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ELECTROWEAK PHYSICS IN 1988

by

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Electroweak Physics in 1988 *

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Abstract

I review some salient features of the standard electroweak model and describe the impact of some recent experimental results on family issues and flavor mixing. Test of the standard model beyond tree level and prospects for detecting Higgs bosons are also discussed.

1 Introduction

In 1988, before the start up of SLC and LEP, the $SU(2) \times U(1)$ model of Glashow, Salam and Weinberg [1] continues to be the paradigm for describing the electroweak interactions. This does not mean that the model has remained static. Rather, in the last year, some issues have sharpened. Most notably, it now looks much more likely that there are only three families of quarks and leptons. At the same time, some other issues have become better defined, but in a sense more intriguing. For instance, radiative corrections appear to be on line with standard model expectations, provided that the top quark mass is not too large. Yet, the substantial flavor mixing seen in the B system argues for a sizable top mass. Finally, some issues continue to be as mysterious as ever, particularly those connected with the symmetry breaking sector of the theory. Given this state of affairs, in this talk I shall try to give a snapshot of the status of the standard electroweak theory, as it is today - a theory not fossilized, but one which is still in a settling process.

2 Three Families and No More?

In the standard electroweak model the fundamental fermions have repetitive $SU(2) \times U(1)$ properties: all left handed fields are in $SU(2)$ doublets, while their right handed counterparts are $SU(2)$ singlets. To date, with the exception of only one excitation (the top quark), we have evidence for three such "families" of quarks and leptons. Although we have no direct experimental evidence for the top quark, we have excellent indirect confirmation that there exist an $SU(2)$ partner for the left handed bottom quark and hence, inferentially, for the top quark. The best indication for this [2] comes from the measurement of the axial charge of b quarks in e^+e^- experiments at PEP and PETRA. The

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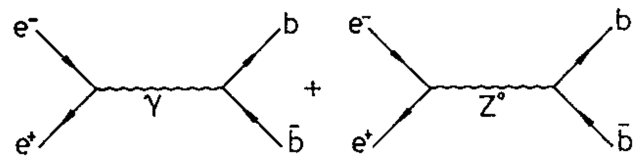

Figure 1: Diagrams contributing to $e^-e^- \to b\bar{b}$ in lowest order.

differential cross section for the process $e^+e^- \to b\bar{b}$ receives, in lowest order, contributions from both photon and Z exchange, as shown in Fig. 1. This latter graph gives rise to an asymmetry in the number of b ($\bar{b}$) quarks produced in the hemisphere along the direction of the incident e^- (e^+) in the CM system. The forward-backward asymmetry - with forward being along the direction of the incoming e^- - for b quarks can be well approximated, in the PEP/PETRA energy range, by retaining only the interference term between the photon and Z contributions and one finds:

$$A^b_{FB}(s) = \frac{9G_F a^b}{16\pi\sqrt{2}\alpha}\left[\frac{sM_Z^2}{M_Z^2 - s}\right] \tag{1}$$

where s is the square of the CM energy and a^b is the axial charge of the b quark. This charge takes different values, depending on the $SU(2)$ assignments of the b quark. For instance:

$$a^b = \begin{cases} -1 & b_L \textit{ doublet};\ b_R \textit{ singlet} \\ 0 & b_L \textit{ singlet};\ b_R \textit{ singlet} \\ 1 & b_L \textit{ singlet};\ b_R \textit{ doublet} \end{cases} \tag{2}$$

The forward-backward asymmetry for b-quarks has been measured by a number of experiments at PEP and PETRA, most notably by JADE [3], whose results are shown in Fig. 2. The b-quarks are identified through the large transverse momentum leptons they produce in their semileptonic decays. The values for the asymmetry obtained by the various PEP and PETRA experiments have been compiled by Greenshaw and Marshall [4] and reviewed by Wu [5]. The averaged value obtained [4]

$$a^b = -0.84 \pm 0.21 \tag{3}$$

needs to be corrected for $B - \bar{B}$ mixing, which dilutes somewhat the signal. This has been done by Wu [5], who arrives at a final value of

$$a^b = -1.08 \pm 0.29 \tag{4}$$

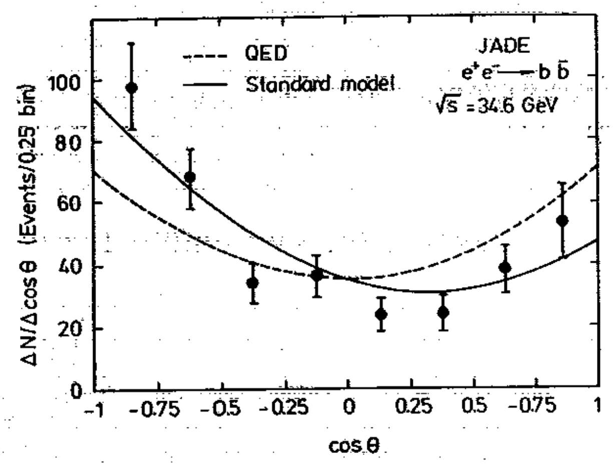

Figure 2: The measured angular distribution for the process $e^+e^- \to b\bar{b}$ measured by the JADE collaboration [3].

Although the errors are still large, this result is in perfect agreement with the supposition that the left handed b quark has a partner. So the top quark exists; the only question remaining is what is its mass!

There are two lower bounds on the top mass, one experimental and the other theoretical, which suggest that top is at least a factor of ten heavier than bottom[1]. The experimental bound on top comes from the UA1 collaboration, working at the Sp$\bar{p}$S collider [7]. In $p\bar{p}$ collisions top can be produced either as a byproduct of W decay, if its mass is below the $W \to t\bar{b}$ threshold, or in association with a $\bar{t}$, by gluon-gluon fusion. The first process has a well determined rate, which can be inferred from the experimentally measured W production and the known $W \to t\bar{b}$ branching fraction in the standard model. The rate for the second process, however, is more uncertain, since it depends on the knowledge of the gluon structure functions and of higher order QCD corrections. The signal for both these sources of top is an isolated muon with large transverse momentum, from the semileptonic decay of top, plus jets resulting from the accompanying debris. This signal, although characteristic, is not without background and a careful Monte Carlo study is needed to ascertain the presence of top.

Fig. 3, taken from [7], shows the expected signal, as a function of m_t, for the cuts imposed by the UA1 collaboration. As can be seen, for $m_t \geq 50$ GeV the main source of

[1] There is a safer, but smaller, experimental lower bound from TRISTAN [6], $m_t \geq 26.4$ GeV which requires very little theory input.

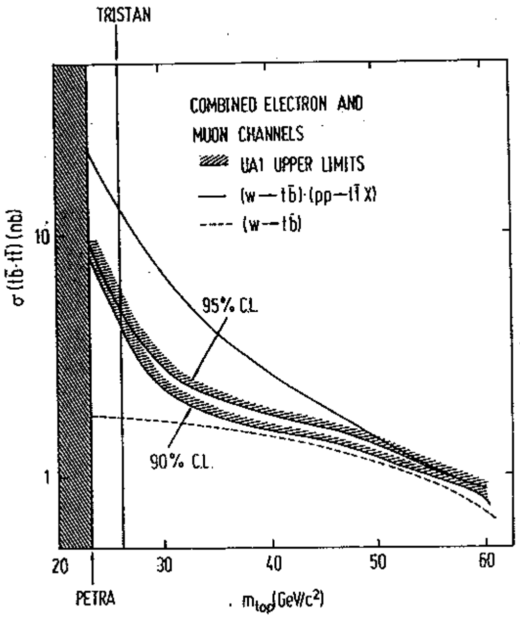

Figure 3: Expected top signal for the UA1 cuts [7], as a function of m_t. The dashed line is the signal from $W \to t\bar{b}$.

the signal is the gluon-gluon fusion process. The signal in Fig. 3 lies above the calculated background for $m_t \geq 56$ GeV (at 90% confidence) and this allows the setting of a lower bound on m_t. Because the predicted rate is dominated by the (somewhat) theoretically uncertain QCD process, this bound is not the most conservative one can set. UA1 sets a more cautious lower bound of $m_t \geq 44$ GeV [7] by varying the QCD inputs for the theoretically predicted top cross section and selecting the lowest value among these. A similar result has also been obtained recently by Altarelli et al. [8], who incorporated the effects of higher order QCD corrections for heavy quark production [9] and obtained $m_t \geq 41$ GeV. In summary, from the nonobservation of a clear signal in high energy $p\bar{p}$ scattering one can deduce - subject to the above mentioned theoretical ambiguities - an experimental lower bound for the top mass, in the range

$$m_t \geq (41 - 56)\ GeV \tag{5}$$

The observation of large $B_d - \bar{B}_d$ mixing by the ARGUS collaboration [10] allows one to infer, theoretically, a lower bound of a comparable magnitude for m_t to that given in (5). As I shall discuss this topic below in some detail, I will not comment further on this bound here. It is also possible, theoretically, to establish an upper bound for m_t, from the consistency of the calculations of electroweak radiative corrections with high precision neutral current experiments and with the W/Z mass measurements. The upper bound one obtains - which will be more fully examined in the next Section - lies in the range

$$m_t \leq (180 - 200)\ GeV, \tag{6}$$

depending on the value of the Higgs mass. The range bracketed by Eqs. (5) and (6) will be, hopefully, further reduced (or top will be found!) by the 1988/89 operation of the Tevatron and the upgraded CERN collider. These experiments should be sensitive to top masses up to around 80 GeV.

Given the reasonable certitude of the existence of top, it is natural to ask if there are more than three families of quarks and leptons. My personal prejudice is that the answer to this question is no. This is based on a, mostly unscientific, feeling that the masses of quarks and leptons should be bounded by roughly M_W and that, since top is already heavy, there is no room left. Experimentally, however, one is beginning to tackle this question more quantitatively. On the one hand, there are some lower bounds, from UA1 [7], on the masses of 4th generation leptons and charge $-\frac{1}{3}$ quarks.

$$m_{l'} \geq 41\ GeV \qquad (90\%\ C.L.) \tag{7}$$

$$m_{b'} \geq 32\ GeV \qquad (90\%\ C.L.) \tag{8}$$

On the other hand, rather sharp limits are beginning to emerge from e^+e^- experiments and from the CERN collider on the number of light neutrinos. These limits, which are becoming of comparable quality to the cosmological bound from nucleosynthesis [11] ($N_\nu \leq 4$), can be reasonably taken to be limits on the number of families. Although we do not really understand the origin of masses at all, it would seem to be a perverse world in which the fourth generation neutrinos would pick up a mass greater than $\frac{1}{2}M_Z$, while their lighter generation cousins are essentially massless!

PETRA and PEP have set limits on the number of neutrinos, N_ν, by looking for the process $e^+e^- \to \gamma\ Nothing$, where "*Nothing*" corresponds here to a neutrino-antineutrino pair. Since the Z couples universally to neutrinos of each generation, the rate for $e^+e^- \to \gamma\ Nothing$ is proportional to N_ν [2]. Even though in these experiments the CM energy is well below the Z mass, one obtains quite good limits on the number of neutrinos. The most sensitive results come from ASP at PEP and CELLO at PETRA, giving 90% C.L. limits of $N_\nu \leq 7.5$ [12] and $N_\nu \leq 8.7$ [13], respectively. An analysis of all e^+e^- data, done by the CELLO collaboration [13], gives a stronger combined limit of

$$N_\nu < 4.6 \qquad (90\%\ C.L.) \tag{9}$$

Neutrino counting limits from the $Sp\bar{p}S$ collider make use of an upper bound on the ratio R for the probability of producing W's and observing them in their $e\nu_e$ mode, to the probability of producing Z 's and observing them in their e^+e^- mode:

$$R = \frac{\sigma_W B(W \to e\nu_e)}{\sigma_Z B(Z \to e^+e^-)} \tag{10}$$

Experimentally, the UA1 and UA2 collaborations obtain the values $R = 9.1\ {}^{+1.7}_{-1.2}$ [14] and $R = 7.2\ {}^{+1.7}_{-1.2}$ [15], respectively, for this ratio, which gives an average [16]

$$R = 8.4\ {}^{+1.2}_{-0.9} \tag{11}$$

[2]This is not precisely correct, since for the process $e^+e^- \to \gamma\nu_e\bar{\nu}_e$ there is also an additional W-exchange contribution.

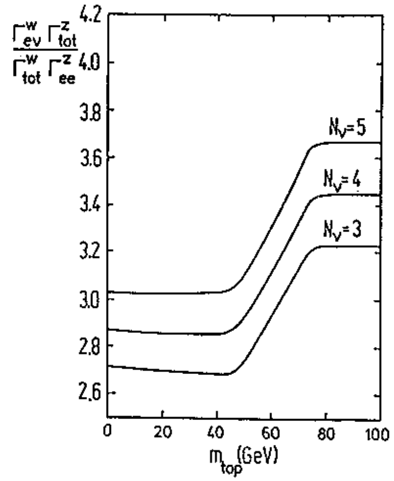

Figure 4: Branching fraction ratio as a function of m_t and N_ν.

and a 90% C.L. upper bound

$$R < 10.1 \qquad (90\%\ C.L.) \tag{12}$$

It is this upper bound that can be used to bound N_ν.

By definition, the branching fraction ratio is given by

$$\frac{B(W \to e\nu_e)}{B(Z \to e^+e^-)} = \frac{\Gamma(W \to e\nu_e)\ \Gamma^Z_{tot}}{\Gamma(Z \to e^+e^-)\ \Gamma^W_{tot}} \tag{13}$$

Both ratios above are calculable in the standard model. However, the ratio of the total widths depends on both N_ν and the top quark mass. Clearly Γ^Z_{tot} will go up if there are more neutrinos [3]. With the present limits on m_t , it is unlikely that the $t\bar{t}$ mode can contribute to Γ^Z_{tot}. However, it is quite possible that the decay $W \to t\bar{b}$ is kinematically allowed and, obviously, the lighter top is, the larger the W total width will be. The behaviour of the branching fraction ratio (13) is plotted in Fig. 4, as a function of m_t and the number of neutrinos.

To bound N_ν from the bound on R of Eq. (12), one has to know the production ratio σ_W/σ_Z of W's to Z's at the collider. This ratio can be estimated theoretically, but there are some uncertainties related to the structure functions one uses and to higher order QCD corrections [4]. A recent analysis by Colas, Denegri and Stubenrauch [16], which considers both BCDMS and EMC structure functions, gives for this ratio

$$\frac{\sigma_W}{\sigma_Z} = 3.25 \pm 0.10 \tag{14}$$

[3]Each extra neutrino species contributes approximately 170 MeV to the total Z width.

[4]These latter uncertainties, however, are not very large since one is dealing with a cross section ratio.

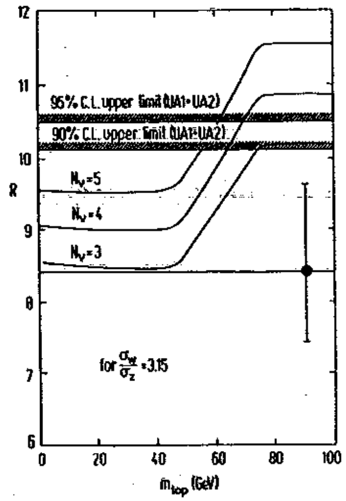

Figure 5: Bounds on N_ν as a function of m_t. The experimental point shown is the average UA1-UA2 result. From [16].

Using the lower 1σ value for the ratio σ_W/σ_Z above and the branching fractions plotted in Fig. 4, the 90% C.L. of Eq. (12) bounds $N_\nu \leq 5$, if $m_t \leq 70$ GeV, but for heavier m_t this bound is $N_\nu \leq 3$. These bounds are shown in Fig. 5. To keep the theoretical uncertainty in the above bounds in perspective, one should note that a 5% decrease in the 1σ lower limit for the ratio σ_W/σ_Z corresponds to an additional neutrino species. So, even though it is likely that there are only three families, it is not sensible to turn the above argument around and try to put a bound on m_t this way!

3 Probing Electroweak Radiative Effects

The measurements of electroweak parameters are now sufficiently precise that one can begin to test the presence of radiative corrections. The Weinberg angle Θ_W can be extracted from high precision ν_μ and $\bar{\nu}_\mu$ deep inelastic scattering data. In lowest order, if the scattering is done off an isoscalar target, this angle is directly related to the ratio R_ν of ν_μ neutral current scattering to ν_μ charged current scattering:

$$R_\nu = \{[\frac{1}{2} - sin^2\Theta_W^0] + \frac{5}{9} sin^4\Theta_W^0 [1+r]\} \tag{15}$$

where r is the ratio of $\bar{\nu}_\mu$ to ν_μ charged current scattering. A global analysis of all high precision experiments [CDHS, CHARM, CCFRR, FMM] done by Amaldi et al. [17], including all corrections, except electroweak radiative corrections, gives

$$sin^2\Theta_W^0 = 0.242 \pm 0.006 \tag{16}$$

Electroweak radiative corrections - adopting the convenient definition of $sin^2\Theta_W$ [18], $sin^2\Theta_W = 1 - \frac{M_W^2}{M_Z^2}$ - change the value (16) by an amount Δs^2. This shift depends very mildly on the (as yet unknown) Higgs mass M_H, as long as $M_H \leq$ TeV. Its dependence on m_t, however, is stronger and one has [19]

$$(\Delta s^2)_{top} \simeq 4.4 \times 10^{-4} [\frac{m_t}{M_W}]^2 \tag{17}$$

Using as canonical values $M_H = 100$ GeV and $m_t = 45$ GeV, one finds [17]

$$\Delta s^2 = -0.009 \pm 0.001 \tag{18}$$

and the result of the global analysis of Amaldi et al. [17] yields

$$(sin^2\Theta_W)_{DIS} = 0.233 \pm 0.003 \pm [0.005] \tag{19}$$

Here the last error in brackets is an estimate of the theoretical error of the whole analysis. Comparing (19) to (16), one sees that the radiative corrections have lowered the value of $sin^2\Theta_W$ by about 4%.

Radiative effects also change the predictions for M_W and M_Z. Using the same definition of the Weinberg angle as above [18], one has [20]:

$$M_W^2 = \frac{\pi\alpha}{\sqrt{2}G_F sin^2\Theta_W} [\frac{1}{1-\Delta r}] = \frac{(37.281 GeV)^2}{sin^2\Theta_W [1-\Delta r]} \tag{20}$$

where Δr is a theoretically calculated radiative correction. The measured values of the W and Z masses by the UA1 and UA2 collaborations, in principle, determine $sin^2\Theta_W$ directly, via its definition. However, the errors are such that this is not the most accurate value of $sin^2\Theta_W$ one can obtain from these measurements. Rather, it is better to extract $sin^2\Theta_W$ from Eq. (20) directly, imputing a value for the radiative correction Δr. This quantity, like Δs^2, depends very mildly on M_H, for $M_H \leq$ TeV, but more significantly on m_t [19], with the latter dependence going roughly as:

$$(\Delta r)_{top} \simeq -0.007 (\frac{m_t}{M_W})^2 \tag{21}$$

For the canonical values of M_H and m_t previously adopted, one finds [17] [20]

$$(\Delta r) = 0.0713 \pm 0.0013, \tag{22}$$

which provides a significant 7% change upwards in the value of $sin^2\Theta_W$ extracted from the W (and Z) masses.

Using the most recent values of the W and Z masses obtained by the UA1 and UA2 collaborations [21], Eq. (20) yields for $sin^2\Theta_W$ the value [17]

$$(sin^2\Theta_W)_{M_W/M_Z} = 0.228 \pm 0.007 \pm [0.002], \tag{23}$$

with the bracketed error again being theoretical. Several remarks are in order:

1. The radiative corrections help to bring the value of $sin^2\Theta_W$ obtained in deep inelastic scattering and the one gotten through the measurement of the W and Z masses in better agreement with each other. This can be seen in the Table below

Table I: Radiative effects for Θ_W

	$sin^2\Theta_W^0$	$sin^2\Theta_W$
Deep inelastic	0.242 ± 0.006	$0.233 \pm 0.003 \pm [0.005]$
W/Z Masses	0.212 ± 0.007	$0.228 \pm 0.007 \pm [0.002]$

2. A global fit of all neutral current experiments, of which the most precise are the ones we have discussed, finds a consistent value for $sin^2\Theta_W$. For the canonical choices of M_H and m_t adopted one has [17]

$$sin^2\Theta_W = 0.230 \pm 0.0048 \tag{24}$$

3. One cannot, however, allow the value of m_t to become too large since the corrections $(\Delta r)_{top}$ and $(\Delta s^2)_{top}$ of Eqs. (17) and (21) act in opposite direction to each other. Thus, if m_t is too big, the excellent agreement shown in Table I above will be spoiled. This phenomenon has been analyzed in detail by Amaldi et al. [17] and it provides a theoretical upper bound for m_t. This bound is in the range indicated in Eq. (6), with the uncertainty in the exact 90% C.L. limit obtained depending on the value of the Higgs mass, as shown in Fig. 6

4. One can extract a value of (Δr) by using the measurement of the W mass and the accurate value of $sin^2\Theta_W$ obtained in deep inelastic scattering. The value obtained

$$(\Delta r)_{exp} = 0.077 \pm 0.037 \tag{25}$$

is in nice agreement with the theoretical expectation (22), but is only 2σ away from zero. One will have to wait for the turn on of SLC and LEP to obtain a more accurate value. I note that the accuracy expected there [22], $\delta(\Delta r) \leq 0.01$, will begin to be sensitive to the effects of very heavy Higgs boson, since [19]

$$(\Delta r)_{Higgs} \sim 0.0024 \ln(\frac{M_H}{M_Z})^2 \tag{26}$$

4 Flavor Physics Constraints

1987 brought three important new experimental inputs for the physics of flavor:

1. Large mixing was seen in the B_d system [10]

2. The ARGUS collaboration reported [23] the first evidence for decays involving a direct $b \to u$ quark transition

3. The NA31 collaboration at CERN presented the first indication for $\Delta S = 1$ CP violation [24]

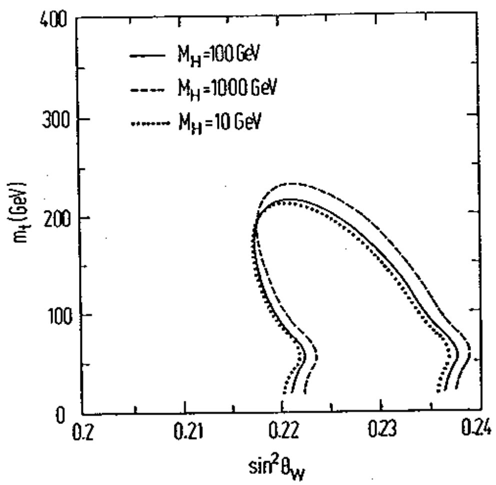

Figure 6: Allowed 90% C.L. regions in the $sin^2\Theta_W$-m_t plane, for different values of M_H. From [17].

In 1988 we are still digesting this information. Here I want to make some remarks on its theoretical impact [5]. I will assume in what follows, for the most part, that there are only three families of quarks and leptons since, although the data does not exclude more than 3 families, all data is consistent with the quark charged current weak interactions being described by a 3x3 unitary matrix V: the Cabibbo Kobayashi Maskawa matrix.

It is useful to adopt a parametrization of the CKM matrix first suggested by Maiani [27]. In this parametrization there is a natural hierarchy between the mixing angles and one can write, following Wolfenstein [28]:

$$sin\theta_1 = \lambda; \quad sin\theta_2 = A\lambda^2; \quad sin\theta_3 = A\rho\lambda^3, \tag{27}$$

where the new parameters A and ρ are of $0(1)$ and $\lambda \simeq 0.22$ is, essentially, the Cabibbo angle. To $O(\lambda^4)$, V takes the simple form:

$$V = \begin{vmatrix} 1-\frac{1}{2}\lambda^2 & \lambda & A\rho\lambda^3 e^{-i\delta} \\ -\lambda & 1-\frac{1}{2}\lambda^2 & A\lambda^2 \\ A\lambda^3(1-\rho e^{i\delta}) & -A\lambda^2 & 1 \end{vmatrix} \tag{28}$$

The parameter A is reasonably well fixed from measurements of the B lifetime. The parameter ρ and the CP violating phase δ, on the other hand, are still quite uncertain. However, their ranges are more restricted as a result of the new 1987 data. I will briefly summarize our present knowledge of all these quantities below.

[5]Some new experimental information became available at the time of the writing of this report. The CLEO collaboration at CESR confirmed the large mixing seen by ARGUS [25]. On the other hand, CLEO did not observe the decays $B_d \to p\bar{p}\pi(\pi)$ [26] and their 90% C.L. on these modes lies about a factor of 2 below the ARGUS observation [23].

Because $|V_{ub}|^2 << |V_{cb}|^2$, the B-lifetime and the B semileptonic branching ratios can be used to fix A. A typical recent analysis, by Altarelli and Franzini [29], gives

$$|V_{bc}|^2 = A^2\lambda^4 = \frac{(2.9 \pm 0.6) \times 10^{-13}}{\tau_B(10^{-12}sec)} \tag{29}$$

which, using the world average for τ_B [30], yields

$$A = 1.05 \pm 0.17 \tag{30}$$

That is, one knows A at the 10-20 % level.

Semileptonic B decays show, at present, no evidence for the direct quark transition $b \to u$. Hence one can use existing bounds on charmless semileptonic B decays to infer a bound on ρ. What is usually quoted are bounds on

$$R = \frac{\Gamma(b \to u)}{\Gamma(b \to c)} \simeq 2\frac{|V_{ub}|^2}{|V_{cb}|^2} = 2(\lambda\rho)^2, \tag{31}$$

where the factor of 2 above is a phase space factor. Using the bound from CLEO [31], $R \leq 0.08$ [6], this implies

$$\rho \leq 0.9 \tag{32}$$

The observation by ARGUS of charmless B decays [23]

$$\begin{aligned} BR(B^+ \to p\bar{p}\pi^+) &= (3.7 \pm 1.3 \pm 1.4) \times 10^{-4} \\ BR(B^0 \to p\bar{p}\pi^+\pi^-) &= (6.0 \pm 2.0 \pm 2.2) \times 10^{-4} \end{aligned} \tag{33}$$

provides the first evidence that $V_{ub} \neq 0$ and hence gives a lower bound for ρ. Unfortunately, translating an exclusive branching ratio into a value of V_{ub} is difficult. The ARGUS collaboration [23] gives a (very) conservative lower bound from (33) of $\frac{|V_{ub}|}{|V_{cb}|} \geq 0.07$, which implies

$$\rho > 0.3 \tag{34}$$

In fact, it is relatively easy to get a much bigger answer for ρ than (34) [7]. For present purposes, however, I will assume that ρ lies in the range indicated by Eqs. (32) and (34).

The parameter space allowed for the CP violating phase δ has been impacted by the ARGUS measurement of $B_d - \bar{B}_d$ mixing [10]. What was measured by ARGUS is the ratio of wrong sign dileptons in B_d decay:

$$r_d = \frac{\Gamma(B_d \to l^- X)}{\Gamma(B_d \to l^+ X)} = 0.21 \pm 0.08 \tag{35}$$

This ratio is related directly to the mixing parameter $x_d = (\frac{\Delta m}{\Gamma})_{B_d}$ via [8]

$$r_d \simeq \frac{x_d^2}{2 + x_d^2} \tag{36}$$

[6] This bound is rather conservative, as emphasized by [32]. Similar bounds have been also obtained by ARGUS and the XTAL BALL [33].

[7] For instance, Shifman's [34] first analysis of the ARGUS data suggested $\frac{|V_{ub}|}{|V_{cb}|} \geq 0.3$, which implies $\rho > 1.3$ - which is above the upper bound (32)! Thus the non observation by CLEO [26] of charmless B decays, at the level claimed by ARGUS, may help alleviate a minicrisis.

[8] This formula assumes that $\Delta\Gamma << \Gamma$, for the B_d system. Thus in (36) I have altogether neglected the contributions of $\Delta\Gamma$ relative to Δm.

Figure 7: Box graph contributing to x_d

The ARGUS result (35) implies $x_d = 0.78 \pm 0.18$. If one takes into account of the existing CLEO upper bound on r_d [35], the combined ARGUS and CLEO results yield a 90% C.L. range for x_d:

$$0.44 \leq x_d \leq 0.78 \tag{37}$$

Theoretically, the mixing parameter x_d in the Cabibbo Kobayashi Maskawa model arises from the box graph of Fig. 7, with the dominant contribution coming from the t quark loop. A straightforward evaluation of this graph yields for x_d a formula which is quadratically dependent on both m_t and V_{td} [9]

$$x_d \sim |V_{td}|^2 m_t^2 \tag{38}$$

Thus, one of the byproducts of the observation of a rather large $B_d - \bar{B}_d$ mixing by ARGUS [10] is that m_t cannot be too small. Indeed, soon after the announcement of the ARGUS result a flood of theoretical papers all, more or less, agreed that the range (37) necessitated $m_t \gtrsim 50$ GeV. However, Eq. (38) also makes it clear that the ARGUS measurement has implications on ρ and the phase δ, since

$$|V_{td}|^2 = A^2\lambda^6[1 + \rho^2 - 2\rho cos\delta] \tag{39}$$

Clearly, if m_t is not too big, one wants ρ and δ to be near $\rho_{max} \simeq 0.9$ and π, respectively, so as to maximize (39).

Although the large $B_d - \bar{B}_d$ mixing observed argues for the phase δ to be as close to π as possible, the observation of CP violation in the Kaon system requires some non trivial phase δ to exist. In collaboration with Krawczyk, Steger and London [36], I have performed recently a combined fit of x_d and the CP violating parameter ϵ in the Kaon system [10], to determine the allowed range for the parameters in the CKM matrix. Apart

[9] This is not quite correct, for large m_t.

[10] ϵ is proportional to $sin\delta$ and so it will vanish if $\delta \to \pi$.

from the existing experimental errors on x_d and ϵ, our analysis is uncertain theoretically since one has to estimate the hadronic matrix elements of the $\Delta B = 2$ and $\Delta S = 2$ quark amplitudes, entering in these quantities. In our work, to be conservative, we assumed a rather broad range of theoretical uncertainty in the calculation of the hadronic matrix elements. Furthermore, we let m_t vary over the full allowed theoretical range, from 40 GeV to 200 GeV. We obtained, in this way, regions in the $\rho - \delta$ plane, allowed by the present experimental results. These regions are in the shape of "moons", which fatten up as one allows the theoretical uncertainty to increase. Our results are shown in Fig. 8.

The lower left-hand figure in Fig. 8 perhaps most closely represents the present "best guesstimate" on the allowed $\rho - \delta$ range. It corresponds to a bag parameter $B_K = 2/3$ in the Kaon system and assumes $100\ MeV \leq (f_{B_d}^2 B_{B_d})^{\frac{1}{2}} \leq 200\ MeV$ in the B_d system. Furthermore, the top mass is varied over the presently permitted range. Note that if $\rho > 0.3$, δ is indeed rather large: $\delta > 2$. Thus the lower limit on the allowed value of ρ is important. However, it should be pointed out that a minimum value of $\rho > 0.15$ is already demanded by the requirement that $\epsilon \neq 0$, since $\epsilon \sim \rho$. Because of the theoretical uncertainties inherent in calculating x_d and ϵ, the standard model with three generations has still quite a large allowed region in the $\rho - \delta$ plane. A measurement of the top quark mass, however, would do much to reduce this range. This is shown in Fig. 9, where the allowed regions in the $\rho - \delta$ plane are detailed for fixed values of m_t, but otherwise allowing for the full range of the other theoretical uncertainties

The positive result obtained by the NA31 experiment at CERN for the CP violating parameter ϵ' in the Kaon system [24]

$$\frac{\epsilon'}{\epsilon} = (3.5 \pm 1.1) \times 10^{-3} \tag{40}$$

is also consistent with the expectations of the standard model. The importance of this result, however, is that:

1. It shows for the first time that $\Delta S = 1$ CP violating process exist, as predicted by the standard model. The ϵ parameter, measured in K decay, could be due to a purely $\Delta S = 2$, superweak transition [37]. ϵ', however, since it is proportional to the matrix element of H_{weak} between a Kaon and two pions, measures directly the amount of $\Delta S = 1$ CP violation. Thus the result (40) rules out all superweak theories, where CP violation occurs only through the mass matrix.

2. The NA31 result [24] does not further restrict the $\rho - \delta$ plane, because of additional theoretical uncertainties in estimating the hadronic matrix element needed to calculate ϵ'. However, the central range for ρ and δ, as well as the central range of the theoretical calculations of the relevant matrix elements, precisely give the central value in (40). Thus the measurement of ϵ' gives additional impetus to the notion that the three generation standard model is, essentially, correct.

There is a corollary to the last point which is perhaps worth noting. This concerns the strong CP problem [38], connected with why the effective CP violating parameter $\bar{\Theta} = \Theta + Arg\det M$, of the combined electroweak and QCD theory, is so small: $\bar{\Theta} \leq 10^{-9}$. It is possible that $\bar{\Theta}$ vanishes due to some additional high energy symmetry [39]. However, one can also imagine that $\bar{\Theta}$ is tiny because the quark mass matrix is essentially real and

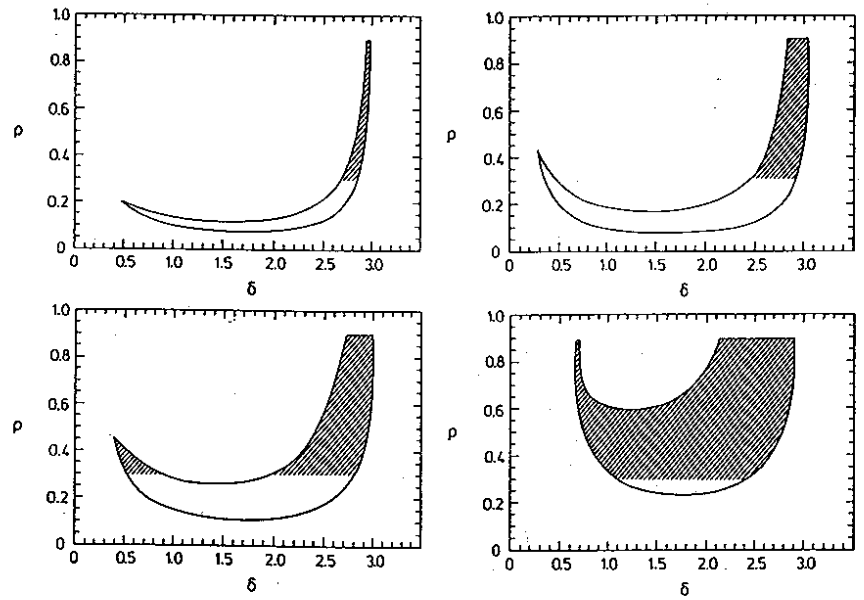

Figure 8: Allowed regions in the $\rho - \delta$ plane for different theoretical assumptions. The crosshatched area lies in the range of Eqs. (32) and (34). For more details, see [36]

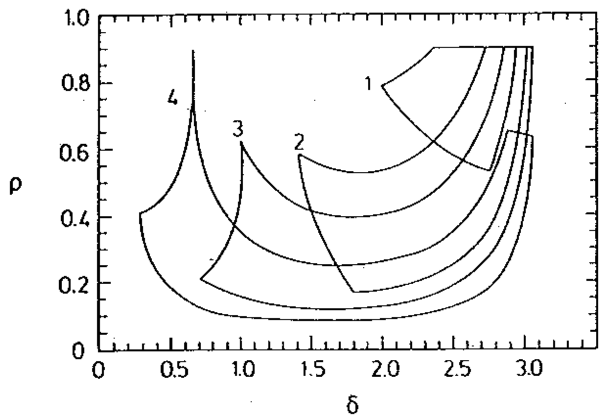

Figure 9: Allowed regions in the $\rho-\delta$ plane for fixed values of m_t, but for the full theoretical uncertainty. Regions 1-4 correspond to m_t = 60, 90, 120, 180 GeV, respectively. From [36]

CP is broken by the vacuum only [40]. In this latter case, the observed CP violation in the Kaon sector requires another explanation, than that provided by the Cabibbo Kobayashi Maskawa mixing. This necessitates, in general, already a fair bit of fine tuning to get ϵ right [41]. The requirement now to also get a non zero value for ϵ' - and of the right magnitude - make the problem much more difficult. Thus the NA31 result has important implications for alternative models of CP violation

5 The Higgs Enigma

The simplest way to break $SU(2) \times U(1)$ to $U(1)_{em}$ involves a complex scalar doublet Φ and an asymmetric potential

$$V = \lambda[\Phi^\dagger\Phi - \frac{1}{2}v^2]^2 \tag{41}$$

The Higgs boson H is the debris that remains in the theory, after the symmetry breakdown. Because the parameter λ in (41) is unknown, the mass of this excitation is not determined by the theory. There are, however, various theoretical arguments which provide both lower and upper bounds for M_H. Considerations related to the triviality of the $\lambda\Phi^4$ theory [42] have been much discussed in the last years and allow one to set a reliable upper bound for M_H of the order of a TeV [43]. If M_H is above this value, the simple symmetry breaking sector assumed is inconsistent, even as an effective field theory, since excitations appear above the physical cutoff necessary to define the theory. Although the triviality upper bound is theoretically reliable, it is also a long long way away from present experimental probing. Indeed, it will require a machine like the SSC before one can say much about the

Higgs, if its mass is near this bound!

The lower bounds on the Higgs mass which exist [44] are less theoretically reliable, since they are based on less firmly established principles (like demanding that radiative effects do not change the nature of the vacuum) and necessitate some constraints on the existing fermion spectrum. The typical value for these bounds, around 10 GeV, furthermore is, or will be soon, accessible experimentally. Thus it is perhaps more worthwhile here to discuss what are the present known experimental bounds on M_H and what are the prospects for improving them, in the near future.

Perhaps the most reliable of the (low mass) bounds on M_H was pointed out by Barbieri and Ericson [45] long ago. It comes from low energy neutron-nucleus scattering, where the presence of a very low mass Higgs alters the shape of the measured angular distribution beyond what is expected, unless

$$\frac{h^2_{HNN}}{M^4_H} \leq 4.3 \times 10^{-10} MeV^{-4} \tag{42}$$

Here h_{HNN} is the Higgs coupling to nucleons. This coupling can be inferred by current algebra methods and is proportional to the pion-nucleon σ-term [46]: $h_{HNN} = \frac{\sigma}{v}$, with $v = (\sqrt{2}G_F)^{-\frac{1}{2}} \simeq 250\ GeV$. Using the most recent determination of σ [47] ($\sigma \simeq 50$ MeV) yields

$$M_H \geq 3.1\ MeV \tag{43}$$

There is another low energy bound on Higgs bosons usually quoted, coming from the study of the nuclear deexcitation of the 20.1 MeV state of ^4He [48]. This bound also depends on the value of h_{HNN}. Unfortunately, the excluded region given by this experiment [48]: 3 MeV $\leq M_H \leq$ 14 MeV uses a value for $h_{HNN} \simeq 1.1 \times 10^{-3}$. Using the σ term evaluation above for h_{HNN} gives a much smaller value for this coupling and no bound is set by [48].

There are additional bounds on Higgs bosons which come from direct flavor changing transitions at the quark level ($s \to dH$; $b \to sH$), induced by Penguin operators. These transitions, which are shown in Fig. 10, are particularly important if m_t is large. The experimental branching ratio [49] $B(K^+ \to \pi^+e^+e^-) = (2.7 \pm 0.5) \times 10^{-7}$ is below the Penguin prediction for the decay $K^+ \to \pi^+H$; $H \to e^+e^-$ of Willey [50] which, for m_t = 45 GeV, is

$$B(K^+ \to \pi^+H; H \to e^+e^-) \simeq 7 \times 10^{-6} \tag{44}$$

This excludes Higgs bosons from the region

$$50MeV \leq M_H \leq 211MeV \tag{45}$$

where the lower limit is due to an experimental cut in [49] and the upper limit is where the process $H \to \mu^+\mu^-$ becomes dominant.

In the case of B mesons, the decay $B \to HX$ is dominated entirely by the t-quark Penguin and the rate is very large indeed for large m_t [51] [11]

$$B(B \to HX) \simeq 0.3(\frac{m_t}{M_W})^4[1 - \frac{M^2_H}{m^2_B}]^2 \tag{46}$$

[11] The extra factor of $[1 - \frac{M_H^2}{m_B^2}]$ in Eq. (46), comes from the decay amplitude. This additional kinematical suppression was pointed out recently by Bertolini, Borzumati and Masiero [52]

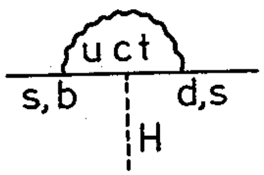

Figure 10: Penguin diagram for Higgs emission.

If $m_t \simeq M_W$, simple charm counting in B decays excludes all $M_H \leq M_B$. For lighter m_t's one must rely on particular exclusive modes and the bounds are more model dependent. In particular, to make use of the CLEO bound on the process $B \to K\mu^+\mu^-$ [54] of roughly 2×10^{-4} to bound M_H, requires an accurate estimate of the relative importance of the $H \to \mu^+\mu^-$ and $H \to \pi^+\pi^-$ modes.

Perhaps the strongest bound to date on Higgs bosons in the GeV range comes from the process $\Upsilon \to H\gamma$. The branching ratio for this process to that for $\Upsilon \to \mu^+\mu^-$ is calculable [55] but, unfortunately, it has very large QCD corrections [56]:

$$\frac{\Gamma(\Upsilon \to H\gamma)}{\Gamma(\Upsilon \to \mu^+\mu^-)} = [\frac{G_F^2 m_b^2}{\sqrt{2}\pi\alpha}(1 - \frac{M_H^2}{M_\Upsilon^2})][\frac{1 - \frac{40}{3\pi}\alpha_s F(\frac{M_H}{M_\Upsilon})}{1 - \frac{16}{3\pi}\alpha_s}] \tag{47}$$

In the above the factor in the first square bracket is the Wilczek result [55], while the second bracket contains the QCD corrections. Note that for light M_H the function F, which is given in detail in [56], is close to unity. Taking this formula at face value CUSB [57] obtains a 90% C.L. bound on M_H of

$$M_H \geq 4.3\ GeV \tag{48}$$

However, one must worry about the theoretical reliability of this result since the QCD corrections in (47) change the Wilczek value by 60%!

This somewhat somber situation with the present Higgs bounds will change dramatically with the turn on of SLC and LEP. LEP 100 and SLC should be able to discover the Higgs boson if $M_H \leq 45$ GeV, while LEP 200 can explore Higgses up to $M_H \simeq 80$ GeV. At the Z, the most favorable way to look for Higgs bosons appears to be the Bjorken process [58] $Z \to Hl^+l^-$, which has both a reasonable rate and a good signature. The process proceeds through the HZZ coupling and the presence of the virtual Z produces a dilepton pair which is highly peaked towards large masses. Because of the clean signature, one can

hope to see a signal even with only a few events [59], so that with 10^6 Z's one should be able to push the Higgs bound to about 45 GeV (or discover the Higgs!)

At LEP 200 one will look for Higgs bosons in associated production with a Z, where the process $e^+e^- \to ZH$, for Higgs bosons lighter than 60-70 GeV, has a rate of the order of pb. With an expected integrated luminosity of around 500 pb^{-1} per year, this process could produce of the order of 500 Higgs bosons. However, one will have to be clever to dig out the signal! A promising possibility [60] is to look for events in which the Z decays into neutrinos, which produces rather spectacular events with two unbalanced jets in one hemisphere, plus nothing in the other. Detailed Monte Carlo studies suggest that at LEP 200 a Higgs boson as heavy as 80 GeV will be detectable [60].

By 1995 it appears to me possible that we will know if a Higgs boson exists, with mass less than 80 GeV or so. If no such signal is seen, one will have to wait for the next generation of colliders (SSC, LHC, CLIC) to renew the hunt for the Higgs. Although quite clear strategies to pursue this more distant goal have been mapped out [61], their discussion here would take me too far afield.

I would like to end with a provocative remark. Although finding the Higgs will undoubtedly be an experimental triumph, it may well be a theoretical tragedy. The discovery of a single, light, Higgs will tell us that the asymmetric potential (41) is an adequate description for the symmetry breaking sector of the standard model. This is informative and nice, but it will not help theorists really to understand where this potential comes from! Indeed, if the Higgs is light enough ($M_H \leq 150 - 200$ GeV) the potential (41) is a correct effective description, up to a cutoff lying much above the Planck mass. There is no way to find out, therefore, what the true underlying theory for the symmetry breakdown might be. The enigma of symmetry breaking is perennially cloaked by the discovery of a light Higgs. So, please do not find the Higgs!

6 Conclusions

In 1988 the standard model of the electroweak interactions is in great shape! Let me recapitulate some of the points made, which allow me to be so bullish:

- Although there is still some room for additional families of quarks and leptons, one is getting close to assert that there are only three generations

- Top must exist and its mass range is reasonably well bounded between 50 and 200 GeV, with my own personal favorite value being $m_t \sim 80$ GeV.

- One is beginning to check the standard model radiative corrections. Indeed, for light top, these corrections are needed, at the 3σ level, to obtain agreement with experiment.

- Although there have been some exciting developments in flavor physics in the last year, the flavour structure looks still very compatible with the simple Cabibbo Kobayashi Maskawa mixing scheme.

- No sign of Higgs is yet in sight, but this is not a real worry since realistic bounds on M_H are of $O(1\ TeV)$.

This bullish state may or may not persist when the new high luminosity e^+e^- colliders at the Z will come in operation. Clearly, everyone is waiting eagerly for SLC and LEP and their potential for doing precision electroweak physics. Although I am also very curious to see what will happen at SLC and LEP, I would trade this information for the knowledge of what will be said - if anything! - about the standard electroweak model when the University of Bologna will celebrate its millesium!

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