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STATUS OF PRECISION TESTS OF THE STANDARD MODEL

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CONTENT

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INTRODUCTION

The running of LEP1 was terminated in 1995 and close-to-final results of the data analysis are now available and were presented at the Warsaw Conference in July 1996[1],[2]. LEP and SLC started in 1989 and the first results from the collider run at the Tevatron were also first presented at about that time. I went back to my rapporteur talk at the Stanford Conference in August 1989[3] and I found the following best values quoted there for some of the key quantities of interest for the Standard Model (SM) phenomenology: m_Z = 91120(160) MeV; m_t = 130 (50) GeV; $\sin^2\theta_{eff}$ = 0.23300(230) and $\alpha_s(m_Z)$ = 0.110(10). Now, after seven years of experimental and theoretical work (in particular with 16 million Z events analysed altogether by the four LEP experiments) the corresponding numbers, as quoted at the Warsaw Conference, are: $m_Z = 91186.3(2.0)$ MeV; $m_t = 175(6)$ GeV; $\sin^2\theta_{eff} = 0.23165(24)$ and $\alpha_s(m_Z) =$ 0.118(3). The progress is quite evident. The top quark has been at last found and the errors on m_Z and $\sin^2\theta_{eff}$ went down by two and one orders of magnitude respectively. At the start the goals of LEP, SLC and the Tevatron were to: a) perform precision tests of the SM at the level of a few per mille accuracy; b) count neutrinos ($N_\nu = 2.989(12)$); c) search for the top quark ($m_t = 175(6)$ GeV); d) search for the Higgs ($m_H > 65$ GeV); e) search for new particles (none found). While for most of the issues the results can be summarized in very few bits, as just shown, it is by far more complex for the first one. The validity of the SM has been confirmed to a level that I can say was unexpected at the beginning. This is even more true after Warsaw. Contrary to the situation presented at the winter '96 Conferences we are now left with no significant evidence for departures from the SM. The discrepancy on R_c has completely disappeared, that on R_b has been much reduced, and so on, and no convincing hint of new physics is left in the data (also including the first results from LEP2). The impressive success of the SM poses strong limitations on the possible forms of new physics. Favoured are models of the Higgs sector and of new physics that preserve the SM structure and only very delicately improve it, as is the case for fundamental Higgs(es) and Supersymmetry. Disfavoured are models with a nearby strong non-perturbative regime that almost inevitably would affect the radiative corrections, as for composite Higgs(es) or for technicolor and its variants.

STATUS OF THE DATA

The relevant new electroweak data together with their SM values are presented in table 1. The SM values correspond to a fit in terms of m_t, m_H and $\alpha_s(m_Z)$, described later in sect. 3, eq. (14), of all the available data including the CDF/D0 value of m_t . A number of comments on the novel aspects of the data are now in order.

Table 1

Quantity	Data (Warsaw '96)	Standard Model	Pull
m_Z (GeV)	91.1863(20)	91.1861	0.1
Γ_Z (GeV)	2.4946(27)	2.4960	-0.5
σ_h (nb)	41.508(56)	41.465	0.8
R_h	20.788(29)	20.757	0.7
R_b	0.2178(11)	0.2158	1.8
R_c	0.1715(56)	0.1723	-0.1
A^l_{FB}	0.0174(10)	0.0159	1.4
A_τ	0.1401(67)	0.1458	-0.9
A_e	0.1382(76)	0.1458	-1.0
A^b_{FB}	0.0979(23)	0.1022	-1.8
A^c_{FB}	0.0733(48)	0.0730	0.1
A_b	SLD direct 0.863(49)	0.935	-2.2
	LEP indir. 0.895(23)		
	Average 0.889(21)		
A_c	SLD direct 0.625(84)	0.667	-0.2
	LEP indir. 0.670(44)		
	Average 0.660(39)		
$\sin^2\theta_{eff}$(LEP-combined)	0.23200(27)	0.23167	1.2
$A_{LR} \to \sin^2\theta_{eff}$	0.23061(47)	0.23167	-2.2
m_W (GeV)	80.356(125)	80.353	0.3
m_t (GeV)	175(6)	172	0.5

What happened to R_c? The tagging method for charm is based on the reconstruction of exclusive final channels. This is rather complicated and depends on branching ratios and on the probability that a charm quark fragments into given hadrons. A shift in the measured value of the branching ratio for $D^0 \to K^-\pi^+$ and the measurement at LEP of $P(c \to D^*)$, acting on R_c in the same direction, have been sufficient to restore a perfect agreement with the SM.

What happened to R_b? The old result at the winter '96 Conferences was (assuming the SM value for R_c) $R_b = 0.2202(16)$. The present official average, shown in table 1, is much lower and only 1.8σ away from the SM value. The essential difference is the result of a new-from-scratch, much improved analysis from ALEPH, which is given by [1],[2]

$$R_b = 0.2161 \pm 0.0014 \qquad \text{(ALEPH)} . \tag{1}$$

In fact if one combines the average of the "old" measurements, given above, with the "new" ALEPH result one practically finds the official average given by the electroweak LEP working group and reported in table 1. This happens to be true in spite of the fact that in the correct procedure one has to take away the ALEPH contribution, now superseded, from the "old" average and add to it some newly presented refinements to some of the "old" analyses. In view of this, it is clear that the change is mainly due to the new ALEPH result. There are objective improvements in this new analysis. Five mutually exclusive tags are simultaneously used in order to decrease the sensitivity to individual sources of systematic error. Separate vertices are reconstructed in the two hemispheres of each event to minimize correlations between the hemispheres. The implementation of a mass tag on the tracks from each vertex reduces the charm background that dominates the systematics. As a consequence it appears to me that the weight of the new analysis in the combined value should be larger than what is obtained from the stated errors. In view of the ALEPH result the necessity of new physics in R_b has disappeared, while the possibility of some small deviation (more realistic than before) of course is still there. In view of the importance of this issue the other collaborations will go back to their data and freshly reconsider their analyses with the new improvements taken into account.

It is often stated that there is a 3σ deviation on the measured value of A_b with respect to the SM expectation[1],[2]. But in fact that depends on how the data are combined. In my opinion one should rather talk of a 2σ effect. Let us discuss this point in detail. A_b can be measured directly at SLC, taking advantage of the beam longitudinal polarization. SLD finds

$$A_b = 0.863 \pm 0.049 \qquad \text{(SLD direct : } -1.5\sigma) \ , \tag{2}$$

where the discrepancy with respect to the SM value, $A_b^{SM} = 0.935$, has also been indicated. At LEP one measures $A_b^{FB} = 3/4\ A_e A_b$. As seen in table 1, the value found is somewhat below the SM prediction. One can then derive A_b by using the value of A_e obtained, using lepton universality, from the measurements of A_l^{FB}, A_τ, A_e: $A_e = 0.1466(33)$:

$$A_b = 0.890 \pm 0.029 \qquad (\mathrm{LEP,\ A_e\ from\ LEP : -\ 1.6}\sigma) \ . \tag{3}$$

By combining the two above values one obtains

$$A_b = 0.883 \pm 0.025 \qquad (\mathrm{LEP + SLD, A_e\ from\ LEP : -2.1}\sigma) \ . \tag{4}$$

The LEP electroweak working group combines the SLD result with the LEP value for A_b modified by adopting for A_e the SLD+LEP average value, which also includes A_{LR} from SLD, $A_e = 0.1500(25)$:

$$A_b = 0.867 \pm 0.020 \qquad (\mathrm{LEP + SLD,\ A_e\ from\ LEP + SLD : -\ 3.1}\sigma) \ . \tag{5}$$

There is nothing wrong with that but, in this case, the well-known $\sim 2\sigma$ discrepancy of A_{LR} with respect to A_e measured at LEP and also to the SM, which is not related to the b couplings, further contributes to inflate the number of σ's. Since the b couplings are more suspect than the lepton couplings it is perhaps wiser to obtain A_b from LEP by using the SM value for A_e: $A_e^{SM} = 0.1458(16)$, which gives

$$A_b = 0.895 \pm 0.023 \qquad (\mathrm{LEP,\ A_e = A_e^{SM} : -1.7}\sigma) \ . \tag{6}$$

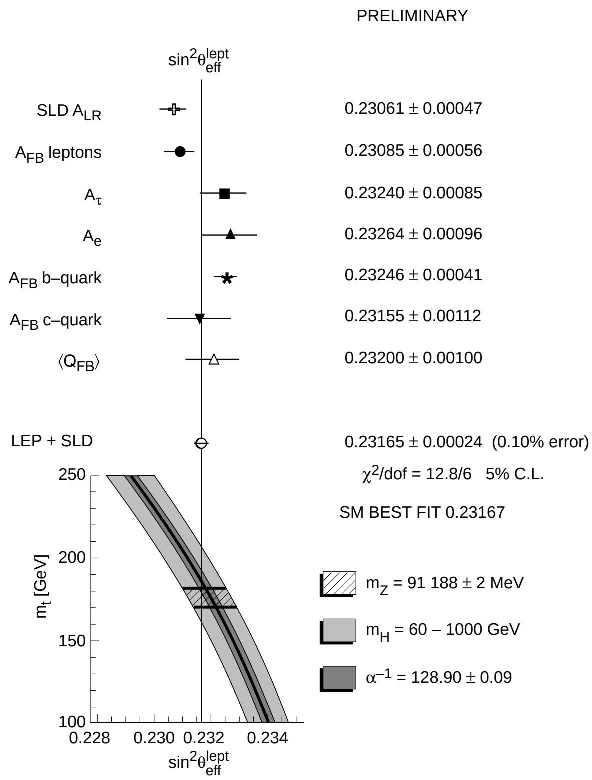

Figure 1

Finally, combining the last value with SLD we have

$$A_b = 0.889 \pm 0.021 \qquad (\mathrm{LEP + SLD,\ A_e = A_e^{SM}} : -2.2\sigma)\ . \tag{7}$$

Note that these are the values reported in table 1.

Finally if one looks at the values of $\sin^2\theta_{eff}$ obtained from different observables, shown in fig. 1 , one notices that the value obtained from A_l^{FB} is somewhat low (indeed quite in agreement with its determination by SLD from A_{LR}). Looking closer, this is due to the FB asymmetry of the τ lepton that, systematically in all four LEP experiments, has a central value above that of e and μ [1],[2]. The combined value for the τ channel is A_τ^{FB} =0.0201(18) while the combined average of e and μ is $A_{e/\mu}^{FB} = 0.0162(11)$. On the other hand A_τ and Γ_τ appear normal. In principle these two facts are not incompatible, because the FB lepton asymmetries are very small. The extraction of A_τ^{FB} from the data on the angular distribution of τ's could be biased if the imaginary part of the continuum was altered by some non-universal new physics effect[4]. But a more trivial experimental problem is at the moment more plausible.

The distribution of measured values of $\sin^2\theta_{eff}$, as it is summarized in fig. 1, is somewhat wide (χ^2/d.o.f. $= 2.13$) with A_l^{FB} and A_{LR} far on one side and A_b^{FB} on the other side. In view of this it would perhaps be appropriate to enlarge the error on the average from ± 0.00024 up to $\pm\sqrt{2.13}\ 0.00024 = \pm 0.00034$, according to the recipe adopted by the Particle Data Group. Thus from time to time in the following we will use the average

$$\sin^2\theta_{eff} = 0.23165 \pm 0.00034 \tag{8}$$

For the analysis of electroweak data in the SM one starts from the input parameters: some of them, α, G_F and m_Z, are very well measured, some other, $m_{f_{light}}$, m_t and $\alpha_s(m_Z)$ are only approximately determined, while m_H is largely unknown. With respect to m_t the situation has much improved since the CDF/D0 direct measurement of the top quark mass[5]. From the input parameters one computes the radiative corrections[6],[7] to a sufficient precision to match the experimental capabilities. Then one compares the theoretical predictions and the data for the numerous observables that have been measured, checks the consistency of the theory and derives constraints on m_t, $\alpha_s(m_Z)$, and hopefully also on m_H.

Some comments on the least known of the input parameters are now in order. The only practically relevant terms where precise values of the light quark masses, $m_{f_{light}}$, are needed are those related to the hadronic contribution to the photon vacuum polarization diagrams that determine $\alpha(m_Z)$. This correction is of order 6%, much larger than the accuracy of a few per mille of the precision tests. Fortunately, one can use the actual data to in principle solve the related ambiguity. But we shall see that the left-over uncertainty is still one of the main sources of theoretical error. As is well known[8]–[18], the QED running coupling is given by:

$$\begin{aligned} \alpha(s) &= \frac{\alpha}{1-\Delta\alpha(s)} \\ \Delta\alpha(s) &= \Pi(s) = \Pi_\gamma(0) - \mathrm{Re}\Pi_\gamma(s) \ , \end{aligned} \tag{9}$$

where $\Pi(s)$ is proportional to the sum of all 1-particle irreducible vacuum polarization diagrams. In perturbation theory $\Delta\alpha(s)$ is given by

$$\Delta\alpha(s) = \frac{\alpha}{3\pi} \sum_f Q_f^2 N_{Cf} \left(\log \frac{2}{m_f^2} - \frac{5}{3} \right) \ , \tag{10}$$

where $N_{Cf} = 3$ for quarks and 1 for leptons. However, the perturbative formula is only reliable for leptons, not for quarks (because of the unknown values of the effective quark masses). Separating the leptonic, the light quark and the top quark contributions to $\Delta\alpha(s)$ we have:

$$\Delta\alpha(s) = \Delta\alpha(s)_1 + \Delta\alpha(s)_h + \Delta\alpha(s)_t \tag{11}$$

with[18]:

$$\Delta\alpha(s)_1 = 0.0331421 \ ; \quad \Delta\alpha(s)_t = \frac{\alpha}{3\pi} \frac{4}{15} \frac{m_Z^2}{m_t^2} = -0.000061 \ . \tag{12}$$

Note that in QED there is decoupling so that the top quark contribution approaches zero in the large m_t limit. For $\Delta\alpha(s)_h$ one can use (9) and the Cauchy theorem to obtain the representation:

$$\Delta\alpha(m_Z^2)_h = -\frac{\alpha m_Z^2}{3\pi} \mathrm{Re} \int_{4m_\pi^2}^{\infty} \frac{ds}{s} \frac{R(s)}{s - m_Z^2 - i\epsilon} \tag{13}$$

where $R(s)$ is the familiar ratio of the hadronic to the point-like $\ell^+\ell^-$ cross-section from photon exchange in e^+e^- annihilation. At s large, one can use the perturbative expansion for $R(s)$ while at small s one can use the actual data.

Recently there has been a lot of activity on this subject and a number of independent new estimates of $\alpha(m_Z)$ have appeared in the literature[8]. In table 2 we report the results of these new computations together with the most significant earlier determinations (previously the generally accepted value was that of Jegerlehner in 1991[12]).

Table 2

Author	Year and Ref.	$\Delta\alpha(m_Z^2)_h$	$\alpha(m_Z^2)^{-1}$
Jegerlehner	1986 [9]	0.0285 ± 0.0007	128.83 ± 0.09
Lynn et al.	1987 [10]	0.0283 ± 0.0012	128.86 ± 0.16
Burkhardt et al.	1989 [11]	0.0287 ± 0.0009	128.80 ± 0.12
Jegerlehner	1991 [12]	0.0282 ± 0.0009	128.87 ± 0.12
Swartz	1994 [13]	0.02666 ± 0.00075	129.08 ± 0.10
Swartz (rev.)	1995 [14]	0.0276 ± 0.0004	128.96 ± 0.06
Martin et al.	1994 [15]	0.02732 ± 0.00042	128.99 ± 0.06
Nevzorov et al.	1994 [16]	0.0280 ± 0.0004	128.90 ± 0.06
Burkhardt et al.	1995 [17]	0.0280 ± 0.0007	128.89 ± 0.09
Eidelman et al.	1995 [18]	0.0280 ± 0.0007	128.90 ± 0.09

The differences among the recent determinations are due to the procedures adopted for fitting the data and treating the errors, for performing the numerical integration, etc. The differences are also due to the threshold chosen to start the application of perturbative QCD at large s and to the value adopted for $\alpha_s(m_Z)$. For example, in its first version Swartz[13] used parametric forms to fit the data, while most of the other determinations use a trapezoidal rule to integrate across the data points. It was observed that the parametric fitting introduces a definite bias[14]. In fact Swartz gets systematically lower results for all ranges of s. In its revised version[14] Swartz improves his numerical procedure. Martin et al.[15] use perturbative QCD down to $\sqrt{s} = 3$ GeV (except in the upsilon region) with $\alpha_s(m_Z) = 0.118 \pm 0.007$. Eidelman et al.[18] only use perturbative QCD for $\sqrt{s} > 40$ GeV and with $\alpha_s(m_Z) = 0.126 \pm 0.005$, i.e. the value found at LEP. They use the trapezoidal rule. Nevzorov et al.[16] make a rather crude model with one resonance per channel plus perturbative QCD with $\alpha_s(m_Z) = 0.125 \pm 0.005$. Burkhardt et al.[17] use perturbative QCD for $\sqrt{s} > 12$ GeV, but with a very conservative error on $\alpha_s(m_Z) = 0.124 \pm 0.021$. This value was determined [19] from e^+e^- data below LEP energies. The excitement produced by the original claim by Swartz[13] of a relatively large discrepancy with respect to the value obtained by Jegerlehner[12] resulted in a useful debate. As a conclusion of this re-evaluation of the problem the method of Jegerlehner has proved its solidity. As a consequence I think that the recent update by Eidelman and Jegerlehner[18] gives a quite reliable result (which is the one used by the LEP groups and in the following). Also, I do not think that a smaller error than quoted by these authors can be justified.

As for the strong coupling $\alpha_s(m_Z)$ we will discuss in detail the interesting recent developments in sect. 4. The world average central value is quite stable around 0.118, before and after the most recent results. The error is going down because the dispersion among the different measurements is much smaller in the most recent set of data. The error is taken to be between ± 0.003 and ± 0.005, depending on how conservative one wants to be. Thus in the following our reference value will be $\alpha_s(m_Z) = 0.118 \pm 0.005$.

Finally a few words on the current status of the direct measurement of m_t. The error is rapidly going down. It was ± 9 GeV before the Warsaw Conference, it is now ± 6 GeV [5]. I think one is soon approaching a level where a more careful investigation

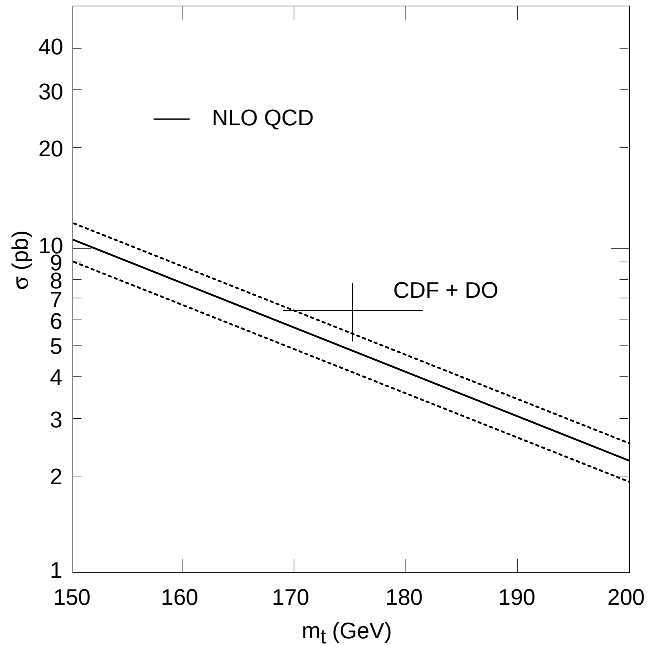

Figure 2

of the effects of colour rearrangement on the determination of m_t is needed. One wants to determine the top quark mass, defined as the invariant mass of its decay products (i.e. $b+W+$ gluons $+$ γ's). However, due to the need of colour rearrangement, the top quark and its decay products cannot be really isolated from the rest of the event. Some smearing of the mass distribution is induced by this colour crosstalk, which involves the decay products of the top, those of the antitop and also the fragments of the incoming (anti)protons. A reliable quantitative computation of the smearing effect on the m_t determination is difficult because of the importance of non-perturbative effects. An induced error of the order of a few GeV on m_t is reasonably expected. Thus further progress on the m_t determination demands tackling this problem in more depth.

The measured top production cross section is in fair agreement with the QCD prediction, but the central value is a bit large (see fig. 2)[20]. The world average for the cross section times branching ratio is $\sigma B = 6.4 \pm 1.3$ pb and the QCD prediction for σ is $\sigma_{QCD} = 4.75 \pm 0.65$ pb [21]. Thus the branching ratio $B = B(t \to bW)$ cannot be far from 100% unless there is also some additional production mechanism from new physics.

In order to appreciate the relative importance of the different sources of theoretical errors for precision tests of the SM, I report in table 3 a comparison for the most relevant observables, evaluated using ref. [22].

What it is important to stress is that the ambiguity from m_t, once by far the largest one, is by now smaller than the error from m_H. We also see from table 3 that the error from $\Delta\alpha(m_Z)$ is expecially important for $\sin^2\theta_{eff}$ and, to a lesser extent, is also sizeable for Γ_Z and ϵ_3.

We now discuss fitting the data in the SM. As the mass of the top quark is now rather precisely known from CDF and D0 one must distinguish between two different types of fits. In one type one wants to answer the question: Is m_t from radiative corrections in agreement with the direct measurement at the Tevatron? For answering

Table 3: Errors from different sources: Δ^{exp}_{now} is the present experimental error; $\Delta\alpha^{-1}$ is the impact of $\Delta\alpha^{-1} = \pm 0.09$; Δ_{th} is the estimated theoretical error from higher orders; Δm_t is from $\Delta m_t = \pm 6$GeV; Δm_H is from Δm_H = 60–1000 GeV; $\Delta\alpha_s$ corresponds to $\Delta\alpha_s = \pm 0.005$. The epsilon parameters are defined in ref. [23].

Parameter	Δ^{exp}_{now}	$\Delta\alpha^{-1}$	Δ_{th}	Δm_t	Δm_H	$\Delta\alpha_s$
Γ_Z (MeV)	± 2.7	± 0.7	± 0.8	± 1.4	± 4.6	± 2.7
σ_h (pb)	56	1	4.3	3.3	4	2.7
$R_h \cdot 10^3$	29	4.3	5	2	13.5	34
Γ_l (keV)	110	11	15	55	120	6
$A^l_{FB} \cdot 10^4$	10	4.2	1.3	3.3	13	0.3
$\sin^2\theta \cdot 10^4$	$\sim$3	2.3	0.8	1.9	7.5	0.15
m_W (MeV)	125	12	9	37	100	4
$R_b \cdot 10^4$	11	0.1	1	2.1	0.25	0
$\epsilon_1 \cdot 10^3$	1.3		$\sim$0.1			0.4
$\epsilon_3 \cdot 10^3$	1.4	0.6	$\sim$0.1			0.25
$\epsilon_b \cdot 10^3$	3.2		$\sim$0.1			2

this interesting but somewhat limited question, clearly one must exclude the CDF/D0 measurement of m_t from the input set of data. Fitting the data in terms of m_t, m_H and $\alpha_s(m_Z)$ one finds the results shown in table 4[2].

Table 4

Parameter	LEP	LEP + SLD	All $\neq m_t$
$\alpha_s(m_Z)$	0.1211(32)	0.1200(32)	0.1202(33)
m_t (GeV)	155(14)	156(11)	157(10)
m_H (GeV)	86(+202 − 14)	48(+83 − 26)	149(+148 − 82)
$(m_H)_{\rm MAX}$ at 1.64σ	417	184	392
χ^2/dof	5/8	18/11	18/13

The extracted value of m_t is typically a bit too low. For example, from LEP data alone one finds $m_t = 155(14)$ GeV. But this is simply due to R_b being taken from the official average: $R_b = 0.2178(11)$. If m_H is not fixed the fit prefers lower values of m_t to adjust R_b. In fact by removing R_b from the input data one increases the central value of m_t from 155 to 171 GeV. In this context it is important to remark that fixing m_H at 300 GeV, as is often done, is by now completely obsolete, because it introduces a strong bias on the fitted value of m_t. The change induced on the fitted value of m_t when moving m_H from 300 to 65 or 1000 GeV is in fact larger than the error on the direct measurement of m_t.

In a more general type of fit, e.g. for determining the overall consistency of the SM or the best present estimate for some quantity, say m_W, one should of course not ignore the existing direct determination of m_t. Then, from all the available data, including $m_t = 175(6)$ GeV, by fitting m_t, m_H and $\alpha_s(m_Z)$ one finds (with χ^2/d.o.f. = 19/14) [2] (see also[24]):

$$m_t = 172 \pm 6 \text{ GeV} ,$$

$$
\begin{aligned}
m_H &= 149 + 148 - 82 \text{ (or } m_H < 392 \text{ GeV at } 1.64\sigma) \\
\alpha_s(m_Z) &= 0.1202 \pm 0.0033 \ .
\end{aligned}
\tag{14}
$$

This is the fit reported in table 1. The corresponding fitted values of $\sin^2\theta_{eff}$ and m_W are:

$$
\begin{aligned}
\sin^2\theta_{eff} &= 0.23167 \pm 0.0002 \\
m_W &= 80.352 \pm 0.034 \text{ GeV} \ .
\end{aligned}
\tag{15}
$$

The error of 34 MeV on m_W clearly sets up a goal for the direct measurement of m_W at LEP2 and the Tevatron.

STATUS OF $\alpha_s(m_Z)$

There are important developments in the experimental determination of $\alpha_s(m_Z)$[25]. There is now a much better agreement between the different methods of measuring $\alpha_s(m_Z)$. In fact the value of $\alpha_s(m_Z)$ from the Z line shape went down and the values from scaling violations in deep inelastic scattering and from lattice QCD went up. We will discuss these developments in detail in the following.

The value of $\alpha_s(m_Z)$ from the Z line shape (assuming that the SM is valid for Γ_h, which is not completely evident in view of R_b) went down for two reasons[1],[2]. First the value extracted from R_h only, which was $\alpha_s(m_Z) = 0.126(5)$, is now down to $\alpha_s(m_Z) = 0.124(5)$. Second the value from all the Z data changed from $\alpha_s(m_Z) = 0.124(5)$ down to $\alpha_s(m_Z) = 0.120(4)$, which corresponds to the fit in eq. (14) . The main reason for this decrease is the new value of σ_h (with a sizeably smaller error than in the past) that prefers a smaller $\alpha_s(m_Z)$. However this determination depends on the assumption that Γ_b is given by the SM. We recall that R_b itself with good approximation is independent of α_s, but its deviation from the SM would indicate an anomaly in Γ_b hence in Γ_h. Taking a possible anomaly in R_b into account the Z line shape determination of $\alpha_s(m_Z)$ becomes approximately:

$$
\alpha_s(m_Z) = (0.120 \pm 0.004) - 4\delta R_b \ . \tag{16}
$$

If the ALEPH value for R_b (see eq .(1)) is adopted, the central value of $\alpha_s(m_Z)$ is not much changed, but of course the error on δR_b is transferred on $\alpha_s(m_Z)$, which becomes

$$
\alpha_s(m_Z) = 0.119 \pm 0.007 \ . \tag{17}
$$

If, instead, one takes R_b from table 1 one obtains a much smaller central value:

$$
\alpha_s(m_Z) = 0.112 \pm 0.006 \tag{18}
$$

Summarizing: the Z line shape result for $\alpha_s(m_Z)$, obtained with the assumption that Γ_h is given by the SM, went down a bit. The central value could be shifted further down if R_b is in excess with respect to the SM.

While $\alpha_s(m_Z)$ from LEP goes down, $\alpha_s(m_Z)$ from the scaling violations in deep inelastic scattering goes up. To me the most surprising result from Warsaw was the announcement by the CCFR collaboration that their well-known published analysis of $\alpha_s(m_Z)$ from xF_3 and F_2 in neutrino scattering off Fe target is now superseded by a re-analysis of the data based on better energy calibration[26]. We recall that their previous result, $\alpha_s(m_Z) = 0.111(3\,\mathrm{exp})$, being in perfect agreement with the value obtained from e/μ beam data by BCDMS and SLAC combined[27], $\alpha_s(m_Z) = 0.113(3\,\mathrm{exp})$, convinced most of us that the average value of $\alpha_s(m_Z)$ from deep inelastic scattering was close to 0.112. Now the new result presented in Warsaw is [26], [25]

$$\alpha_s(m_Z) = 0.119 \pm 0.0015(\mathrm{stat}) \pm 0.0035(\mathrm{syst}) \pm 0.004(th) \quad (\mathrm{CCFR-revised})\ , \tag{19}$$

where the error also includes the collaboration estimate of the theoretical error from scale and renormalization scheme ambiguities. As a consequence the new combined value of $\alpha_s(m_Z)$ from scaling violations in deep inelastic scattering is given by

$$\alpha_s(m_Z) = 0.115 \pm 0.006\ , \tag{20}$$

with my more conservative estimate, of the common theoretical error (Schmelling, the rapporteur in Warsaw quotes ± 0.005 [25]). If we compare eq. (20) with LEP eq. (14), we see that, whatever our choice of theoretical errors is, there is no need for any new physics in R_b to fill the gap between the two determinations of $\alpha_s(m_Z)$.

Finally $\alpha_s(m_Z)$ from lattice QCD is also going up [28]. The main new development is a theoretical study of the error associated with the extrapolation from unphysical values of the light quark masses, which is used in the lattice extraction of $\alpha_s(m_Z)$ from quarkonium splittings. According to ref. [29] this effect amounts to a shift upward of $+0.003$ in the value of $\alpha_s(m_Z)$. From the latest unquenched determinations of $\alpha_s(m_Z)$, Flynn, the rapporteur in Warsaw[28], gives an average of 0.117(3). But the lattice measurements of $\alpha_s(m_Z)$ moved very fast over the last few years. At the Dallas conference in 1992, the quoted value (from quenched computations) was $\alpha_s(m_Z) = 0.105(4)$ [30], while at Beijing in 1995 the claimed value was $\alpha_s(m_Z) = 0.113(2)$ but the error was estimated to be ± 0.007 by the rapporteur Michael[31]. So, with the present central value, I will keep this more conservative error in the following:

$$\alpha_s(m_Z) = 0.117 \pm 0.007\ . \tag{21}$$

To my knowledge, there are no other important new results on the determination of $\alpha_s(m_Z)$. Adding a few more well-established measurements of $\alpha_s(m_Z)$ we have table 5, where the errors denote my personal view of the weights the different methods should have in the average (in brackets Th and Exp are labels that indicate whether the dominant error is theoretical or experimental).

The average value given as

$$\alpha_s(m_Z) = 0.118 \pm 0.003 \tag{22}$$

is very stable. The same value was quoted by Schmelling, a rapporteur at the Warsaw Conference[25], with a different treatment of errors. Had we used $\alpha_s(m_Z)$ from the Z line shape assuming the SM value for R_b, i.e. $\alpha_s(m_Z) = 0.120 \pm 0.004$, the average value would have been 0.119. To be safe one could increase the error to ± 0.005.

A MORE MODEL-INDEPENDENT APPROACH

We now discuss an update of the epsilon analysis[23]. The epsilon method is more complete and less model-dependent than the similar approach based on the variables

Table 5

Measurements	$\alpha_s(m_Z)$
R_τ	0.122 ± 0.007 (Th)
Deep Inelastic Scattering	0.115 ± 0.006 (Th)
Y_{decay}	0.112 ± 0.010 (Th)
Lattice QCD	0.117 ± 0.007 (Th)
$Re^+e^-(\sqrt{s} < 62$ GeV)	0.124 ± 0.021 (Exp)
Fragmentation functions in e^+e^-	0.124 ± 0.010 (Th)
Jets in e^+e^- at and below the Z	0.121 ± 0.008 (Th)
Z line shape (taking R_b from ALEPH)	0.119 ± 0.007 (Exp)

S, T and U [32]–[35] which, from the start, necessarily assumes dominance of vacuum polarization diagrams from new physics and truncation of the q^2 expansion of the corresponding amplitudes. In a completely model-independent way we define[23] four variables, called ϵ_1, ϵ_2, ϵ_3 and ϵ_b, that are precisely measured and can be compared with the predictions of different theories. The quantities ϵ_1, ϵ_2, ϵ_3 and ϵ_b are defined in ref. [23] in one-to-one correspondence with the set of observables m_W/m_Z, Γ_l, A_l^{FB} and R_b. The four epsilons are defined without need of specifying m_t and m_H. In the SM, for all observables at the Z pole, the whole dependence on m_t and m_H arising from one-loop diagrams only enters through the epsilons. The same is true for any extension of the SM such that all possible deviations only occur through vacuum polarization diagrams and/or the $Z \to b\bar{b}$ vertex.

The epsilons represent an efficient parametrization of the small deviations from what is solidly established, in a way that is unaffected by our relative ignorance of m_t and m_H. The variables S, T, U depend on m_t and m_H because they are defined as deviations from the complete SM prediction for specified m_t and m_H. Instead the epsilons are defined with respect to a reference approximation, which does not depend on m_t and m_H. In fact the epsilons are defined in such a way that they are exactly zero in the SM in the limit of neglecting all pure weak loop-corrections (i.e. when only the predictions from the tree level SM plus pure QED and pure QCD corrections are taken into account). This very simple version of improved Born approximation is a good first approximation according to the data. Values of the epsilons in the SM are given in table 6 [22],[23].

By combining the value of m_W/m_Z with the LEP results on the charged lepton partial width and the forward–backward asymmetry, all given in table 1, and following the definitions of ref. [23], one obtains:

$$\begin{aligned} \epsilon_1 &= \Delta\rho = (4.3 \pm 1.4) \times 10^{-3} \\ \epsilon_2 &= (-6.9 \pm 3.4) \times 10^{-3} \\ \epsilon_3 &= (3.0 \pm 1.8) \times 10^{-3} \;. \end{aligned} \tag{23}$$

Finally, by adding the value of R_b listed in table 1 and using the definition of ϵ_b given in ref. [23] one finds (note that ϵ_b is defined through R_b and the expression of R_b as a function of ϵ_b is practically independent of α_s):

$$\epsilon_b = (-1.1 \pm 2.8) \times 10^{-3} \qquad (R_b \text{ from table 1}) \;. \tag{24}$$

Table 6: Values of the epsilons in the SM as functions of m_t and m_H as obtained from recent versions[22] of ZFITTER and TOPAZ0. These values (in 10^{-3} units) are obtained for $\alpha_s(m_Z) = 0.118$, $\alpha(m_Z) = 1/128.87$, but the theoretical predictions are essentially independent of $\alpha_s(m_Z)$ and $\alpha(m_Z)$ [23].

m_t (GeV)	ϵ_1 m_H (GeV) =			ϵ_2 m_H (GeV) =			ϵ_3 m_H (GeV) =			ϵ_b All m_H
	65	300	1000	65	300	1000	65	300	1000	
150	3.47	2.76	1.61	−6.99	−6.61	−6.4	4.67	5.99	6.66	−4.45
160	4.34	3.59	2.38	−7.29	−6.9	−6.69	4.6	5.91	6.55	−5.28
170	5.25	4.46	3.21	−7.6	−7.2	−6.97	4.52	5.82	6.43	−6.13
180	6.2	5.37	4.1	−7.93	−7.51	−7.24	4.42	5.72	6.34	−7.02
190	7.2	6.33	5.07	−8.29	−7.81	−7.49	4.31	5.6	6.26	−7.95
200	8.26	7.34	6.1	−8.65	−8.12	−7.75	4.19	5.49	6.19	−8.92

This is the value that corresponds to the official average reported in table 1 which I have criticized. Here in this epsilon analysis we prefer to use the ALEPH value for R_b, ($R_b = 0.2161(14)$), which leads to

$$\epsilon_b = (-5.7 \pm 3.4) \times 10^{-3} \qquad (\mathrm{R_b \ from\ ALEPH}) \tag{25}$$

To proceed further and include other measured observables in the analysis, we need to make some dynamical assumptions. The minimum amount of model dependence is introduced by including other purely leptonic quantities at the Z pole such as $A_{\tau_{pol}}$, A_e (measured from the angular dependence of the τ polarization) and A_{LR} (measured by SLD). For this step, one is simply relying on lepton universality. Note that the choice of A_l^{FB} as one of the defining variables appears at present not particularly lucky, because the corresponding determination of $\sin^2\theta_{eff}$ markedly underfluctuates with respect to the average value (see fig. 1). We then use the combined value of $\sin^2\theta_{eff}$ obtained from the whole set of asymmetries measured at LEP and SLC, with the error increased according to eq. (8) and the related discussion. At this stage the best values of ϵ_1, ϵ_2, ϵ_3 and ϵ_b are modified according to

$$\begin{aligned}
\epsilon_1 &= \Delta\rho = (4.7 \pm 1.3) \times 10^{-3} \\
\epsilon_2 &= (-7.8 \pm 3.3) \times 10^{-3} \\
\epsilon_3 &= (4.8 \pm 1.4) \times 10^{-3} \\
\epsilon_b &= (-5.7 \pm 3.4) \times 10^{-3}
\end{aligned} \tag{26}$$

In fig. 3 we report the 1σ ellipse in the ϵ_1–ϵ_3 plane that correspond to this set of input data.

All observables measured on the Z peak at LEP can be included in the analysis, provided that we assume that all deviations from the SM are only contained in vacuum polarization diagrams (without demanding a truncation of the q^2 dependence of the corresponding functions) and/or the $Z \to b\bar{b}$ vertex. From a global fit of the data on m_W/m_Z, Γ_T, R_h, σ_h, R_b and $\sin^2\theta_{eff}$ (for LEP data, we have taken the correlation matrix for Γ_T, R_h and σ_h given by the LEP experiments[2], while we have considered the additional information on R_b and $\sin^2\theta_{eff}$ as independent), we obtain:

$$\epsilon_1 = \Delta\rho = (4.7 \pm 1.3) \times 10^{-3}$$

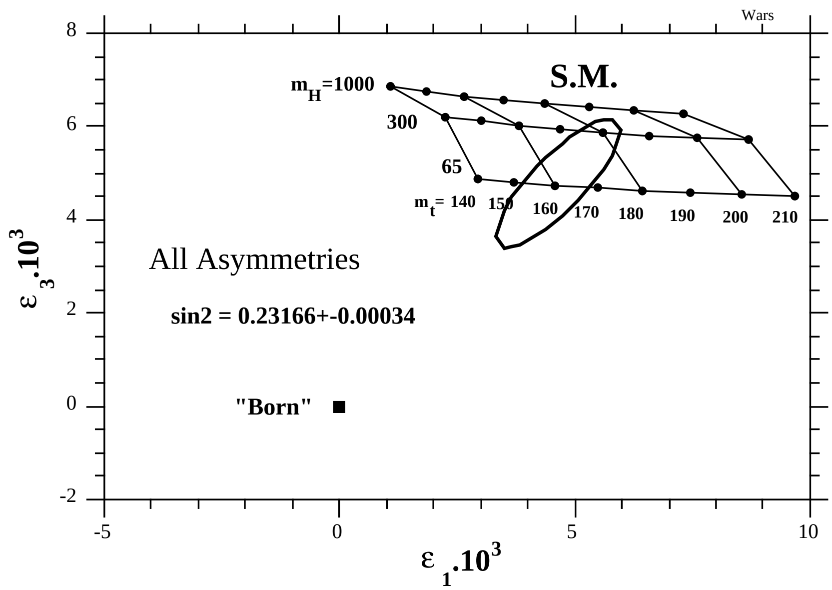

Figure 3

$$\begin{aligned} \epsilon_2 &= (-7.8 \pm 3.3) \times 10^{-3} \\ \epsilon_3 &= (4.7 \pm 1.4) \times 10^{-3} \\ \epsilon_b &= (-4.8 \pm 3.2) \times 10^{-3} \end{aligned} \quad (27)$$

The comparison of theory and experiment in the planes ϵ_1–ϵ_3 and ϵ_b–ϵ_3 is shown in figs. 4 and 5, respectively. Note that adding the hadronic quantities hardly makes a difference in the ϵ_1–ϵ_3 plot in comparison with fig. 3 which only included the leptonic variables. In other words the inclusive hadronic quantities do not show any peculiarity. A number of interesting features are clearly visible from this plot. First, the good agreement with the SM and the evidence for weak corrections, measured by the distance of the data from the improved Born approximation point (based on tree level SM plus pure QED or QCD corrections). Second, we see the preference for light Higgs or, equivalently, the tendency for ϵ_3 to be rather on the low side (both features are now somewhat less pronounced than they used to be). Finally, if the Higgs is light the preferred value of m_t is somewhat lower than the Tevatron result (which in this analysis is not included in the input data). The data ellipse in the ϵ_b–ϵ_3 plane is consistent with the SM and the CDF/D0 value of m_t. This is because we have taken the ALEPH value for R_b. For comparison, we also show in figs. 6 and 7 the same plots as in figs.4 and 5, but for the official average values of R_b and $\sin^2\theta_{eff}$ as reported in table 1. The main difference is the obvious displacement of ϵ_b and the smaller errors in the ϵ_1–ϵ_3 plot. Finally, the status of ϵ_2 is presented in fig. 8. The agreement is very good. ϵ_2 is sensitive to m_W and a more precise test will only be possible when the measurement of m_W will be much improved at LEP2 and the Tevatron.

To include in our analysis lower-energy observables as well, a stronger hypothesis needs to be made: vacuum polarization diagrams are allowed to vary from the SM only in their constant and first derivative terms in a q^2 expansion[33]–[35], a likely picture, e.g. in technicolor theories[36]–[38]. In such a case, one can, for example, add to the analysis the ratio R_ν of neutral to charged current processes in deep inelastic neutrino scattering on nuclei[39], the "weak charge" Q_W measured in atomic parity violation experiments on

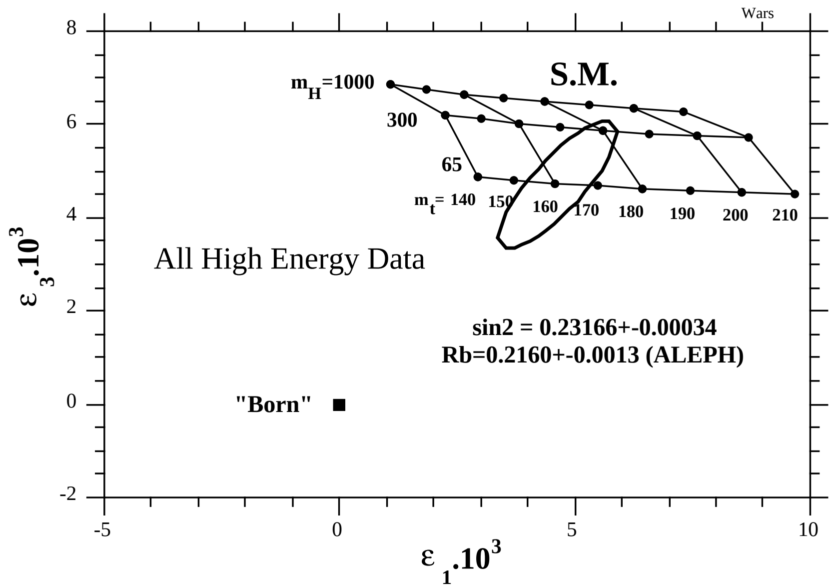

Figure 4

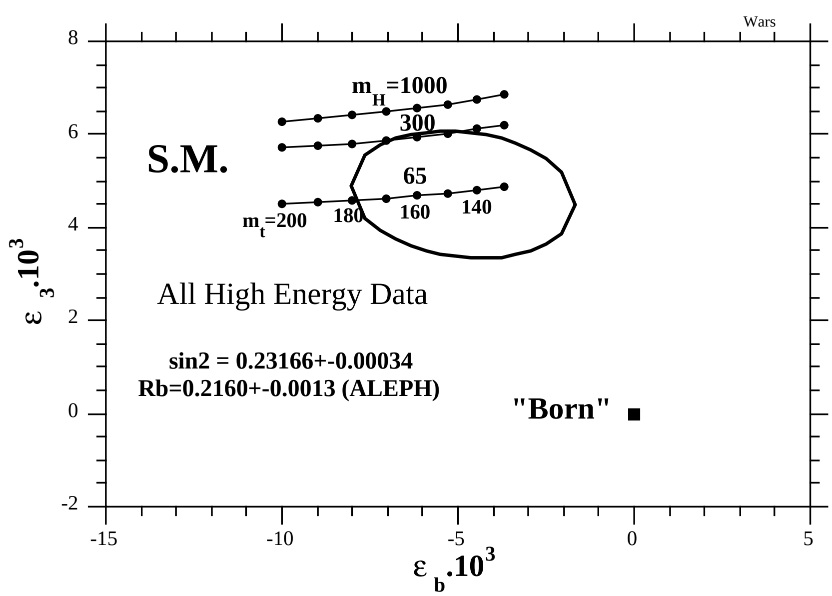

Figure 5

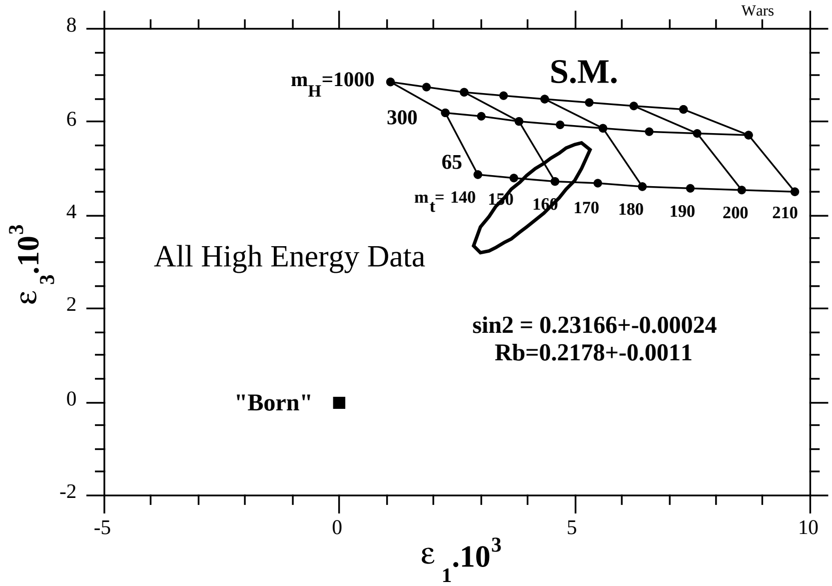

Figure 6

Cs [40], and the measurement of g_V/g_A from $\nu_\mu e$ scattering[41]. In this way one obtains the global fit (R_b from ALEPH, $\sin^2\theta_{eff}$ with enlarged error as in eq. (8)):

$$\begin{aligned} \epsilon_1 &= \Delta\rho = (4.3 \pm 1.2) \times 10^{-3} \\ \epsilon_2 &= (-8.0 \pm 3.3) \times 10^{-3} \\ \epsilon_3 &= (4.4 \pm 1.3) \times 10^{-3} \\ \epsilon_b &= (-4.6 \pm 3.2) \times 10^{-3}\ . \end{aligned} \tag{28}$$

With the progress of LEP, the low-energy data, while important as a check that no deviations from the expected q^2 dependence arise, play a lesser role in the global fit. Note that the present ambiguity on the value of $\delta\alpha^{-1}(m_Z) = \pm 0.09$ [18] corresponds to an uncertainty on ϵ_3 (the other epsilons are not much affected) given by $\Delta\epsilon_3\ 10^3 = \pm 0.6$ [23]. Thus the theoretical error is still comfortably less than the experimental error.

To conclude this section I would like to add some comments. As is clearly indicated in figs. 3 to 8 there is by now solid evidence for departures from the "improved Born approximation" where all the epsilons vanish. In other words a strong evidence for the pure weak radiative corrections has been obtained, and LEP/SLC are now measuring the various components of these radiative corrections. For example, some authors [42] have studied the sensitivity of the data to a particularly interesting subset of the weak radiative corrections, i.e. the purely bosonic part. These terms arise from the virtual exchange of gauge bosons and Higgses. The result is that indeed the measurements are sufficiently precise to require the presence of these contributions in order to fit the data.

CONCEPTUAL PROBLEMS WITH THE STANDARD MODEL

Given the striking 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

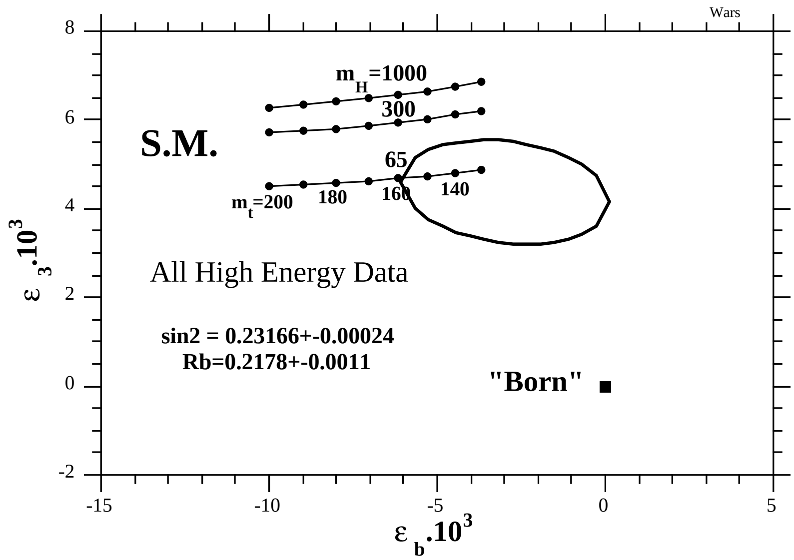

Figure 7

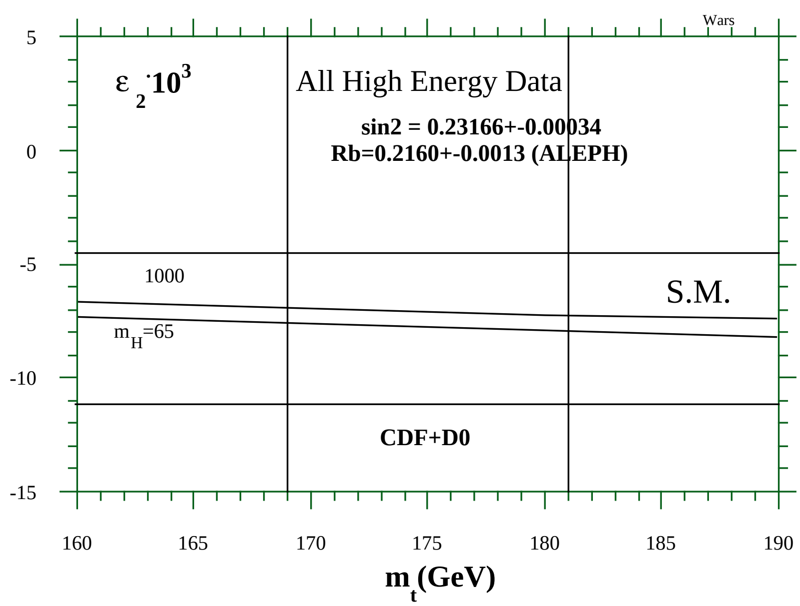

Figure 8

is closed? The main reason is that there are strong conceptual indications for physics beyond the SM.

It is considered highly implausible that the origin of the electroweak symmetry breaking can be explained by the standard Higgs mechanism, without accompanying new phenomena. New physics should be manifest at energies in the TeV domain. This conclusion follows fron an extrapolation of the SM at very high energies. The computed behaviour of the $SU(3) \otimes SU(2) \otimes U(1)$ couplings with energy clearly points towards the unification of the electroweak and strong forces (Grand Unified Theories: GUTs) at scales of energy $M_{GUT} \sim 10^{14}$–10^{16} GeV, which are close to the scale of quantum gravity, $M_{Pl} \sim 10^{19}$ GeV [43]. One can also imagine a unified theory of all interactions also including gravity (at present superstrings [44] provide the best attempt at such a theory). Thus GUTs and the realm of quantum gravity set a very distant energy horizon that modern particle theory can no longer ignore. Can the SM without new physics be valid up to such large energies? This appears unlikely because the structure of the SM could not naturally explain the relative smallness of the weak scale of mass, set by the Higgs mechanism at $m \sim 1/\sqrt{G_F} \sim 250$ GeV, G_F being the Fermi coupling constant. The weak scale m is $\sim 10^{17}$ times smaller than M_{Pl}. Even if the weak scale is set near 250 GeV at the classical level, quantum fluctuations would naturally shift it up to where new physics starts to apply, in particular up to M_{Pl} if there was no new physics up to gravity. This so-called hierarchy problem[45] is related to the presence of fundamental scalar fields in the theory with quadratic mass divergences and no protective extra symmetry at $m = 0$. For fermions, first, the divergences are logaritmic and, second, at $m = 0$ an additional symmetry, i.e. the 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 is absorbed in a redefinition of masses and couplings. Rather the hierarchy problem is one of naturalness. If we consider the cut off as a manifestation of 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 cancellation arise.

According to the above argument the observed value of $m \sim 250$ GeV is indicative of the existence of new physics nearby. There are two main possibilities. Either there exist fundamental scalar Higgses, but the theory is stabilized by supersymmetry, the boson–fermion symmetry that would downgrade the degree of divergence from quadratic to logarithmic. For approximate supersymmetry the cut off is replaced by the splitting between the normal particles and their supersymmetric partners. Then naturalness demands that this splitting (times the size of the weak gauge coupling) is of the order of the weak scale of mass, i.e. the separation within supermultiplets should be of the order of no more than a few TeV. In this case the masses of most supersymmetric partners of the known particles, a very large menagerie of states, would fall, at least in part, in the discovery reach of the LHC. There are consistent, fully formulated field theories constructed on the basis of this idea, the simplest one being the MSSM[46]. Note that all normal observed states are those whose masses are forbidden in the limit of exact $SU(2) \otimes U(1)$. Instead, for all SUSY partners the masses are allowed in that limit. Thus when supersymmetry is broken in the TeV range, but $SU(2) \otimes U(1)$ is intact only spartners take mass while all normal particles remain massless. Only at the lower weak scale the masses of ordinary particles are generated. Thus a simple criterion exists to understand the difference between particles and sparticles.

The other main avenue is compositeness of some sort. The Higgs boson is not elementary but either a bound state of fermions or a condensate, due to a new strong force, much stronger than the usual strong interactions, responsible for the attraction. A plethora of new "hadrons", bound by the new strong force, would exist in the LHC range. A serious problem for this idea is that nobody so far has been able to build up a realistic model along these lines, which could eventually be explained by a lack of ingenuity on the theorists side. The most appealing examples are technicolor theories[36],[37]. These models where inspired by the breaking of chiral symmetry in massless QCD induced by quark condensates. In the case of the electroweak breaking new heavy techniquarks must be introduced and the scale analogous to Λ_{QCD} must be about three orders of magnitude larger. The presence of such a large force relatively nearby has a strong tendency to clash with the results of the electroweak precision tests[38]. Another interesting idea is to replace the Higgs by a $t\bar{t}$ condensate[47]. The Yukawa coupling of the Higgs to the $t\bar{t}$ pair becomes a four-fermion $\bar{t}t\bar{t}t$ coupling with the corresponding strength. The strong force is in this case provided by the large top mass. At first sight this idea looks great: no fundamental scalars, no new states. But, looking closely, the advantages are largely illusory. First, in the SM the required value of m_t is too large: $m_t \geq 220$ GeV or so. Also a tremendous fine-tuning is required, because m_t would naturally be of the order of M_{GUT} or M_{Pl} if no new physics is present (the hierarchy problem in a different form!). Supersymmetry could come to the rescue in this case also. In a minimal SUSY version the required value of the top mass is lowered[48], $m_t \sim 195 \sin\beta$ GeV. But the resulting theory is physically indistinguishable from the MSSM with small $\tan\beta$, at least at low energies[49]. This is because a strongly coupled Higgs looks the same as a $t\bar{t}$ pair.

The hierarchy problem is certainly not the only conceptual problem of the SM. There are many more: the proliferation of parameters, the mysterious pattern of fermion masses and so on. But while most of these problems can be postponed to the final theory that will take over at very large energies, of order M_{GUT} or M_{Pl}, the hierarchy problem arises from the instability of the low-energy theory and requires a solution at relatively low energies. A supersymmetric extension of the SM provides a way out that is well defined, computable and that preserves all virtues of the SM. The necessary SUSY breaking can be introduced through soft terms that do not spoil the stability of scalar masses. Precisely those terms arise from supergravity when it is spontaneoulsly broken in a hidden sector[50]. But alternative mechanisms of SUSY breaking are also being considered[94]. As we shall now discuss, there are also experimental and phenomenological hints that point in this direction.

At present the most important phenomenological evidence in favour of supersymmetry is obtained from the unification of couplings in GUTs. Precise LEP data on $\alpha_s(m_Z)$ and $\sin^2\theta_W$ confirm what was already known with less accuracy: standard one-scale GUTs fail in predicting $\sin^2\theta_W$ given $\alpha_s(m_Z)$ (and $\alpha(m_Z)$), while SUSY GUTs[51] are in agreement with the present, very precise, experimental results. According to the recent analysis of ref. [52], if one starts from the known values of $\sin^2\theta_W$ and $\alpha(m_Z)$, one finds for $\alpha_s(m_Z)$ the results:

$$\begin{aligned} \alpha_s(m_Z) &= 0.073 \pm 0.002 \qquad \text{(Standard GUTS)} \\ \alpha_s(m_Z) &= 0.129(+0.010, -0.008) \qquad \text{(SUSY GUTS)} \end{aligned} \tag{29}$$

to be compared with the world average experimental value $\alpha_s(m_Z) = 0.118(5)$.

A very elegant feature of the GUT-extended supersymmetric version of the SM is that the occurrence of the $SU(2) \otimes U(1)$ electroweak symmetry breaking is naturally and automatically generated by the large mass of the top quark[53]. Assuming that all scalar masses are the same at the GUT scale, the effect of the large Yukawa coupling of the top quark in the renormalization group evolution down to the weak energy scale, drives one of the Higgs squared masses negative (that Higgs which is coupled to the up-type quarks). The masses of sleptons and of the Higgs coupled to the down-type quark are much less modified, while the squark masses are increased due to the strongly interacting gluino exchange diagrams. The negative value of the squared mass corresponds to the onsetting of the electroweak symmetry breaking. That the correct mass for the weak bosons is obtained as a result of the breaking implies constraints on the model, more stringent if no fine-tuning is allowed to a given level of accuracy. Various fine-tuning criteria have been analysed in the literature[54],[55]. Typically no more than a factor 10 fine tuning is allowed. With this assumption and realistic values of m_t one obtains the bounds shown in fig. 9[56]. These upper bounds give a quantitative specification of the constraints implied by a natural solution of the hierarchy problem in the context of the GUT-extended MSSM. They look very promising for LEP2 (but the bounds scale with the inverse square root of the fine-tuning factor...).

Many of the simpler GUTs predict the unification at M_{GUT} of the b and τ Yukawa couplings, or, equivalently, that for the running masses $m_b(M_{GUT}) = m_\tau(M_{GUT})$ [57]. The observed difference of the b and τ masses arises from the evolution due to the different interactions of quarks and leptons. Many authors studied the combined constraints from coupling unification and b and τ Yukawa unification[58]. The result is that there are a small $\tan\beta$ solution (typically in the range $\tan\beta$ =0.5–3) and a large $\tan\beta$ solution (with $\tan\beta$ =40–60). However the large $\tan\beta$ solution is somewhat disfavoured by a natural implementation of the electroweak symmetry breaking, according to the mechanism discussed above. In fact at large values of $\tan\beta \geq m_t/m_b$, the dominance of the top over the bottom Yukawa coupling, which is an important ingredient for that mechanism, is erased or even inverted. A closer look at the small $\tan\beta$ solution shows that the top mass is close to its fixed-point solution $m_t \sim 195 \sin\beta$ GeV so that $m_t \sim 175$ GeV corresponds to $\tan\beta \sim 2$. Correspondingly the mass of the lightest Higgs is relatively small[59], as discussed in sect. 9, which is good for LEP2.

In the MSSM the lightest neutralino is stable and provides a very good cold dark matter candidate. It is interesting that if the constraint $\Omega = 1$, which corresponds to the critical density for closure of the Universe, is added to the previous ones, consistency can still be achieved in a sizeable domain of the parameter space[58]–[60].

In conclusion, gauge coupling unification, natural $SU(2) \otimes U(1)$ electroweak symmetry breaking, b and τ Yukawa unification and a plausible amount of dark matter all fit together for realistic values of m_t, $\alpha_s(m_Z)$ and m_b. SUSY GUTs, with a single step of symmetry breaking from m_{GUT} down to m_W, appear to work well.

PRECISION ELECTROWEAK TESTS AND THE SEARCH FOR NEW PHYSICS

We now concentrate on some well-known extensions of the SM, which not only are particularly important per se but also are interesting in that they clearly demonstrate the constraining power of the present level of precision tests.

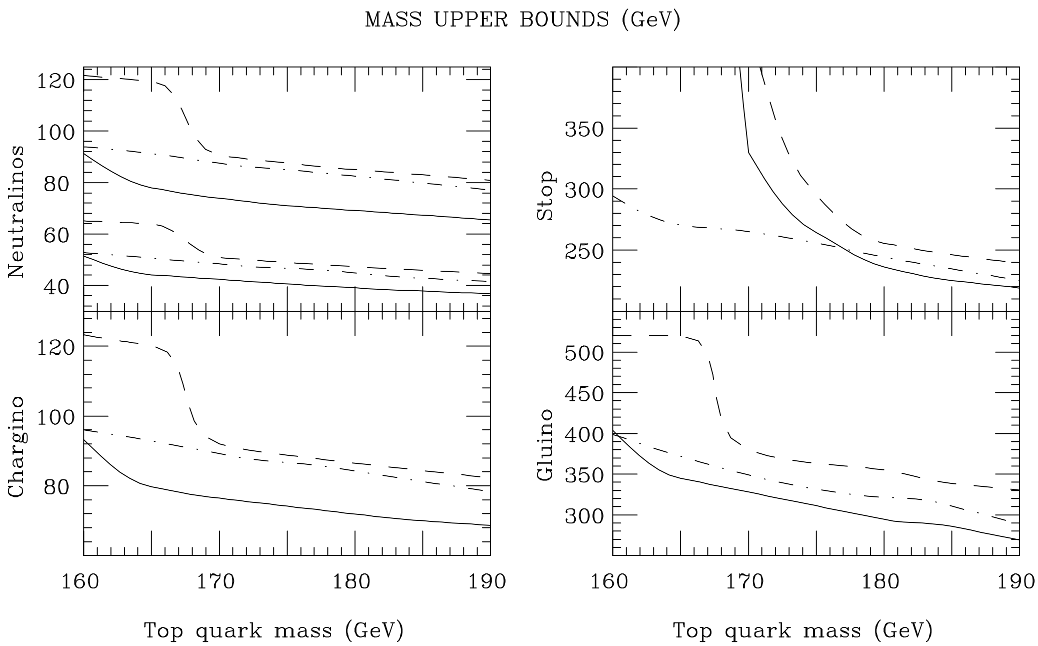

Figure 9

Upper bounds on gluino, lightest and next-to-lightest neutralino, and lightest chargino and stop masses based on the requirement of no fine tuning larger than 10%. The solid (dashed) lines refer to the minimal supersymmetric standard model with universal boundary conditions at M_{GUT} for the soft supersymmetry-breaking terms, without (with) the inclusion of the one-loop effective potential. The dot-dashed lines show the mass upper limits, for non-universal boundary conditions at M_{GUT}, without the includion of the one-loop effective potential.

Minimal Supersymmetric Standard Model

The MSSM[46] is a completely specified, consistent and computable theory. There are too many parameters to attempt a direct fit of the data to the most general framework. So one can consider two significant limiting cases: the "heavy" and the "light" MSSM.

The "heavy" limit corresponds to all sparticles being sufficiently massive, still within the limits of a natural explanation of the weak scale of mass. In this limit a very important result holds[61]: for what concerns the precision electroweak tests, the MSSM predictions tend to reproduce the results of the SM with a light Higgs, say $m_H \lesssim 100$ GeV.

In the "light" MSSM option, some of the superpartners have a relatively small mass, close to their experimental lower bounds. In this case the pattern of radiative corrections may sizeably deviate from that of the SM. The most interesting effects occur in vacuum polarization amplitudes and/or the $Z \to b\bar{b}$ vertex and therefore are particularly suitable for a description in terms of the epsilons (because in such a case, as explained in ref. [23], the predictions can be compared with the experimental determination of the epsilons from the whole set of LEP data). They are:

i) a threshold effect in the Z wave function renormalization[61], mostly due to the vector coupling of charginos and (off-diagonal) neutralinos to the Z itself. Defining the vacuum polarization functions by $\Pi_{\mu\nu}(q^2) = -ig_{\mu\nu}[A(0) + q^2 F(q^2)] + q_\mu q_\nu$ terms, this is a positive contribution to $\epsilon_5 = m_Z^2 F'_{ZZ}(m_Z^2)$, the prime denoting a derivative with respect to q^2 (i.e. a contribution to a higher derivative term not included in the usual S, T, U formalism). The ϵ_5 correction shifts ϵ_1, ϵ_2 and ϵ_3 by $-\epsilon_5$, $-c^2\epsilon_5$ and $-c^2\epsilon_5$ respectively, where $c^2 = \cos^2\theta_W$, so that all of them are reduced by a comparable amount. Correspondingly all the Z widths are reduced without affecting the asymmetries. This effect falls down particularly fast when the lightest chargino mass increases from a value close to $m_Z/2$. Now that we know, from the LEP1.5 and LEP2 runs, that the chargino mass is not so light, its possible impact is drastically reduced.

ii) A positive contribution to ϵ_1 from the virtual exchange of the scalar top and bottom superpartners[62], analogous to the contribution of the top–bottom left-handed quark doublet. The needed isospin splitting requires one of the two scalars (in the MSSM the stop) to be light. From the value of m_t, not much space is left for this possibility. If the stop is light then it must be mainly a right-handed stop.

iii) A negative contribution to ϵ_b, due to the virtual exchange of a charged Higgs[63]. If one defines, as customary, $\tan\beta = v_2/v_1$ (v_1 and v_2 being the vacuum expectation values of the Higgs doublets giving masses to the down and up quarks, respectively), then, for negligible bottom Yukawa coupling or $\tan\beta \geq m_t/m_b$, this contribution is proportional to $m_t^2\ /\tan^2\beta$.

iv) A positive contribution to ϵ_b due to virtual chargino–stop exchange[64], which in this case is proportional to $m_t^2/\sin^2\beta$ and prefers small $\tan\beta$. This effect again requires the chargino and the stop to be light in order to be sizeable.

v) A positive contribution to ϵ_b due to virtual h and A exchange[65], provided that $\tan\beta$ is so large that the Higgs couplings to the b quarks are as large or more than to the t quark.

If there is really an excess in R_b, it could be explained by either of the last two mechanisms[66]. For small $\tan\beta$ there is a positive contribution to R_b [63]–[77] if charginos and stops are light and the charged Higgs is heavy. Not to spoil the agreement for $\epsilon_1 = \Delta\rho$, we need the right stop to be light, while the left stop and the sbottom are kept heavy and near to one another, which is quite possible. Alternatively, for $\tan\beta$ large, of order 30 to 60, if h and A, the two neutral Higgses that can be lighter than the Z, are particularly light, then one also obtains[65] a substantial positive contribution to R_b. The large $\tan\beta$ value is needed in order to have a large coupling to $b\bar{b}$. However, such large values of $\tan\beta$ are somewhat unnatural. Also in this case having light charginos and stop helps.

Technicolor

It is well known that technicolor models[36]–[38] tend to produce large and positive corrections to ϵ_3. From ref. [38] where the data on ϵ_3 and ϵ_1 are compared with the predictions of a class of simple versions of technicolor models, one realizes that the experimental errors on ϵ_3 are by now small enough that these models are hopelessly disfavoured with respect to the SM.

More recently it has been shown[78] that the data on ϵ_b also produce evidence against technicolor models. The same mechanism that in extended technicolor generates the top quark mass also leads to large corrections to the $Z \to b\bar{b}$ vertex that have the wrong sign. For example, in a simple model with two technidoublets ($N_{TC} = 2$), the SM prediction is decreased by the amount[78],[79]:

$$\Delta\epsilon_b = -28 \times 10^{-3} \left| \frac{\xi}{\xi'} \left(\frac{m_t}{174 \text{ GeV}} \right) \right| \tag{30}$$

where ξ and ξ' are Clebsch-like coefficients, expected to be of order 1. The effect is even larger for larger N_{TC}. In a more sophisticated version of the theory, the so-called "walking" technicolor[79], where the relevant coupling constants walk (i.e. they evolve slowly) instead of running, the result is somewhat smaller[80] but still gigantic. Later it was shown[81] that in order to avoid this bad prediction one could endow the extended technicolor currents with a non-trivial behaviour under the electroweak group.

In conclusion, it is difficult to really exclude technicolor because it is not a completely defined theory, and no realistic model could be built so far out of this idea. Yet, it is interesting that the most direct realizations tend to produce $\Delta\epsilon_3 \gg 0$ and $\Delta\epsilon_b \gg 0$ which are both disfavoured by experiment.

OUTLOOK ON THE SEARCH FOR NEW PHYSICS

As we have seen in the previous sections, the whole set of electroweak tests is at present quite consistent with the SM. The pattern of observed pulls shown in table 1 accurately matches what we expect from a normal distribution of measurement errors. Even the few hints of new physics that so far existed have now vanished: R_c is back to normal and R_b is closer to the SM prediction. We no longer need new physics to explain R_b. Even the faint indication that $\alpha_s(m_Z)$ would prefer an excess in R_b has disappeared. Of course it is not excluded that a small excess of R_b is indeed real. For example the chances of nearby SUSY have not really been hit. Actually, with the absence of chargino signals at LEP1.5 and LEP2, which implies an increase of the lower bound on the chargino mass, the most plausible range for a possible effect on R_b in the MSSM is bounded within $\sim 1\sigma$ or $\sim 1.5\sigma$ of the ALEPH result (or $R_b \leq 0.2175$–0.2180) [66].

What is the status of other possible signals of new physics? The ALEPH multijet signal at LEP1.5[82] awaits confirmation from LEP2 before one can really get excited. So far no such convincing confirmation has been reported from the first $\sim$ 10 pb^{-1} of integrated luminosity collected at LEP2 at $\sqrt{s}$ = 161 GeV. The ALEPH multijet signal[82], if real, cannot be interpreted in the MSSM. But it could be a signal of some more unconventional realization of supersymmetry (e.g. with very light gluinos[83] or, more likely, with R-parity breaking [84]). It is perhaps premature to speculate on these events: in a few months we will know for sure if they are real or not, as soon as LEP2 will collect enough luminosity.

The CDF excess of jets at large transverse energy is not very convincing either[85]. It is presented as an excess with respect to the QCD prediction. But the QCD prediction can be to some extent forced in the direction of the data by modifying the parton densities, in particular the gluon density. At the price of a somewhat unnatural shape of the gluon density one can sizeably reduce the discrepancy without clashing with other data[86]. On the contrary this is not the case for the quark densities, which are tightly constrained by deep inelastic scattering data in the same x range[87]. Also the newly released D0 data do not show any additional evidence for the effect[88]. However, the D0 results are less accurate. Thus on the one hand one can say that D0 is compatible with either QCD or CDF. On the other hand their data are flat so that, to explain the absence of the signal, one should imagine a cancellation between the effect and the variation of systematics with E_T. It was pointed out in refs. [89],[90] that if the effect was real it could be explained in terms of a new vector boson Z' of mass around 1 TeV coupled mainly to quarks rather than leptons. In the presence of simultaneous anomalies in R_b, R_c and the jet yield at large E_T, it was attractive to present a unique explanation for all three effects. Now if only the jet excess is what remains this solution has lost most of its appeal. But in principle it is still possible to reduce the mixing of the Z' to the ordinary Z in such a way that its effect is only pronounced for jets while it remains invisible at LEP[92].

It is representative of the present situation that perhaps the best hint for new physics in the data is given by the single CDF event with $e^+e^-\gamma\gamma \not{E}_T$ in the final state[91]. Indeed this event is remarkable and it is difficult to imagine a SM origin for it. It is true that it is easier to imagine an experimental misinterpretation of the event (e.g. a fake electron, two events in one or the like) than a SM process that generates it. But it is a single event and even an extremely unlikely possibility can occur once. Several papers have already been devoted to this event[93]. In SUSY models two main possibilities have been investigated. Both interpret the event as a selectron pair production followed by decays $\tilde{e} \to eN'$,$N' \to N\gamma$. The observed production rate and the kinematics demand a selectron around 100 GeV and large branching ratios. In the first interpretation, within the MSSM, N' and N are neutralinos. In order to make the indicated modes dominant one has to restrict to a very special domain of the parameter space of the model. Neutralinos and charginos in the LEP2 range are then favoured. The second interpretation is based on the newly revived alternative approach in which SUSY breaking is mediated by ordinary gauge rather than gravitational interactions[46],[94]. In the most familiar approach of the MSSM, SUSY is broken in a hidden sector and the scale of SUSY breaking is very large, of order $\Lambda \sim \sqrt{G_F^{-1/2} M_P}$, where M_P is the Planck mass. But since the hidden sector only communicates with the visible sector through gravitational interactions the splitting of the SUSY multiplets is much smaller, in the TeV energy domain, and the goldstino is practically decoupled. In the alternative scenario the (not so much) hidden sector is connected to the visible one by ordinary gauge interactions. As these are much stronger than the gravitational interactions, Λ can be much smaller, as low as 10–100 TeV. It follows that the goldstino is very light in these models (with mass of order or below 1 eV typically) and is the lightest, stable SUSY particle, but its couplings are observably large. Then, in the CDF event, N' is a neutralino and N is the goldstino. The signature of photons comes out more naturally in this SUSY breaking pattern than in the MSSM. If the event is really due to selectron production it would be a manifestation of nearby SUSY that could be confirmed at LEP2. This is what we all wish. We shall see!

THE LEP2 PROGRAMME

The LEP2 programme started at the end of June'96. At first the energy was fixed at 161 GeV, which is the most favourable energy for the measurement of m_W from the cross-section for $e^+e^- \to W^+W^-$ at threshold. Then gradually the energy will be increased up to a maximum of about 193 GeV to be reached in mid '98. An average integrated luminosity of about 150 pb^{-1} per year is foreseen. LEP2 will run until the end of 1999 at least, before the shutdown for the installation of the LHC. The main goals of LEP2 are the search for the Higgs and for new particles, the measurement of m_W and the investigation of the triple gauge vertices WWZ and $WW\gamma$. A complete updated survey of the LEP2 physics is collected in two volumes[95].

An important competitor of LEP2 is the Tevatron collider. By and around the year 2000 the Tevatron will have collected about 1 fb^{-1} of integrated luminosity at 1.8–2 TeV. The competition is especially on the search of new particles, but also on m_W and the triple gauge vertices. For example, for supersymmetry while the Tevatron is superior for gluinos and squarks, LEP2 is strong on Higgses, charginos, neutralinos and sleptons.

Concerning the Higgs it is interesting to recall that the large value of m_t has important implications on m_H both in the minimal SM [96]–[98] and in its minimal supersymmetric extension[99],[100]. I will now discuss the restrictions on m_H that follow from the CDF value of m_t.

It is well known[96]-[98] that in the SM with only one Higgs doublet a lower limit on m_H can be derived from the requirement of vacuum stability. The limit is a function of m_t and of the energy scale Λ where the model breaks down and new physics appears. Similarly an upper bound on m_H (with mild dependence on m_t) is obtained[101] from the requirement that up to the scale Λ no Landau pole appears. The lower limit on m_H is particularly important in view of the search for the Higgs at LEP2. Indeed the issue is whether one can reach the conclusion that if a Higgs is found at LEP2, i.e. with $m_H \leq m_Z$, then the SM must break down at some scale $\Lambda > 1$ TeV.

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 for our purposes, 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 $\phi'(\mu) = \exp \int^t \gamma(t')dt'\phi$, 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:

$$\frac{d\lambda}{dt} = \frac{3}{4\pi^2}[\lambda^2 + 3\lambda h_t^2 - 9h_t^4 + \text{gauge terms}] \, , \tag{31}$$

with the normalization such that at $t = 0, \lambda = \lambda_0 = m_H^2/2v^2$ and the top Yukawa coupling $h_t^0 = m_t/v$. We see that, for m_H small and m_t large, 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[97]:

$$m_H > 135 + 2.1\,[m_t - 174] - 4.5\,\frac{\alpha_s(m_Z) - 0.118}{0.006}\;. \tag{32}$$

Summarizing, we see from Eq. (32) that indeed for $m_t > 150$ GeV the discovery of a Higgs particle at LEP2 would imply that the SM breaks down at a scale Λ below M_{GUT} or M_{Pl}, smaller for lighter Higgs. Actually, for $m_t \sim 174$ GeV, only a small range of values for m_H is allowed, $130 < m_H < \sim 200$ GeV, if the SM holds up to $\Lambda \sim M_{GUT}$ or M_{Pl} (where the upper limit is from avoiding the Landau pole[101]). As is well known[96] the lower limit is not much relaxed, even if strict vacuum stability is replaced by some sufficiently long metastability. Of course, the limit is only valid in the SM with one doublet of Higgses. It is enough to add a second doublet to avoid the lower limit. A particularly important example of theory where the bound is violated is the MSSM, which we now discuss.

As is well known[46], in the MSSM there are two Higgs doublets, which implies three neutral physical Higgs particles and a pair of charged Higgses. The lightest neutral Higgs, called h, should be lighter than m_Z at tree-level approximation. However, radiative corrections[102] increase the h mass by a term proportional to m_t^4 and logarithmically dependent on the stop mass. Once the radiative corrections are taken into account the h mass still remains rather small: for $m_t = 174$ GeV one finds the limit (for all values of $\tan\beta$) $m_h < 130$ GeV [100]. Actually there are reasons to expect that m_h is well below the bound. In fact, if h_t is large at the GUT scale, which is suggested by the large observed value ot m_t and by a natural onsetting of the electroweak symmetry breaking induced by m_t, then at low energy a fixed point is reached in the evolution of m_t. The fixed point corresponds to $m_t \sim 195 \sin\beta$ GeV (a good approximate relation for $\tan\beta = v_{up}/v_{down} < 10$). If the fixed-point situation is realized, then m_h is considerably below the bound, as shown in Ref. 56.

In conclusion, for $m_t \sim 174$ GeV, we have seen that, on the one hand, if a Higgs is found at LEP the SM cannot be valid up to M_{Pl}. On the other hand, if a Higgs is found at LEP, then the MSSM has good chances, because this model would be excluded for $m_h > 130$ GeV.

For the SM Higgs, which plays the role of a benchmark, also important for a more general context, the LEP2 reach has been studied in detail. Accurate simulations have shown[95] that at LEP2 with 500 pb^{-1} per experiment, or with 150 pb^{-1} if the four experiments are combined, one can reach the 5σ discovery range given by $m_H \leq 82, 95$ GeV for $\sqrt{s} = 175, 192$ GeV respectively, and the 95% exclusion range $m_H \leq 83, 98$ GeV. On the basis of these ranges we understand why a few GeV make a lot of difference. With $\sqrt{s} = 175$ GeV there would be practically no overlap with the LHC, which, even in the most optimistic projections, cannot see the Higgs below $m_H =$ 80 GeV or so. With $\sqrt{s} = 185$ GeV, there starts to be some overlap, but only limited to $m_H \leq$ 85–90 GeV, which is still a very difficult, time-consuming and debatable range for the LHC. With $\sqrt{s} = 195$ GeV there is already a quite reasonable overlap, up to $m_H \leq$ 95–100 GeV. The issue is not only that of avoiding a gap between LEP2 and the LHC, but also of providing an essential independent and complementary channel to study the new particle in a range of mass that is certainly rather marginal for the LHC.

In the MSSM a more complicated discussion is needed because there are several Higgses and the parameter space is multidimensional. Also, through the radiative corrections, the Higgs masses at fixed values of all MSSM parameters sensitively depend on the top quark mass. For decreasing top quark masses the upper bound on the light Higgs mass decreases. We note that the discovery range for LEP2 can be specified in terms of the light Higgs mass with little model dependence. On the contrary the same analysis for the LHC depends very much on the detailed quantitative pattern of the decay branching ratios. The usual plots that are seen in the experimental discussions are based on some typical choice of parameters, which is to some extent indicative.

In Ref. [95], the analysis for the MSSM is presented in great detail, as this case is rather complicated and was not deeply studied previously. With the typical choice of parameters, in the sense specified above, the domains of the $\tan\beta, m_A$ plane which are most difficult for the LHC are a "hole" at moderate values of $\tan\beta$ and m_A (say $\tan\beta < 10, m_A =$ 100–200 GeV) and a "strip" at small $\tan\beta$ and large m_A (typically $\tan\beta =$ 1–3 and $m_A > 300$ GeV). If m_t is not too small, these difficult regions can probably be covered at the LHC, but only with very large integrated luminosities $L = 3 \times 10^5$ pb^{-1}. LEP2 potentially can reduce the "hole" and completely cover the "strip", especially for m_t rather small. But while for $\sqrt{s} = 175$ GeV this is only true for rather extreme values of m_t and the squark mixing, at $\sqrt{s} = 192$ GeV only the central values are required (always with 150 pb^{-1} of integrated luminosity and the four experiments combined). Thus, as in the case of the SM, $\sqrt{s} = 192$ GeV is needed for a reasonable overlap, while less than that appears risky.

We now consider the search for supersymmetry. For charginos the discovery range at LEP2 is only limited by the beam energy for practically all values of the parameters. Thus every increase of the beam energy is directly translated into the upper limit in chargino mass for discovery or exclusion. For the Tevatron the discovery range is much more dependent on the position in parameter space. For some limited regions of this space, with 1 fb^{-1} of integrated luminosity, the discovery range for charginos at the Tevatron goes well beyond $m_\chi =$ 90–100 GeV, i.e. the boundary of LEP2, while in most of the parameter space one would not be able to go that far and only LEP2, with sufficient energy, would find the chargino.

The stop is probably the lightest squark. For a light stop the most likely decay modes are $\tilde{t} \to b\chi^+$ if kinematically allowed, otherwise $\tilde{t} \to c\chi$. A comparative study of these modes at LEP2 and at the Tevatron is presented in Ref. [95]. The result is that in either case at LEP2 the discovery range is up to about ($E_{beam} - 10$) GeV. At the Tevatron there is some difference between the two possible decay modes and some dependence on the position in the $\tilde{t}$–χ or the $\tilde{t}$–χ^+ planes, but it is true that very soon, at the end of the present run, with 100 pb^{-1}, a large region of the potential LEP2 discovery range will be excluded (in particular for the $\tilde{t} \to c\chi$ mode). Some limited regions will require more luminosity at the Tevatron and could be accessible to LEP2.

While on the stop the chances are better at the Tevatron than at LEP2 the converse is true for the sleptons. Here the Tevatron can only compete for a particularly favourable pattern of branching ratios. Finally, for neutralinos there is only a small region of the parameter space where these particles would be the first spartners to be discovered. The discovery ranges are very much parameter-dependent both at the Tevatron and at LEP2. For these reasons no detailed quantitative comparison still exists, although the channel $e^+e^- \to \chi\chi'$ has been extensively studied at the LEP2 workshop[95].

The measurement of m_W will be done at LEP2 from the cross-section at threshold and from direct reconstruction of the jet–jet final state in W decay. At present m_W is known with an error of ± 150 MeV from the direct measurement (see table 1). From the fit to all electroweak data one finds $m_W = 80352 \pm 34$ MeV (see eq. (10)), in agreement with the direct measurement. As a consequence the goal for LEP2 is to measure m_W with an accuracy $\delta m_W \leq \pm(30-40)$ MeV, in order to provide an additional significant check of the theory.

For the threshold method[95] the minimum of the statistical error is obtained for $\sqrt{s} = 2m_W + 0.5$ GeV $= 161$ GeV, which in fact was the initial operating energy of LEP2. The total error of this method is dominated by the statistics. If each of the four experiments will eventually collect 50 pb^{-1} of integrated luminosity (10 already collected and the rest in a possible future comeback at low energy) and the results are combined, then the statistical error will be $\delta m_W = \pm 95$ MeV and the total error $\delta m_W = \pm 108$ MeV. After ~ 10 pb^{-1} the present combined result is $m_W = (80.4 \pm 0.2 \pm 0.1)$ GeV [103]. Thus with realistic luminosity this method is not sufficient by itself.

In principle the direct reconstruction method can use the totally hadronic or the semileptonic final states $e^+e^- \to W^+W^- \to jjjj$ or $jjl\nu$. The total branching ratio of the hadronic modes is 49%, while that of the $\ell = e, \mu$ semileptonic channels is 28%. The hadronic channel has more statistics but could be severely affected by non-perturbative strong interaction effects: colour recombination among the jets from different W's and Bose correlations among mesons in the final state from WW overlap. Colour recombination is perturbatively small. But gluons with $E < \Gamma_W$ are important and non-perturbative effects could be relatively large, of the order of 10–100 MeV. Similarly for Bose correlations. One is not in a position to really quantify the associated uncertainties. Fortunately the direct reconstruction from the semileptonic channels can, by itself, lead to a total error $\delta m_W = \pm 44$ MeV, for the combined four experiments, each with 500 pb^{-1} of luminosity collected at $\sqrt{s} \geq 175$ GeV. Thus the goal of measuring m_W with an accuracy below $\delta m_W = \pm 50$ MeV can be fulfilled, and it is possible to do better if one learns from the data how to limit the error from colour recombination and Bose correlations.

The study of triple gauge vertices is another major task of LEP2. The capabilities of LEP2 in this domain are comparable to those of the LHC and go well below the level of deviations from the tree-level couplings that in the SM are expected from one-loop radiative corrections. LEP2 can push down the existing direct limits considerably. For given anomalous couplings the departures from the SM are expected to increase with energy. For the energy and the luminosity available at LEP2, given the accuracy of the SM established at LEP1, it is however not very likely, to find signals of new physics in the triple gauge vertices.

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