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STATUS OF THE STANDARD MODEL

by

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STATUS OF THE STANDARD MODEL

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ABSTRACT

I illustrate by means of a variety of recent examples how well the standard SU(3)xSU(2)xU(1) model works. Among the topics discussed are: W and Z physics; some aspects of perturbative QCD; theoretical and experimental constraints and prospects for Higgs boson detection; and CP violation within the Kobayashi-Maskawa framework.

1. PREMISES

A report on the standard SU(3)xSU(2)xU(1) model of the strong and electroweak interactions these days can follow two roads. Either it tries to be encyclopedic, and details the extensive evidence that exists supporting the model in a variety of different physical contexts, or it picks and chooses some significant recent results, which exemplify again how well the model works. This report follows the second route. The examples I have chosen, for illustration, are a matter of personal taste, although I believe they fairly represent what might be considered highlights in the field this year. It will be noticed that no sharp distinction exists anymore between strong and weak interaction tests. Hadronic interactions are used to extract properties of the weak bosons and, conversely, the $p_\perp$ distributions of these bosons are used to test QCD. The only sad note, in all this unity, is that the standard model in 1985 works all too well! Alas, there is not much that can be done about this, until more detailed and higher energy experiments find some real traces of disagreement. In 1985, the anomalies of 1984 appear to have been only statistical fluctuations!

2. W AND Z PHYSICS

The first topic which I would like to discuss is the progress made in determining the properties of the weak intermediate bosons in the collider experiments at CERN.

2.1 W and Z Masses

The standard electroweak theory of Glashow, Salam and Weinberg /1/ makes precise predictions for the masses of the intermediate vector bosons, in terms of low energy parameters. It is convenient /2/ to adopt a definition of the Weinberg angle in terms of the W and Z masses:

$$\sin^2\theta_W = 1 - M_W^2/M_Z^2 \tag{1}$$

Then the W mass, including radiative corrections, can be expressed as /3/:

$$M_W^2 = \frac{\pi\alpha}{\sqrt{2}\, G_\mu \sin^2\theta_W (1-\Delta r)} \tag{2}$$

where G_μ is the Fermi constant determined from μ decay, in which certain electromagnetic contributions are explicitly included /4/. Numerically one has /4/ $G_\mu = (1.16638 \pm 0.00002)\times 10^{-5}\,\mathrm{GeV}^{-2}$. The Weinberg angle, defined by Eq. (1), can be extracted from radiatively corrected ν deep inelastic scattering, with the result /5/

$$\sin^2\theta_W = 0.217 \pm 0.014 \tag{3}$$

Finally, the quantity $(1-\Delta r)$ is theoretically calculable and has the value /6/, for $m_t = 40$ GeV, $m_H = m_Z$

$$1 - \Delta r = 0.9304 \pm 0.0020 \tag{4}$$

The dependence of (4) on m_t and m_H is mild and will be commented upon below.

Using the above results, and a precise value for the fine structure constant, gives for the W and Z masses the predictions

$$M_W = 83.0\,^{+2.9}_{-2.7}\ \mathrm{GeV}; \qquad M_Z = 93.8\,^{+2.4}_{-2.2}\ \mathrm{GeV} \tag{5}$$

The main error in Eqs. (5) is due to the error in $\sin^2\theta_W$ from Eq. (3). The most recent UA_1 and UA_2 values, as reported by Di Lella /7/, are in perfect agreement with the above predictions

$$\begin{aligned} M_W &= 83.1\,^{+1.3}_{-0.8} \pm 3 \ \mathrm{GeV} \quad UA_1 \\ M_W &= 81.2 \pm 1.1 \pm 1.3 \ \mathrm{GeV} \quad UA_2 \\ M_Z &= 93.0 \pm 1.6 \pm 3 \ \mathrm{GeV} \quad UA_1 \\ M_Z &= 92.5 \pm 1.3 \pm 1.5 \ \mathrm{GeV} \quad UA_2 \end{aligned} \tag{6}$$

(the first error is statistical, the second is systematic), but the errors in both the theoretical prediction and experiment are too large to test significantly the radiative corrections. I should note that calculating $\sin^2\theta_W$, using Eq. (1), from the collider data one cancels most of the systematic error. The UA_1 and UA_2 average /7/

$$\overline{\sin^2\theta_W} = 0.218 \pm 0.023 \tag{7}$$

agrees very nicely with that obtained from the low energy experiments and the precision of the determination is quite comparable to that in Eq. (3).

It is of obvious theoretical interest to eventually be able to check the radiative corrections embodies in Δr. It turns out, however, that most of the 7 % change in $(1-\Delta r)$ in Eq. (4) is theoretically rather trivial, coming essentially from the effects of the running of α to the W mass scale. The interesting variation in Δr, coming from the properties of the Higgs sector or the existence of widely split fermion doublets, is at the 1 % level. For instance /6/, changing m_H from 10^2 GeV to 10^3 GeV changes Δr by 0.009. To be able to test the electroweak theory to this level necessitates two very precise measurements (cf. Eqs. (1) and (2)). One of these will be provided by SLC and LEP, through the measurement of the Z^0 mass to better than one part per mil ($\delta M_Z / M_Z \sim 10^{-3}$). The other will require either a comparable measurement of the W mass, ($\delta M_W / M_W \sim 10^{-3}$), or an extremely accurate measurement of $\sin^2\theta_W$ ($\delta\sin^2\theta_W \sim 2.5\times 10^{-3}$) by measuring the forward-backward asymmetry at the Z^0 peak to 2 parts per mil. This latter measurement, although very difficult, appears to be feasible at LEP /8/. Interestingly enough, measuring M_W to 100 MeV also appears feasible at LEPII /9/. Let me briefly comment on how this may be done.

Four ways have been suggested /9/ for measuring the W mass in the process $e^+e^- \to W^+W^-$. They involve:

i) Measuring the threshold dependence of $\sigma(e^+e^- \to W^+W^-)$.
ii) Measuring the endpoint in the electron spectrum in $W \to e\nu$ decays.
iii) Measuring the jet-jet invariant mass arising from hadronic decays of the W.
iv) Measuring the $e\nu$ invariant mass in $W \to e\nu$ decays.

The statistical and systematic errors for all these four methods, in one year of running at LEPII, lie in the 100 MeV range. Let me illustrate this for case iv). Here one selects events in which one W decays hadronically and the other leptonically. Having determined the W axis from the jet-jet analysis, as shown in Fig. 1, then all the remaining kinematics is fixed. At LEPII, in contrast to the collider, $M_{e\nu}$ and not only the transverse mass $M^T_{e\nu}$ is determined. Fig. 2 shows a Monte Carlo reconstruction /9/ of $M_{e\nu}$ using an integrated luminosity of 100 pb^{-1}. The statistical error here is ± 55 MeV and there is a systematic shift of the input W mass of 80 MeV.

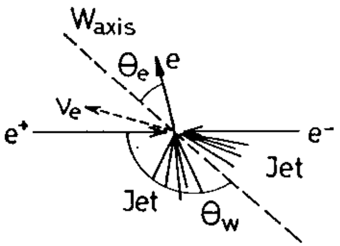

Fig. 1: Kinematic reconstruction of $M_W = M_{e\nu}$

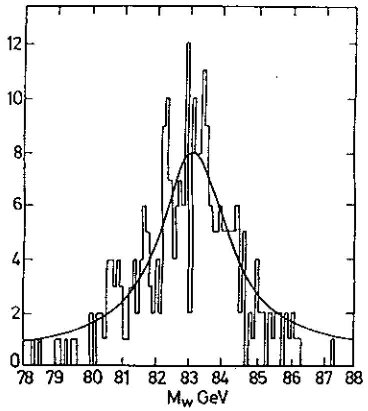

Fig. 2: Monte Carlo reconstruction of $M_{e\nu}$, from Ref. 9

2.2 Neutrino Counting

In the standard model, a precise measurement of the Z^o width gives a determination of the number of neutrino species and, inferentially, of the number of generations. Unfortunately, the present uncertainty in mass resolution, and the scant number of events, preclude a direct measurement of Γ_Z at the collider. Nevertheless, one can infer the number of neutrino species N_ν by using a bit of theory /10/. What is measured at the collider is the ratio

$$R = \frac{\sigma_Z \; B(Z\to e^+e^-)}{\sigma_W \; B(W\to e\nu)} \tag{8}$$

of the production of Z's and W's, multiplied by their decay branching ratio into electrons. The average value for R determined by UA_1 and UA_2 is /7/

$$\bar{R} = 0.125 \pm 0.023 \tag{9}$$

and at 90 % confidence level R > 0.096. Let me display the explicit dependence of R on N_ν. One has

$$R = \frac{\sigma_W}{\sigma_Z} \; \frac{\Gamma(Z^o\to e^+e^-)}{\Gamma(W\to e\nu)} \; \frac{\Gamma_W^{st}}{\Gamma_Z^{st} + (N_\nu - 3)\;\Gamma(Z\to\bar{\nu}\nu)} \tag{10}$$

Using some theoretical imput, therefore, one may compute N_ν from the measured value of R. The cross section ratio can be calculated in QCD rather accurately /11/ (see below) and one finds $\sigma_W/\sigma_Z = 0.30 \pm 0.02$. The ratio of leptonic widths is also well known since it follows from the low energy couplings of the Z^o and W ($\Gamma_Z/\Gamma_W \simeq 0.37$). Finally, the total widths Γ_W^{st}, Γ_Z^{st} in the standard model do depend on m_t but one has, nearly, $\Gamma_W^{st} \simeq \Gamma_Z^{st}$. It is obvious, therefore, from (10) and the experimental result (9) that there is not much room for extra neutrinos.

To be specific, using the choice of parameters of Deshpande et al. /10/ (M_W = 83 GeV, M_Z = 94 GeV, $\sin^2\theta_W = 0.22$, m_t = 40 GeV) one has $\Gamma(Z^o\to e^+e^-)/\Gamma(W\to e\nu) = 0.368$ and Γ_W = 2.82 GeV, Γ_Z = 2.83 GeV. Then the collider measurement of R /7/ implies $N_\nu = 1.3 \pm 2.7$, while using the 90 % confidence limit on R one has

$$N_\nu < 5.4 \pm 1.0 \tag{11}$$

Given the very little amount of theory imputed, this is quite impressive. The collider is already catching up with well known nucleosynthesis bound /12/ $N_\nu \simeq 3 - 4$!

2.3 Weak Boson Production and QCD

The total production cross section, as well as the $P_\perp$ distribution of the weak bosons, at the collider can be reliably computed in perturbative QCD. To lowest order in α_s the weak boson production proceeds by quark - antiquark annihilation. In $O(\alpha_s)$, however, both produced gluons and quark-gluon processes must be included, as shown schematically in Fig. 3 /13/.

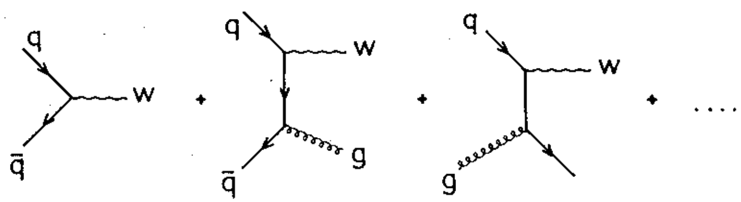

Fig. 3: Lowest order contributions to W production in hadronic collisions

Using the full theoretical machinery of perturbative QCD, prediting σ_W and σ_Z has passed from being an ancient and honorable art /14/ to a science /11/. To a very good approximation, one finds that one can include the $O(\alpha_s)$ corrections by just multiplying the lowest order parton result by an overall factor (K-factor).

Schematically, one can write for the production cross sections

$$\sigma_{W/Z} = K(Q^2)\, N_{W/Z} \int dx_1 dx_2\, \delta(x_1 x_2 - \tau)\, \{q(x_1,M^2)\bar{q}(x_2,M^2) + 1 \leftrightarrow 2\} \quad (12)$$

Here $\tau = M^2/S$, with M being the boson mass; Q^2 is a dynamical scale associated with the K-factor; $N_{W/Z}$ is the appropriate weak vertex factor squared and the quark distributions have been evolved to M^2. In evaluating this formula to predict $\sigma_{W/Z}$ there are two sources of uncertainty:

1) One needs to know the parton densities for all x, evolved to M^2 - which is a large scale. Fortunately, at the CERN collider the dominant values of x_i in (12) are $x_1 \simeq x_2 \simeq \sqrt{\tau} \simeq 0.15$. For this values, valence-valence collisions dominate and the relevant densities are quite well known. Going up in energy, as will happen with Tevatron, is less favorable from this point of view.

2) The scale Q^2 is not really determined until $O(\alpha_s^2)$ terms are computed. The natural choice for Q^2 would be $Q^2 \simeq M^2$, although one could envisage $Q^2 \simeq \langle P_\perp^2 \rangle_{W/Z}$, which is much less than M^2. At any rate, assuming that $Q^2 \simeq M^2$, one expects considerably smaller K-factors than in Drell-Yan processes, where the pair invariant mass is much smaller than the W/Z mass. Typically /11/ $K \simeq 1.3 - 1.4$ here.

The results of the recent calculation of Altarelli, Ellis and Martinelli /11/ are presented in Table I. The lower error in the table is due to effects of variation in the parton density, keeping $K(Q^2=M^2)$ fixed. The larger error takes $Q^2 = \langle P_\perp^2 \rangle_{W/Z}$, thereby increasing the value of K.

Table I Predictions for W and Z production at the CERN collider

$\sqrt{s}$ GeV	σ_W (M_W = 83 GeV) nb	σ_Z (M_Z = 94 GeV) nb
540	$4.2^{+1.3}_{-0.6}$	$1.3^{+0.4}_{-0.2}$
630	$5.3^{+1.6}_{-0.9}$	$1.6^{+0.5}_{-0.3}$

Using $B(W \to e\nu) = 0.089$; $B(Z \to e^+e^-) = 0.032$ one can compare these results with the values of $\sigma . B$ measured at the collider /7/. This is done graphically in Fig. 4, and one sees that the agreement is fine.

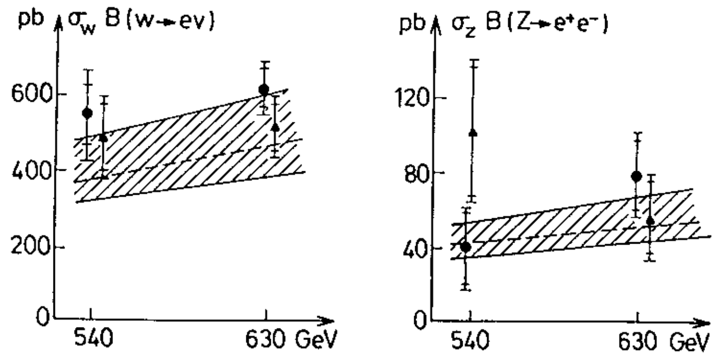

Fig. 4: Comparison of UA_1 ◆ and UA_2 ▲ data on (σ.B) with the QCD calculations of Ref. 11.

More exclusive quantities, than the total production cross section, can also be computed. An example of this is the W $p_\perp$ distribution. The calculation of this distribution is easy for $p_\perp \gg M_W$, but for $p_\perp \simeq M_W$ one must worry about all orders in α_s. The point is that one encounters large logarithms, of the type $\alpha_s^n(p_\perp^2) \ln^m M_W^2/p_\perp^2$, which cannot be ignored. Fortunately, one has been able to resum terms of this type /15/ and one obtains an eikonal-like representation for the differential $p_\perp$ distribution /11/

$$\frac{d\sigma}{dp_\perp^2} = \int \frac{d^2b}{4\pi}\, e^{-i\vec{p}_\perp \cdot \vec{b}}\, \sigma_o(1+A)\, e^{S(b)} + Y(p_\perp) \quad (13)$$

Here σ_o is the lowest order cross section and A contains the multiplicative $O(\alpha_s)$ corrections. The non singular $O(\alpha_s)$ corrections are in $Y(p_\perp)$, while the dangerous terms which have been resumed are in $S(b)$ which is now known to $O(\alpha_s^2)$ /16/. Although $S(b)$ affects the shape of the $p_\perp$ distribution, it does not contribute to the total cross section since $S(0) = 1$.

In Fig. 5, I show the $p_\perp$ distribution for the W boson determined by the UA_2 collaboration /7/, compared to the QCD prediction computed from Eq. (13) by Altarelli, Ellis, Greco and Martinelli /11/. The fit is obviously excellent. Note that the 1983 spectacular UA_2 events A, B, C with more data now appear to be consistent with theory. Very similar fits have been shown by the UA_1 collaboration /17/.

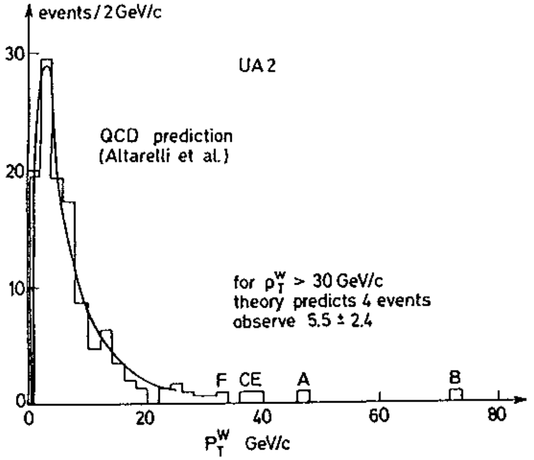

Fig. 5: Transverse momentum distribution of the W boson compared to the QCD calculation of Ref. 11

3. QCD NEWS

In this section I want to discuss some additional QCD tests for which new or more refined data has become available this year. I begin by looking again at some results coming out of the CERN collider.

3.1 Jets at the Collider

The two jet cross section at the collider arises from a combination of many subprocesses: $q\bar{q} \to q\bar{q}$; $qg \to qg$; $gg \to q\bar{q}$; etc. The $\hat{t}$ channel gluon exchange which enters in most of these subprocesses gives a typical Rutherford angular dependence:

$$\frac{d\sigma}{d\cos\theta} \sim \frac{\alpha_s^2}{\hat{s}(1-\cos\theta)^2} \qquad (14)$$

This characteristic dependence, which was already apparent in the early running of the collider /18/, can be clearly seen in Fig. 6, taken from a recent UA_1 publication /19/. In fact, the inclusion of small scale breaking effects, both in $\alpha_s(\hat{t})$ as well as in the evolution of the structure functions, seems to improve the fit. For the di jet mass range $\hat{s} = m_{2j} = 150 - 250$ GeV, initial state gluons are quite important. One finds /19/ that the qq to qg to gg subprocesses are in the

ratio of 36 : 52 : 12. Further, the gluonic contributions are dominant at small x values /20/.

Gluons are not only important in the initial state in the collider. There is clear evidence now for their "presence" in the final state, as three-jet events are clearly distinguishable in the data /21/. Both the UA_1 and the UA_2 collaborations have attempted to extract from the ratio of 3-jet to 2-jet processes a value for α_s. This is not an easy task and at present the inferred results should be considered only as semiquantitative.

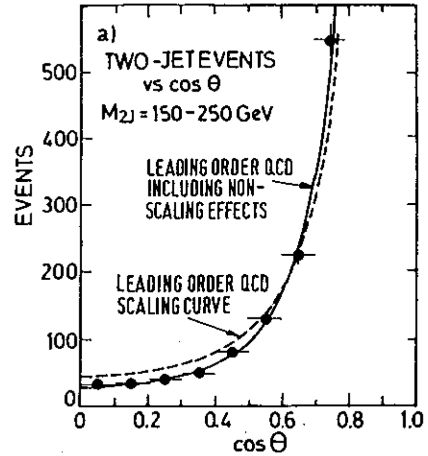

Fig. 6: Rutherford behaviour of the 2-jet cross section

If c_{3j} and c_{2j} are the fraction of 3-jet and 2-jet events one expects, then the cross section ratio is just

$$\frac{\sigma_{3j}}{\sigma_{2j}} = \frac{c_{3j}\,\alpha_s^3}{c_{2j}\,\alpha_s^2} \qquad (15)$$

Thus, if the jet fractions are equal and if the 3-jet and 2-jet processes depended on the same scale, the cross section ratio would directly measure α_s. Unfortunately the fractions c_{3j} and c_{2j} depend crucially on the subprocess and on the experimental cuts one is imposing, so that a careful analysis is needed before one can extract α_s from the cross section ratio. Furthermore, the typical scale appropriate for a 3-jet process is not the same as that for a 2-jet process, so that it is also not possible to simply cancel the α_s factors in (15). Finally, it should be commented that although c_{3j} and c_{2j} are calculable in perturbative QCD, their full calculation is not completed. Virtual processes of higher order need yet to be included, to properly determine the appropriate scale to evaluate numerator and denominator in (15).

To give a feeling of some of the ambiguities one encounters in the present analysis, I show in Fig. 7 a plot of σ_{3j}/σ_{2j} from the UA_1

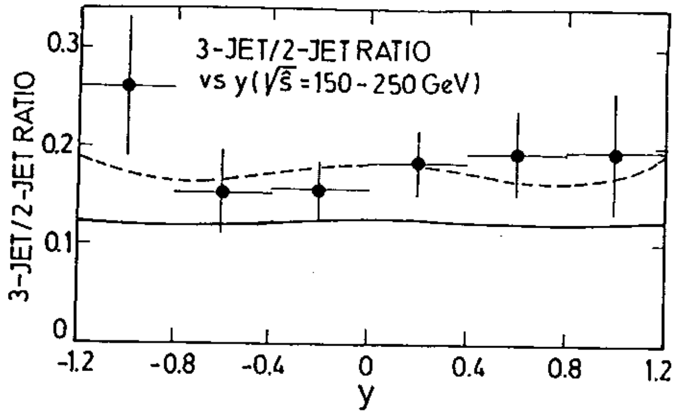

Fig. 7: Ratio of 3-jet to 2-jet cross section from Ref. 19

collaboration /19/, along with two theoretical fits. The solid line corresponds to a QCD fit in which the 2-jet and 3-jet scales are taken to be the same, while for the dotted line one has $\langle q^2\rangle_{3j} \simeq 4/9 \langle q^2\rangle_{2j}$. Clearly the latter curve gives a better fit to the data, for the choice of α_s taken. However, by increasing α_s the solid line could also be brought in agreement with the data. Hence, roughly speaking, the difference between the solid and dashed line gives one an idea on the uncertainty in α_s. This is borne out by the value obtained by the collaboration /19/, after their analysis

$$\alpha_s(4000\ \text{GeV}^2) = 0.16 \pm 0.02 \pm 0.03 \qquad (16)$$

where the first error is statistical and the second is systematic. A similar result was obtained by the UA_2 collaboration, although their value for α_s /22/ is somewhat larger, taking the K factors the same:

$$\alpha_s \frac{K_3}{K_2} = 0.23 \pm 0.01 \pm 0.04 \qquad (17)$$

The discrepancy, in my opinion, reflects nothing more than the uncertainties in the analysis.

3.2 Measurement of $F_L(x)$

One of the nicests tests of QCD which became available this year concerns the longitudinal structure function $F_L(x)$ in deep inelastic

scattering. This structure function is a combination of F_2 and F_1,

$$F_L = F_2 - 2x F_1 \qquad (18)$$

and is a measure of the spin of the constituents of the proton. In the parton model, where only spin 1/2 quarks contribute to the scattering, F_L vanishes. This is the famous Callan Gross relation /23/. However, in QCD, where also gluons enter, $F_L(x)$ acquires a non zero value. In fact, one predicts /24/

$$F_L(x;q^2) = \frac{\alpha_s(q^2)}{2\pi} x^2 \int_x^1 \frac{d\xi}{\xi^2} \{\frac{8}{3} F_2(\xi;q^2) + 16(1-\frac{x}{\xi})\xi\, G(\xi;q^2)\} \qquad (19)$$

where G is the gluon density function in the proton.

Unfortunately F_L is very difficult to extract experimentally and up to recently the available data did not allow for a significant test of (19). This situation has changed this year thanks to the new, high statistics, CDHSW data. Having more than half a million ν and $\bar{\nu}$ events they can extract F_L by studying the y distribution of

$$\frac{d\sigma^{\bar{\nu}}}{dx\,dy} - (1-y)^2 \frac{d\sigma^{\nu}}{dx\,dy}$$

Their results, presented by Feltesse /25/ are shown in Fig. 8, along with a QCD fit. The quality of the data and the theoretical agreement are impressive.

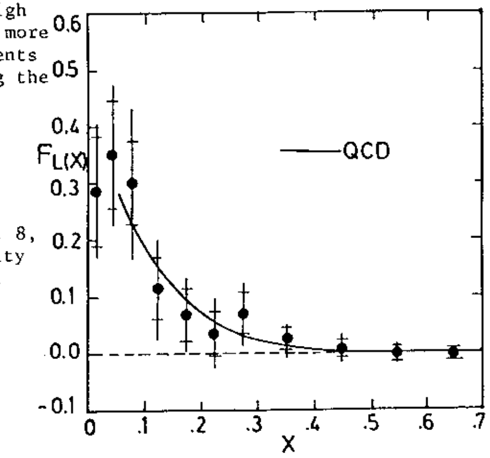

Fig. 8: CDHSW data on $F_L(x)$, from Ref. 25

3.3 Hard Photons and QCD

Another topic in which considerable experimental and theoretical progress has been achieved this year concerns direct photon production in hadronic reactions. Aurenche, Baier, Fontannaz and Schiff /26/ have completed a computation of higher order corrections for this process (see Fig. 9 for the relevant graphs). Interestingly enough, they find that, by optimizing the scales at which $\alpha_s(q^2)$ and the structure functions are evaluated, their results are very stable and lead to K factors very near to unity. Aurenche et al. /26/ have compared their calculations with a host of different prompt photon data, over a wide energy range from $\sqrt{s}$ = 19 GeV to $\sqrt{s}$ = 630 GeV. For $p_\perp \geq 4$ GeV, they find

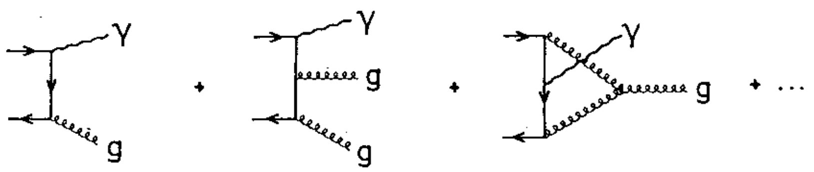

Fig. 9: Processes leading to hard γ production

very nice fits to the data, using for structure functions those of Duke and Owens, SetI with $\Lambda_{\overline{MS}}$ = 200 MeV /27/. I show some of these fits in Figs. 10.

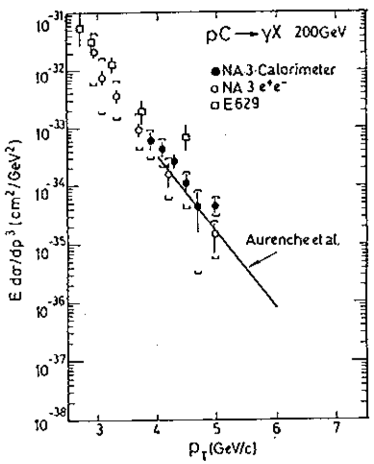

Fig. 10a: pC → γX at $\sqrt{s}$ = 19.4 GeV, from Ref. 28

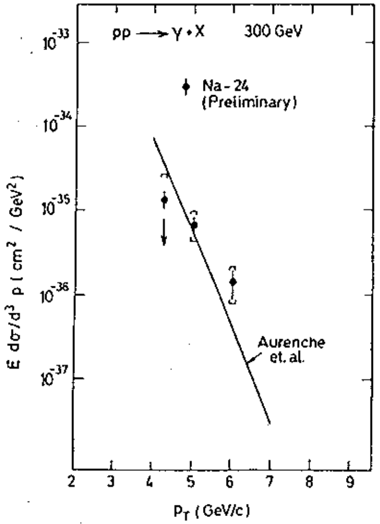

Fig. 10b: pp → γX at $\sqrt{s}$ = 23.8 GeV, from Ref. 29

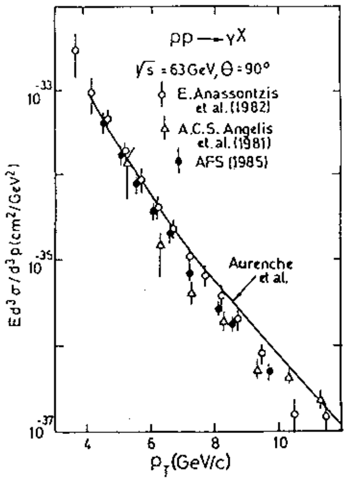

Fig. 10c: pp → γX at $\sqrt{s}$ = 63 GeV, from Ref. 30

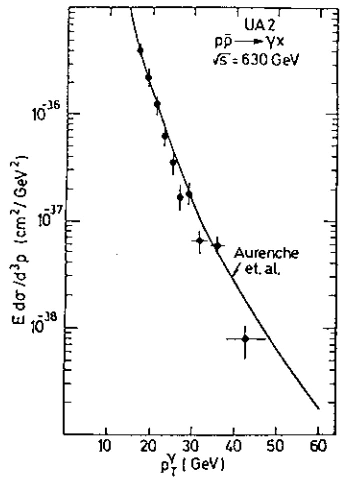

Fig. 10d: $p\bar{p} \to \gamma X$ at $\sqrt{s}$ = 630 GeV, from Ref. 31

Aurenche, Baier, Douiri, Fontannaz and Schiff /32/ have also computed the $O(\alpha_s)$ corrections to Compton scattering $\gamma q \to \gamma q$. Their

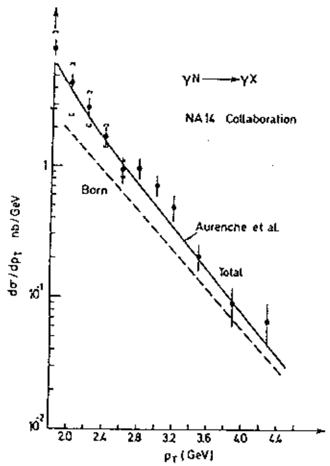

Fig. 11: Deep inelastic Compton scattering $\gamma N \to \gamma X$ data from Ref. 33 at E_γ = 100 GeV

results are in good agreement with the new NA 14 data /33/. Indeed, as shown in Fig. 11, the $O(\alpha_s)$ corrections improve the fit.

4. HIGGS BOSONS

The Higgs boson although an integral part of the standard model, is its most ephemeral entity. It is there to preserve renormalizability, although it could be obviated if some sort of dynamical symmetry breaking of SU(2)xU(1) obtained. Naturally, its theoretical and experimental investigation is of fundamental concern.

4.1 An Experimental Bound on m_H?

The mass of the Higgs boson, m_H, is not predicted by the standard model. Furthermore, since the coupling of the Higgs to fermions is proportional to the fermion's mass, Higgs bosons are difficult to produce. As a result, obtaining any kind of bound on Higgs bosons is a difficult business. Perhaps the most promising way to look for, relatively light, Higgs bosons is through the Wilczek mechanism /34/, in which a heavy quarkonia decays into a Higgs boson plus a photon. Since this rate is proportional to the mass squared of the heavy quark, it is clear that Υ decays are the most reasonable hunting ground, at present, for Higgs bosons.

The ratio between Υ decays into a Higgs boson and a photon and its decay into μ pairs is given by /34/:

$$\frac{\Gamma(\Upsilon\to H\gamma)}{\Gamma(\Upsilon\to\mu^+\mu^-)} = \frac{G_F\, m_b^2}{\sqrt{2}\,\pi\alpha}\left(1-\frac{m_H^2}{m_\Upsilon^2}\right) \simeq \frac{M_\Upsilon^2/M_W^2}{8\sin^2\theta_W}\left(1-\frac{m_H^2}{M_\Upsilon^2}\right) \tag{20}$$

For m_H not too near the kinematic boundary, this ratio is of the order of 2 %, and hence is within experimental reach. Recent CUSB data, presented by J. Lee Franzini at the Autun Meeting /35/, appears to exclude, at the 90 % confidence limit, Higgs bosons lighter than about 4.5 GeV. This is seen in Fig. 12. However, the Wilczek formula Eq. (20) has large QCD corrections and the issue is unsettled.

It has been known for a long time that the QCD corrections to the rate for quarkonia to decay into lepton pairs are rather large /36/. It turns out that the QCD corrections to decays of quarkonia into a Higgs and a photon, calculated by Vysotosky /37/ are even larger. Taking into account of both corrections one has

$$\frac{\Gamma(\Upsilon\to H\gamma)}{\Gamma(\Upsilon\to\mu^+\mu^-)} = R_{\text{Wilczek}}\,\frac{\left\{1-(40/3\pi)\,\alpha_s\,F(m_H^2/M_\Upsilon^2)\right\}}{\left\{1-(16/3\pi)\,\alpha_s\right\}} \tag{21}$$

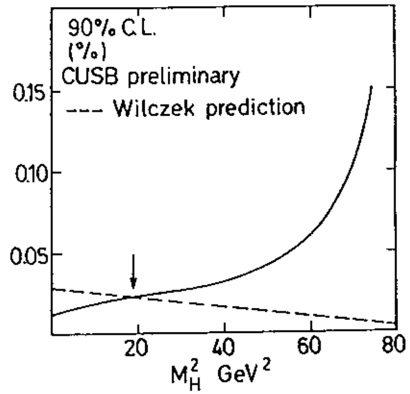

Fig. 12: Comparison of the CUSB data, Ref. 35, with the Wilczek prediction

The function $F(m_H^2/M_\Upsilon^2)$ is explicitly given in Vysotosky's paper /37/. However, for $m_H \lesssim 0.8\, M_\Upsilon$, $F \simeq 1$. Because the corrections in Eq. (21) are so large, one probably cannot trust them. But even proceeding naivly, expanding in α_s, one sees that these corrections vitiate the CUSB bound. Imagining that α_s is small in Eq. (21), one has for the ratio, for light m_H:

$$R = R_{\text{Wilczek}}\left\{1-\frac{8\,\alpha_s(M_\Upsilon^2)}{\pi}\right\} \tag{22}$$

If $\alpha_s(M_\Upsilon^2) = 0.15$ then the square bracket above is 0.62. For $\alpha_s(M_\Upsilon^2) = 0.20$, the Wilczek prediction is reduced by 50 %. Given that the data in Fig. 12 is barely below the original Wilczek prediction, it is clear that including QCD corrections removes the bound on m_H altogether. Of course, if one could experimentally go well below the Wilczek prediction in Υ decays, then even taking into account of the QCD corrections, one could rule out sufficiently light Higgs bosons.

4.2 Toponium as a Higgs Detector

If the top quark mass is really in the range $m_t = 40 \pm 10$ GeV, suggested by the UA_1 collaboration /38/, then toponium will give a very strong bound on the Higgs mass. Recalling that the Wilczek rate grows with the square of the quarkonia mass (cf. Eq. (20)), one sees that for toponium near 80 GeV the rate for decay into $H\gamma$ is of the order of the $\mu^+\mu^-$ rate. Even large QCD corrections will not affect this qualitative fact and one should be able to rule out (or discover!) Higgs bosons with masses within 5 - 10 GeV of the kinematic limit.

This matter has been studied carefully recently, in occasion of the LEP Jamboree /39/. First of all, provided the top quark has a mass less than half the Z^0 mass, its discovery at LEP should be straightforward. (Because of the much larger energy spread, the SLC is not in such a favorable condition.) Z^0 decays should determine the top mass within one GeV. Then at LEP, doing a scan of the relevant 2 GeV region, in 80 MeV steps, one should be able to determine the existence of toponium, by means of topological cuts, in about 2 weeks /39/. Having found toponium, then a Hγ signal is relatively easy to detect, even with rather small luminosity. This is shown in Fig. 13, for the case of a 70 GeV toponium, where the integrated luminosity needed for a 3σ Hγ signal is plotted versus the Higgs mass. Since at LEP an integrated luminosity of around 10 pb^{-1} per month is expected, it is clear that, if toponium is at 70 GeV, then the discovery of a 60 GeV Higgs boson will require only about one month running.

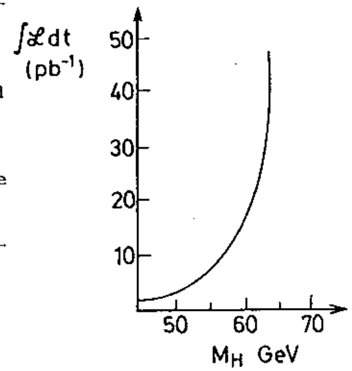

Fig. 13: Integrated luminosity needed for Higgs boson detection, if toponium has a mass of 70 GeV

4.3 Theoretical Guesstimates on m_H

Although, as I have already mentioned, the Higgs mass is not predictable in the standard model, predicting m_H is a favorite theoretical pastime. The results obtained, in general, reflect the prejudices put in. Two examples will suffice. Beg Panagiotakopoulos and Sirlin /40/ by requiring stability of the Higgs self coupling λ in the perturbative renormalization group equations, find that this coupling should be bound by the U(1) coupling constant squared. This implies, immediately, a bound on the Higgs mass:

$$m_H \leq 2\sqrt{2}\cot\theta_W M_W < 130 \text{ GeV} \tag{23}$$

Kubo, Sibold and Zimmermann /41/, on the other hand, solve the renormalization group equations under the assumptions that all couplings are driven by just one coupling. This reduction of couplings means that, effectively, all couplings are driven by the largest gauge coupling. Neglecting the SU(2)xU(1) couplings altogether and all Yukawa couplings, except that of the t quark, they find that either

$$m_H = 61 \text{ GeV}; \quad m_t = 81 \text{ GeV} \tag{24a}$$

or

$$m_H > 40 \text{ GeV}; \quad m_t < 81 \text{ GeV} \tag{24b}$$

if this coupling reduction holds.

Less dependent on particular dynamical assumptions is the study of the Higgs sector on the lattice. This has been undertaken by a number of authors recently /42/, most notably by Montvay and collaborators. What has been investigated is the pure SU(2) Higgs model in the presence of gauge fields and most particularly what happens when the Higgs self coupling λ becomes very large. What Montvay finds /42/ is that as $\lambda \to \infty$ the ratio of the Higgs mass to the W mass remains of O(1). Furthermore, his numerical results seem to show λ independence, over a wide range of λ, suggesting perhaps that λ may be an irrelevant variable. These results hold in a strong coupling region for the gauge coupling and need to be extrapolated to small g to make direct connection with physics. When this extrapolation is attempted Montvay finds that the m_H/M_W ratio grows. A preliminary value obtained this way is that $m_H \simeq 6\ M_W$ /42/.

Although these lattice investigations are really just beginning, they appear extremely interesting theoretically. Obviously, they still need considerable refinement. For instance, at the moment, no U(1) factor is included at all. Nevertheless, if these more refined lattice calculations continued to converge on such a large value for the Higgs mass, this would put the Higgs boson out of reach experimentally until the advent of the SSC!

5. KM MATRIX, CP VIOLATION AND ALL THAT

As a last topic of discussion I want to consider CP violation in the standard model. In particular I want to examine critically whether the usual explanation of CP violation, through the appearance of a phase in the Kobayashi Maskawa mixing matrix, is still tenable, or whether finally one is forced to go beyond the standard model. My answer (unfortunately?) will be that everything is still compatible with the standard model, although the model is being challenged.

5.1 Quark Mixings

The mass matrices for the quarks and leptons are beyond prediction in the standard model and so are the mixings among the quarks. However, the model does predict that the mixing matrix (KM matrix) is unitary. Thus, even though very little is known about the top quark, quite a lot is known about the mixing matrix elements V_{tb}, V_{ts} and V_{td}. In particular, the comparatively recent discovery of a very long B lifetime coupled to stringent bounds on the ratio of $b \to u$ to $b \to c$ transitions have provided important information for the structure of the KM matrix, establishing that $V_{us} >> V_{cb} >> V_{ub}$. Wolfenstein /43/ has given a handy approximate parametrization of the KM matrix, which takes these new facts into account and is easy to remember:

$$V_{KM} = \begin{vmatrix} 1-\frac{1}{2}\lambda_W^2 & \lambda_W & \lambda_W^3 A(\rho - i\eta) \\ -\lambda_W & 1-\frac{1}{2}\lambda_W^2 & \lambda_W^2 A \\ \lambda_3^3 A(1-\rho-i\eta) & -\lambda_W^2 A & 1 \end{vmatrix} \tag{25}$$

where

$$\lambda_W \simeq \sin\theta_c \simeq 0.23; \quad A \simeq 1; \quad \rho^2 + \eta^2 \leq 0.25 \tag{26}$$

5.2 CP Violation - ε Parameter

The smallness of V_{cb} and V_{ub}, along with a possible relatively light top quark /38/, have brought the standard model explanation for the CP parameter ε into question . (In fact, originally, Glashow Ginsparg and Wise /44/ used ε and the smallness of V_{cb}, V_{ub} to get a lower bound on m_t.) Recall that ε, in the standard model, is given by the imaginary part of the box graph shown in Fig. 14, in which all the charge 2/3 quarks enter. For the imaginary part, the t quark contribution is crucial and this is surpressed because V_{td} and V_{ts} are small (c.f. Eq. (25)).

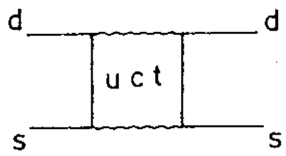

Fig. 14: Box graph whose imaginary part contributes to ε

Instead of expressing ε in terms of the KM matrix elements it is perhaps more useful to detail its dependence directly on other measured quantities /45/. Let me denote the B lifetime by

$$\tau_B = \beta \times 10^{-12} \text{ sec} \tag{27a}$$

and the ratio of the $b \to u$ to $b \to c$ leptonic decays by

$$R = \frac{\Gamma(b \to u e \nu)}{\Gamma(b \to c e \nu)} \tag{27b}$$

Experimentally the world average of all existing experiments, computed by D. Haidt /46/, gives for β:

$$\beta = 1.0 \pm 0.19 \tag{28a}$$

J. Lee Franzini /35/ quotes a 90 % confidence limit for R of

$$R < 0.03 \tag{28b}$$

while Thorndyke /47/ has a more conservative limit:

$$R < 0.04 \tag{28c}$$

The dominant dependence of ε on these two parameters and on the top quark mass scales as

$$\varepsilon \sim \frac{1}{\beta}\sqrt{R}\, m_t^2 \ . \tag{29}$$

Hence a longer B lifetime, a tighter bound on the $b \to u$ to $b \to c$ ratio and a smaller value for m_t all conspire to make ε smaller. Clearly, at some point, the standard model explanation for ε would then cease to be tenable. Even now, the situation is somewhat fluid, since a determination of ε requires besides the above experimental information also a theoretical calculation for the matrix element of $(\bar{d}\gamma_\mu(1-\gamma_5)s)^2$ between K and $\bar{K}$ states. This calculation is also quite uncertain.

It has been customary to denote by B the ratio of the above matrix element to that computed via vacuum insertion. Various approaches have been followed to compute B, ranging from bag model calculations to lattice calculations, with results ranging from about - 0.4 to +2.5 for B. Two values of B appear special: B = 1, which is just the vacuum insertion approximation and B = 1/3, which is a result that follows from current algebra /48/ (see below). At any rate, the predicted value for ε is proportional to B and thus, whether one considers the standard model to be in trouble or not depends in part on the value for B assumed.

Buras /49/ has displayed this interrelation, between experimental and theoretical input for ε, in a nice way. He computes, as a function of β, R and m_t, what is the minimum value of B necessary to fit in the standard model the experimental value for ε. As an illustration, I display in Table II, some of his results /49/ for a selected range of β, R and m_t.

Table II Minimum values of B needed to fit ε in the standard model

m_t (GeV) \ β	0.7	1.0	1.3	
25	0.45	0.75	1.08	R = 0.03
40	0.23	0.41	0.63	R = 0.03
55	0.14	0.24	0.40	R = 0.03
25	0.86	1.42	2.01	R = 0.01
40	0.47	0.82	1.22	R = 0.01
55	0.29	0.53	0.82	R = 0.01

Clearly if $\beta = 1$, $R = 0.03$ and $m_t = 40$ GeV a value of $B = 1/3$ would just be slightly unacceptable. However, if $R = 0.01$, with $\beta = 1$ and $m_t = 40$ GeV, really B must be near unity, for the standard model not to be in trouble. I would characterize the situation as tantalyzing, but not yet critical. Obiously, theoretically, it is important to pin down the value of B.

5.3 B Parameter Controversy

Given the above discussion, it is particularly important to examine the theoretical basis of predictions for the B parameter which yield a small value. Of particular importance, in this respect, is the current algebra prediction of Donoghue, Golowich and Holstein /48/ for B, $B \simeq 1/3$. Their result is rather easy to understand and appears to require very little theoretical input. What these authors realized is that, in the limit of good chiral $SU(3)_L \times SU(3)_R$, the 4-quark operator that enters in the K-$\bar{K}$ matrix element (and therefore in B) is related to that which describes the $K^+ \to \pi^+\pi^0$ decay. Both of these operators transform as $(27_L, 1_R)$. Using an effective chiral Lagrangian the amplitudes for both of these processes are then fixed, save for an overall normalization constant. This constant can, however, be determined from the experimental value for $A(K^+ \to \pi^+\pi^0)$ and it is this procedure which gives $B \simeq 1/3$.

Because the logic above is so simple, considerable credence was given to a small value for B. However, recently, Bijnens, Sonoda and Wise /50/ have brought into question the validity of the chiral limit for the evaluation of B. They have calculated $O(m_\pi^2 \ln m_K^2)$ corrections to the chiral limit of Donoghue et al. /47/ and found that, for the case of the $\Delta s = 2$ operators, these corrections are larger than the zero order result! Thus the chiral limit inference that $B \simeq 1/3$ may be flawed.

This situation has been made even more difficult to interpret by a very recent paper of Pich and de Rafael /51/. These authors have reexamined this issue from a completely different point of view - that of QCD sum rules - and find again that $B \simeq 1/3$! Let me briefly discuss their idea, so as to give a flavor of their calculations. What one needs to calculate is the K-$\bar{K}$ matrix element of the $\Delta s = 2$ operator $O = (\bar{d}\gamma_\mu(1-\gamma_5)s)^2$. What one does, instead, is to calculate the matrix element of an appropriate effective chiral operator, with the same transformation properties as O. To actually get a number, of course, one must know the proportionality constant relating O to this effective operator. To be specific, de Rafael and Pich /51/ replace O by

$$O \to \mathcal{L}^{eff}_{chiral} = G_{\Delta s=2}\left([f_\pi^2\, U\partial_\mu\, U^+]_{23}\right)^2 \qquad (30)$$

where U is an SU(3) matrix of Goldstone boson fields, but $G_{\Delta s=2}$ is unknown. To get a value of B they need to fix $G_{\Delta s=2}$.

Donoghue et al. /47/ fixed $G_{\Delta s=2}$ by relating it to the amplitude for $K^+ \to \pi^+\pi^0$, determined experimentally. Pich and de Rafael /51/, on the other hand, do this by comparing the behaviour of two different two point functions in the fashion of QCD sum rules. What they consider are the two point functions of the operators O and of $\mathcal{L}_{eff}$, $\Delta(q^2)$ and $\Delta_{eff}(q^2)$, and compare integrals over their spectral functions:

$$\int dq^2 \, \mathrm{Im}\, \Delta(q^2) = \int dq^2 \, \mathrm{Im}\, \Delta_{eff}(q^2) \qquad (31)$$

For the integral on the left-hand side above they use the answer obtained from the QCD short distance behaviour. For the right-hand side, they make use of resonance saturation in the chiral model. Matching the results gives a value for $G_{\Delta s=2}$ and therefore B. What Pich and de Rafael find in this way is

$$B = 0.33 \pm 0.09 \qquad (32)$$

It appears rather amazing to me that this calculation should agree so well with the simple chiral result of Donoghue et al./48/. Thus, it could well be that the agreement is fortuitous. However, Guberina, Pich and de Rafael /52/, have made an analogous, but in detail different, calculation for the operator that enters in $K^+ \to \pi^+\pi^0$ and computed the amplituds $A(K^+ \to \pi^+\pi^0)$. Their result $A_{GPR} = 1.8 \times 10^7$ sec^{-1} is in excellent agreement with the experimental value ($A_{exp} = 1.7 \times 10^7$ sec^{-1}). This would argue that their methods are reliable and so that one should trust also their result (32).

5.4 CP Violation - ε'/ε

The situation with ε'/ε is even more uncertain than that with ε. Now, even though some of the dependence on the specific size of the KM matrix elements is milder, there is a new, uncertain, hadronic matrix element to estimate /53/:

$$\langle \pi\pi | O_{Penguin} | K \rangle \sim B' \qquad (33)$$

Also B' estimates can vary by more than a factor of 2, so that the theoretical uncertainty in $\varepsilon'/\varepsilon \sim B'/B$ can be really quite large. Two typical ranges, for $m_t = 40$ GeV, $\beta = 1$, which appear in the recent literature, are

$$10^{-3} \leq \varepsilon'/\varepsilon \leq 15 \times 10^{-3} \qquad \text{/Ref. 54/} \qquad (34a)$$

$$2 \times 10^{-3} \leq \varepsilon'/\varepsilon \leq 8 \times 10^{-3} \qquad \text{/Ref. 49/} \qquad (34b)$$

Thus the wonderful experimental limits /55/

$$\varepsilon'/\varepsilon = (-4.6 \pm 5.3 \pm 2.4) \times 10^{-3} \qquad \text{/Chicago-Saclay/}$$

$$\varepsilon'/\varepsilon = (1.7 \pm 8.2) \times 10^{-3} \qquad \text{/BNL-Yale/}$$

badly need a more accurate theory prediction, to really push the standard model.

6. CONCLUDING REMARKS

I conclude with pretty much the observation I made at the beginning: SU(3)xSU(2)xU(1) works remarkably well! Some hope does exist for finding some discrepancy in this beautiful edifice. For one thing, the Higgs sector is essentially unknown. Here toponium and the lattice calculations may begin to shed some light. It is possible that the CP violation parameter ε may need some new physics for its explanation. But that would require $R < 0.01$, m_t being really light and that the theoretical ambiguities in B were finally resolved! Also, it is clear that a measurmend of ε'/ε at the 10^{-3} level would seriously impact the model, expecially if some of the theoretical ambiguities in B',B were under better control.

My personal conclusion is that the real question should be shifted from: does the standard model work? to, why is the standard model a description of nature? The list of unanswered questions in this latter case is rich and deep. A partial sampling includes:

Why SU(3)xSU(2)xU(1)?
Why chiral fermions?
Why is $(\sqrt{2}\,G_F)^{-1/2} \sim 10^3\,\Lambda_{QCD}$?
Why do the fermions replicate?
What fixes the fermion masses and mixing?

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