g. a baroque Higgs sector, the doublet-triplet splitting problem, fast proton decay etc) are eliminated or improved[69,70]. In these GUT models the metric is factorized, but while for the hierarchy problem $R \sim 1/TeV$, here one considers $R \sim 1/M_{GUT}$ (not so large!). P breaks $N = 2$ SUSY, valid in 5 dimensions, down to $N = 1$ while P' breaks SU(5). At the weak scale there are models where SUSY, valid in $n > 4$ dimensions, is broken by orbifolding[63], in particular the model of ref.[64], where the mass of the Higgs is in principle computable and is predicted to be light.

Symmetry breaking by boundary conditions (BC) is more general than the particular case of orbifolding[65]. Breaking by orbifolding is somewhat rigid: for example, normally the rank remains fixed and it corresponds to Higgs bosons in the adjoint representation (the role of the Higgs is taken by the 5th component of a gauge boson). BC allow a more general breaking pattern and, in particular, can lower the rank of the group. In a simplest version one starts from a 5 dimensional model with two branes at $y = 0; \pi R$. In the action there are terms localised on the branes that also should be considered in the minimization procedure. For a scalar field φ with a mass term (M) on the boundary, one obtains the Neumann BC $\partial_y \varphi = 0$ for $M \to 0$ and the Dirichlet BC $\varphi = 0$ for $M \to \infty$. In gauge theories one can introduce Higgs fields on the brane that take a vev. The crucial property is that the gauge fields take a mass as a consequence of the Higgs mechanism on the boundary but the mass remains finite when the Higgs vev goes to infinity. Thus the Higgs on the boundary only enters as a way to describe and construct the breaking but actually can be removed and still the gauge bosons associated to the broken generators take a finite mass. One is then led to try to formulate "Higgsless models" for EW symmetry breaking based on BC[66]. The RS warped geometry can be adopted with the Planck and the infrared branes. There is a larger gauge symmetry in the bulk which is broken down to different subgroups on the two branes so that finally of the EW symmetry only $U(1)_Q$ remains unbroken. The W and Z take a mass proportional to $1/R$. Dirac fermions are on the bulk and only one chirality has a zero mode on the SM brane. In Higgsless models unitarity, which in general is violated in the absence of a Higgs, is restored by exchange of infinite KK recurrences, or the breaking is delayed by a finite number, with cancellations guaranteed by sum rules implied by the 5-dim symmetry. Actually no compelling, realistic Higgsless model for EW symmetry breaking emerged so far. There are serious problems from EW precision tests [68] because the smallness of the W and Z masses forces R to be rather small and, as a consequence, the spectrum of KK recurrences is quite close. However these models are interesting as rare examples where no Higgs would be found at the LHC but instead new signals appear (new vector bosons, i.e. KK recurrences of the W and Z).

An interesting model that combines the idea of the Higgs as a Goldstone boson and warped extra dimensions was proposed and studied in references[72] with a sort of composite Higgs in a 5-dim AdS theory. It can be considered as a new way to look at walking technicolor[73] using AdS/CFT correspondence. In a RS warped metric framework all SM fields are in the bulk but the Higgs is localised near the TeV brane. The Higgs is a pseudo-Goldstone boson (as in Little Higgs models) and EW symmetry breaking is triggered by top-loop effects. In 4-dim the bulk appears as a strong sector. The 5-dimensional theory is weakly coupled so that the Higgs potential and EW observables can be computed. The Higgs is rather light: $m_H < 185$ GeV. Problems with EW precision tests and the Zbb vertex have been fixed in latest versions. The signals at the LHC for this model are a light Higgs and new resonances at 1-2 TeV

In conclusion, note that apart from Higgsless models (if any?) all theories discussed here have a Higgs in LHC range (most of them light).

10. Effective theories for compositeness

In this approach[74] a low energy theory from truncation of some UV completion is described in terms of an elementary sector (the SM particles minus the Higgs), a composite sector (including the Higgs, massive vector bosons and new fermions) and a mixing sector. The Higgs is a pseudo Goldstone boson of a larger broken gauge group, with the corresponding massive vector bosons. Mass eigenstates are mixtures of elementary and composite states, with light particles mostly elementary and heavy particles mostly composite. But the Higgs is totally composite (perhaps also the right-handed top quark). New physics in the composite sector is well hidden because light particles have small mixing angles. The Higgs is light because only acquires mass through interactions with the light particles from their composite components. This general description can apply to models with a strongly interacting sector as arising from little Higgs or extra dimension scenarios.

11. The anthropic solution

The apparent value of the cosmological constant poses a tremendous, unsolved naturalness problem[36]. Yet the value of is close to the Weinberg upper bound for galaxy formation[75]. Possibly our Universe is just one of infinitely many (Multiverse) continuously created from the vacuum by quantum fluctuations. Different physics takes place in different Universes according to the multitude of string theory solutions (10^{500}). Perhaps we live in a very unlikely Universe but the only one that allows our existence[76]. I find applying the anthropic principle to the SM hierarchy problem excessive. After all we can find plenty of models that easily reduce the fine tuning from 10^{14} to 10^2: why make our Universe so terribly unlikely? By comparison the case of the cosmological constant is a lot different: the context is not as fully specified as the for the SM (quantum gravity, string cosmology, branes in extra dimensions, wormholes through different Universes....)

12. Conclusion

Supersymmetry remains the standard way beyond the SM. What is unique to SUSY, beyond leading to a set of consistent and completely formulated models, as, for example, the MSSM, is that this theory can potentially work up to the GUT energy scale. In this respect it is the most ambitious model because it describes a computable framework that could be valid all the way up to the vicinity of the Planck mass. The SUSY models are perfectly compatible with GUT's and are actually quantitatively supported by coupling unification and also by what we have recently learned on neutrino masses. All other main ideas for going beyond the SM do not share this synthesis with GUT's. The SUSY way is testable, for example at the LHC, and the issue of its validity will be decided by experiment. It is true that we could have expected the first signals of SUSY already at LEP, based on naturality arguments applied to the most minimal models (for example, those with gaugino universality at asymptotic scales). The absence of signals has stimulated the development of new ideas like those of extra dimensions and "little Higgs" models. These ideas are very interesting and provide an important reference for the preparation of LHC experiments. Models along these new ideas are not so completely formulated and studied as for SUSY and no well defined and realistic baseline has sofar emerged. But it is well possible that they might represent at least a part of the truth and it is very important to continue the exploration of new ways beyond the SM. New input from experiment is badly needed, so we all look forward to the start of the LHC.

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