WEAK INTERACTIONS AND GAUGE THEORIES

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I shall report on the status of the electroweak gauge theory, also known as quantum asthenodynamics (QAD: astheno = weak). The major result is that the standard[1,2] WS-GIM model describes the data well, although one should still look for signs of further complexity and better tests of its gauge theory aspect. A second important result is that the measured values of the three basic coupling constants of present energy physics, g_s, g and $\sqrt{5/3}\,g'$ of $SU(3)_c \times SU(2)_2 \times U(1)$, are compatible with the idea that these interactions are unified at high energies as will be discussed by Wilczek. Most of my talk will be devoted to open questions. We know little about the Higgs sector of the theory. We know something about the fermion sector, but we don't understand it. We know little about the origin of CP violation. Some of the answers to these questions may be hidden in the very high energy regime of the theory, but we must exploit as best we can our low energy laboratories. One important endeavor is to measure accurately the parameters available to us by, for example, studies of b and hopefully soon t-quark decay, and also by taking a harder look at CP violation in the neutral kaon system. There is a body of weak interaction data that has been available to us for many years but never fully understood, namely non-leptonic weak decay amplitudes. New data is becoming available with the measured decays of the Ω^- and charmed particles, and we can ask whether the technology of gauge theories allows a better understanding of these processes. This field has seen renewed activity, much of it centered around penguinology. Since the penguin diagram will appear at various places in my talk, I define it at the outset in Fig. 1.
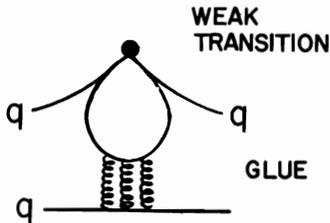

WEAK TRANSITION

GLUE

Fig. 1 Penguin Diagram

1. The Status of QAD

Accumulating experimental data has been steadily confirming the standard model of the electroweak interactions. Fits to the data with the exception of atomic physics results, are shown in Fig. 2. The striking success of the model is best illustrated by the essentially unconstrained fit[3] in the top half of the figure. A vector-axial vector current-current interaction was assumed; the factorization hypothesis is necessary for combining electron-deuterium data with neutrino data, but was also independently checked. The only other assumption was equality of strange and down quark couplings, which plays a minor role in the analysis. The couplings are normalized so that the neutrino coupling ν_ρ and the electron axial vector coupling g_A^e are 1 and 1/2, respectively in the standard model. The remaining parameters relevant to the u,d and e couplings are displayed so that the plotted data points measure the weak angle in the standard model. In other words, they should be equal and lie at a value between zero

and one if the standard model is correct. Considering that they could take a priori any values, the result is impressive. The lower half of Fig. 2 shows fits[3,4] to the $SU(2)_L \times U(1)$ model with no assumption on the Higgs structure (open circles) and assuming weak iso-doublet Higgs multiplets (closed circles) which constrains the overall normalization of the neutral current Fermi coupling. The last two points are from the Peking-Fermilab νe elastic scattering data[5] and the electron-deuterium data[6] presented at this conference. The shaded bar represents the prediction[7,8] for $\sin^2\theta_w$ if the electro-weak and strong interactions become unified, assuming that no new physics enters in such a way as to modify the coupling constant evolution prior to the unification energy, and that the strong interaction coupling strength is characterized by $\Lambda \simeq 0.5$ GeV as extracted[9] from deep inelastic lepton-nucleon scattering data.

The conclusion to be drawn from the data points of Fig. 2 is that neutral current transitions at present energies are described by the simple four-fermion coupling

$$H_{NC} = \frac{G_F}{\sqrt{2}} \, (I_3^w - \sin^2\theta_w \, Q)^2 \qquad (1.1)$$

as written down[1] by Weinberg and Salam. In addition, although the results are less precise, data from both neutrino-induced dilepton production[10] and the study of charmed particle decays[11] strongly suggest that the charged charm-changing current is the one written down[2] by Glashow, Iliopoulos and Maiani: A V-A current satisfying a $\Delta C = \Delta Q$ rule and dominated by $\Delta C = \Delta S$ transitions. If these two statements are correct they imply that the weak interaction couplings are those of a renormalizable theory[12], and moreover they have the simplest possible form compatible with renormalizability. Are there other possibilities? An obvious modification compatible with the data is

$$SU(2)_L \times U(1) \times G \qquad (1.2)$$

where G is an arbitrary group under which ordinary particles are invariant. This hypothesis is not as content-free as it may seem because intermediate boson mixing could lead to unexpected structure in their propagators[13]. The data[14] from PETRA so far shows no hint of such effects. A more popular alternative is

$$SU(2)_L \times SU(2)_R \times U(1). \qquad (1.3)$$

In this case the data is telling us that the effective Fermi constant for $SU(2)_R$ is negligibly weak at present energies, which means

$$M_{W_R} \gg M_{W_L} \qquad (1.4)$$

Since most theorists believe anyway that there are super heavy vector bosons in addition to the "moderately" heavy $W^\pm$ and Z of $SU(2)_L \times U(1)$, this possibility cannot be dismissed. However, the lesson of the data[6,15] is that if electroweak interactions are described by a gauge theory, the gauge group relevant[16] at present energies is $SU(2)_L \times U(1)$.

The question remains: is it a gauge theory[17]? All that has been confirmed so far is the structure (1.1)
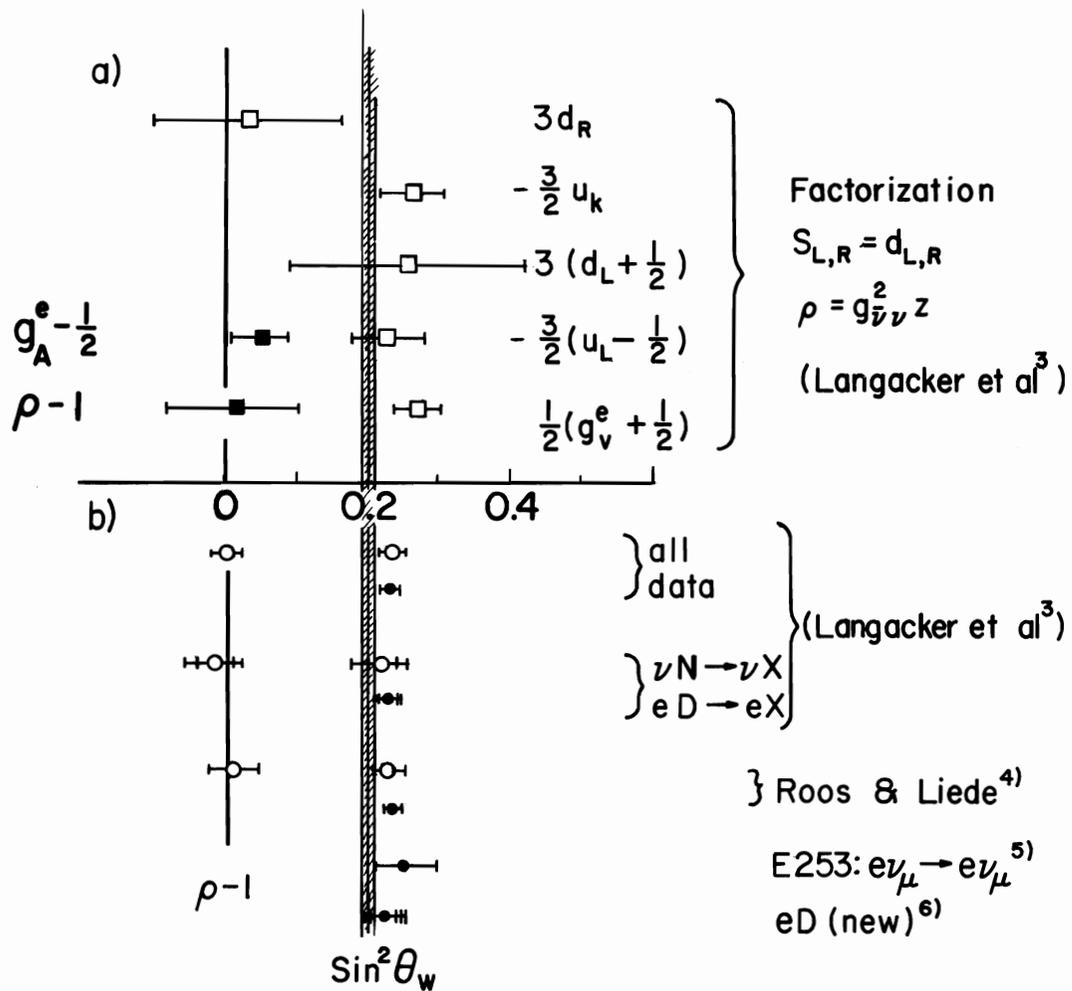
Fig. 2 Fits to Neutral Current Data

  a) Unconstrained fit to neutral current couplings compared with standard model predictions: zero (full squares) and $\sin^2\theta_w$ (open squares).

  b) Fits to $SU(2)_L \times U(1)$ with $\rho = G_F^{NC}/G_F^{CC}$ unconstrained (open circles) and with $\rho = 1$ (full circles). Multiple error bars separate statistical (systematic) and theoretical errors. The shaded bar allows unification with strong interactions if $\Lambda \simeq 0.5$ GeV.

of the effective Fermi interaction. More direct confirmation of the gauge theory aspect will come with the anticipated observation of the $W^{\pm}$ and Z in $p\bar{p}$ colliders and possibly of their propagator effects[18] in PETRA and PEP. A crucial test of the gauge theory is the measurement of the vector boson self coupling vertex of Fig. 3 which probably must await a LEP 100. The primary motivation for believing in a gauge theory is that higher order effects are calculable. QAD is a theory as respectable as QED, and its real test is the confrontation of higher order effects with experiment. The need for such tests has been particularly emphasized by Veltman, and considerable effort has been devoted[19-22] to the calculation of observable deviations from the lowest order theory. The difficulty is that higher order corrections are dominated by soft photon effects which are simply proportional to the Born approximation result and are uninteresting from the standpoint of testing QAD. One has to look for
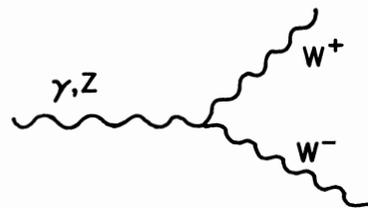

Fig. 3 Tri-Vector Coupling of QAD

special cases where the Born term is suppressed, so that "hard photon" corrections, which are inextricably

tied together with W and Z exchange effects, may become observable. In $e^+e^- \to \mu^+\mu^-$ the lowest order γ and Z exchange diagrams interfere destructively in the forward direction at energies just below the Z mass and in the backward direction just above. In these regions higher order "hard photon" or QAD effects become appreciable[19]:

$$\frac{\Delta\sigma_{QAD}}{\sigma_{Born}} \simeq \begin{array}{ll} 0.3, & \theta_{cm} = 20^o, \quad E_{cm} = (73\text{-}80) \text{ GeV} \\ 0.4, & \theta_{cm} = 160^o, \quad E_{cm} = (110\text{-}114) \text{ GeV} \end{array} \tag{1.5}$$

for $\sin^2\theta_w = 0.22$ which gives $M_Z \simeq 90$ GeV. Another interesting process[20] is $e^+e^- \to W^+W^-$ where radiative corrections are sensitive to the Higgs mass. Increasing the Higgs mass from 10 to 1000 GeV increases $\Delta\sigma_{QAD}$ from about 6% to 10% at $E_{cm} = 200$ GeV. A higher order QAD effect which might be obserable at lower energies is an induced isoscalar axial vector component in the hadronic nuclear current which is absent in lowest order in the standard model. In this case the effect is dominated by gluon exchange effects[23], Fig. 4, but if one can isolate the amplitude governed by the electron vector coupling, both the Born term and the gluon exchange effects are suppressed because of the suppression of the electron vector coupling

$$g_V^e = 2\sin^2\theta_w - 1/2 \tag{1.6}$$

for values of $\sin^2\theta_w$ close to 0.25. In this case the QAD induced effects are found to be[22]

$$(g_A^{I=0})_{QAD} / (g_V^{I=1})_{Born} \simeq .10 \tag{1.7}$$

In addition, precision measurements of the W and Z masses together with a determination of θ_w from the current structure alone would determine the deviation from the standard model prediction

$$m_Z/m_W\cos\theta_w = 1 \tag{1.8}$$

which is also sensitive to radiative corrections.
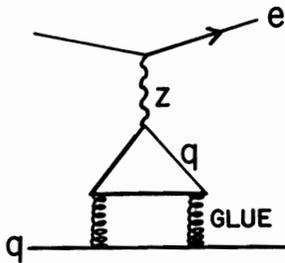

Fig. 4 Gluon exchange mechanism for induced isoscalar axial vector hadronic neutral current.

2. The Scalar Meson Spectrum

The experimental confirmation of the overall normalization of the neutral current coupling (1) tells us that the scalars which contribute to the W and Z masses are doublets under weak isospin. However, their number is arbitrary. There are philosophical arguments for and against a proliferation of Higgs scalars. The assumption of a single Higgs doublet immediately insures a) natural suppression of flavor changing neutral current transitions and b) conservation of electric charge: the third axis in weak isospin space is defined so that

$$\langle H\rangle \equiv \begin{pmatrix} \langle\phi^+\rangle \\ \langle\phi^o\rangle \end{pmatrix} \equiv \begin{pmatrix} o \\ v \end{pmatrix} \tag{2.1}$$

for a single multiplet, but in the general case vacuum expectations values need not be aligned in weak isospin space. Both of these arguments can be evaded for multi-Higgs systems by imposing discreet symmetries and/or a suitable choice of Higgs potential. Arguments for more than one Higgs include a) the possibility of imposing discreet symmetries which restrict the fermion mass matrix and b) the possibility of "soft CP violation" as a solution to the "strong CP violation" problem. These issues will be discussed in more detail below. If more than one Higgs doublet is introduced, there are necessarily physical charged scalar particles, and the best way to decide the issue is to look for their production in e^+e^- annihilation or in semi-weak decays[24] like $t \to b + H^+$, which could complete favorably[25] with $t\bar{t} \to 3g$ in toponium decay.

What can be said about the Higgs mass? If there is only one Higgs doublet, the potential takes the form

$$V(\phi) = -\mu^2|\phi|^2 + \lambda|\phi|^4 + A|\phi|^4 \ell n (|\phi|^2/m^2) + \ldots\ldots \tag{2.2}$$

where the last term is the leading contribution from radiative corrections[26]: $A = O(g^4, \lambda^2)$. If $\lambda \gg A$ radiative corrections are negligible and the potential has a minimum at

$$|\phi|^2 = \mu^2/2\lambda \equiv \frac{1}{2\sqrt{2}G_F} \tag{2.3}$$

and the mass of the physical Higgs particle is

$$m_H^2 = 4\mu^2 = 2\sqrt{2}\lambda/G_F = \frac{4\sqrt{\lambda}}{g} m_W, \tag{2.4}$$

which depends on the scalar self coupling strength. If λ gets too small, however, the radiative corrections dominate, and stability of the vacuum, $V(\langle\phi\rangle) < V(0)$, requires[27]

$$m_H \gtrsim 7 \text{ GeV}. \tag{2.5}$$

A smaller value is possible if the "observed" vacuum is only a local minimum; then vacuum tunneling is possible[28], and the lifetime of the universe implies[29]

$$m_H \gtrsim 260 \text{ MeV} \tag{2.6}$$

If λ gets too large perturbation theory becomes inapplicable; requiring a convergent perturbative expansion gives bounds[30]

$$m_H \lesssim 1 \text{ TeV} \tag{2.7}$$

A "natural value" (e.g. from an underlying supersymmetry) relating 4-point scalar and vector couplings might be $\lambda = O(g^2)$ in which case one would get

$$m_H \sim m_W. \tag{2.8}$$

A recent guess for the Higgs mass is related to the "hierarchy problem" in grand unified theories. The symmetry breaking from the unified gauge group G_V down to the "observed" exact symmetry goes in two stages

$$G_V \xrightarrow{V} SU(3)_c \times SU(2)_L \times U(1) \xrightarrow{v} SU(3)_c \times U(1)_{e.m.}, \tag{2.9}$$

characterized by the ratio

$$v/V \sim 10^{-12} \qquad (2.10)$$

where V and v are the vacuum expectation values of the scalar fields which are responsible for the two stages of symmetry breaking. Since we cannot let λ be arbitrarily large we necessarily have

$$\mu^2 <<< v^2 \qquad (2.11)$$

and an attractive assumption is that some underlying principle dictates that μ should vanish. Then there is no symmetry breakdown unless $\lambda \lesssim O(g^2)$ in which case radiative corrections dominate. In fact, the parameter λ appearing in Eq. (2.2) is not really a constant but has an implicit dependence on m; m is an artificial normalization parameter which cannot effect physics:

$$\partial V/\partial m = 0 \qquad (2.12)$$

Together with the known form of the radiative corrections, Eq. (2.12) determines the m-dependence of λ, defined for example by ($\mu^2 \equiv 0$):

$$V(\phi) \equiv \lambda(m)|\phi|^4 + A|\phi|^4 \ln(\phi^2/m^2) + \ldots \qquad (2.13)$$

so that $\lambda(m)$ is the effective scalar coupling, at $|\phi| = m$. Since $A > 0$ we get $\partial\lambda(m)/\partial m > 0$ so that if $\lambda = O(g^2)$, say at a mass scale m = V characteristic of the first symmetry breakdown, it will vanish at some smaller value, m = v, as illustrated in Fig. 1. Since the dependence is logarithmic and $A = O(g^2)$ the ratio
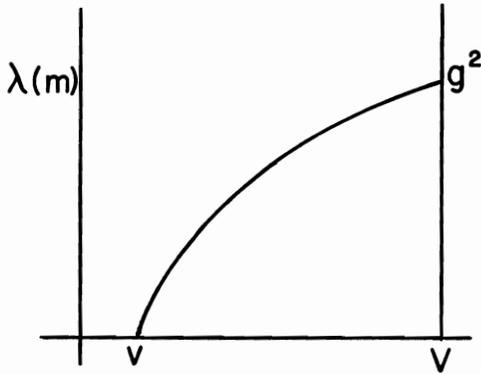
$$v/V = e^{-O(1/g^2)} \qquad (2.14)$$

Fig. 5 Evolution of the scalar self coupling constant

can be quite small. Explicit calculation[31] shows that if the conditions $\mu \equiv 0$ and $\lambda = O(g^2)$ are met at the grand unification scale, $V = O(10^{15}$ GeV), the coupling λ will vanish at the mass scale where $SU(2)_L \times U(1)$ is known to break down. In specific models studied[32], the condition $\mu \equiv 0$ is unnatural; radiative corrections of the type in Fig. 6 lead to a natural mass scale $\mu^2 \sim g^4V^2$. Nevertheless one can speculate that $\mu^2 = 0$ on the grounds that this is more plausible (i.e. there is a further, unknown symmetry) than a non-zero value which is accidentally tiny with respect to the unification mass scale. Then at the $SU(2)_L \times U(1)$ level there are no unknown parameters and the mass of the Higgs particle is determined[33,34]:

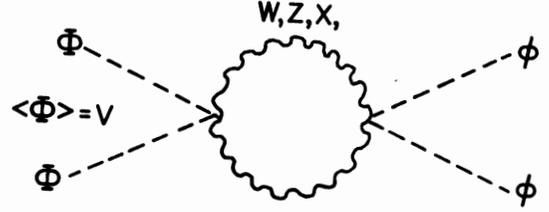$$m_H = (\frac{\partial^2 V}{\partial H^2})_{\phi = <\phi>} \simeq 10.4 \text{ GeV} \qquad (2.15)$$

Fig. 6 Diagram contributing large mass term to Higgs doublet in GUTS.

for $\sin^2\theta_w \simeq 0.20$ if there are no heavy fermions, $m_f \sim m_W$. This possibility is interesting in that it may be accessible to present experiments and also because its proximity in mass to the $b\bar{b}$ bound state system can lead to interesting mixing effects[35,34] with a scalar $b\bar{b}$ state as illustrated[34] in Fig. 7. The best source of such a particle would be in toponium (T) radiative decay[36] should it be found at PETRA or PEP. Figure 8 summarizes theoretical guesses at the Higgs mass in terms of a probability distribution[25]; the drop-off below $m_H = 15$ MeV is the only constraint provided so far by experiment.
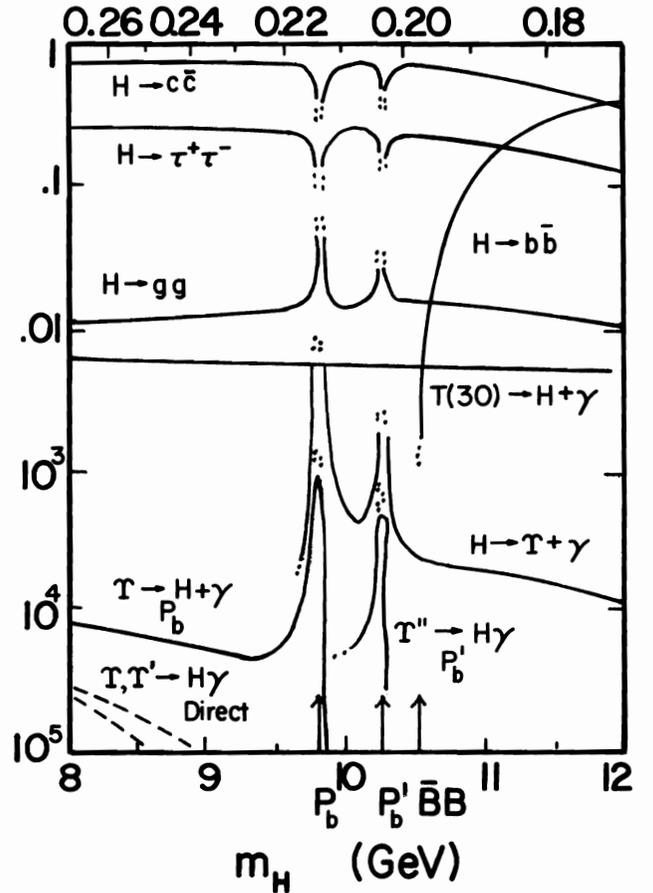

Fig. 7 Profile of a 10 GeV Higgs[34]

3. The Fermion Spectrum

Up to now there is very little understanding as to why there is a repetition of fermions with identical quantum numbers nor of the observed pattern of masses and mixing angles. I shall briefly review some of the arguments in the literature which attempt to limit the number of fermion generations or to relate mass matrix parameters.

3.1 How many generations? There are general arguments based on the requirement that the theory be self-consistent and calculable in perturbation theory. Fermion doublets with large mass splittings can give large corrections, via the diagram of Fig. 9, to the mass relation (1.8), which follows from the assumption of weak isodoublet Higgs scalars, and insures the equality of the effective Fermi couplings for neutral and charged current interactions. The fits of Fig. 1 give

$$1 - G_F^{NC}/G_F^{CC} \lesssim 0.04 \qquad (3.1)$$

allowing two standard deviations, whereas the diagram of Fig. 9 induces a contribution[37]

$$1 - G_F^{NC}/G_F^{CC} > 0.05 \qquad (3.2)$$

if there is a lepton doublet

$$M_{\nu_L} = 0, \; M_L \gtrsim 400 \text{ GeV} \qquad (3.3)$$

or a color triplet quark doublet

$$M_{b'} = 1/3 \, M_{t'}, \gtrsim 100 \text{ GeV} \qquad (3.3)$$

This analysis gives the only firm bounds based on data, but there are arguments based on the desired validity of perturbation theory, similar in the spirit to the argument bounding the Higgs mass from above. For example unitarity breaks down[38] in the Born approximation in the fermion Higgs sector for fermion masses greater than a TeV because the fermion-Higgs coupling constant grows with M_f in the minimal (one Higgs) model. A better restriction[39] comes from examining the fermion loop contribution to the effective Higgs potential. The constant A in Eq. (2.2) is proportional to

$$A \propto M_H^4 + 6M_W^4 + 3M_Z^4 - 4M_f^2 \qquad (3.4)$$

where M_H is the Higgs mass in the tree approximation (2.4), and the bounds (2.5) and (2.6) were obtained neglecting possible heavy fermion contributions. For $A < 0$, the potential becomes negative for large values of $|\phi|$, so the vacuum would be unstable. From the unitarity bound (2.7), the condition $A > 0$ gives

$$M_f \lesssim 800 \text{ GeV} \qquad (3.5)$$

for a heavy lepton. Tree unitarity only restricts $\lambda \leq 16\pi$; the more stringent condition $\lambda < 1$ would give

$$M_f \lesssim 135 \text{ GeV} \qquad (3.6)$$

For a color triplet of quarks the bounds (3.5) and (3.6) are reduced by about 30%, and for n heavy fermions they are reduced by a factor $n^{-1/4}$.

Further restrictions on the number of generations are suggested by the study of grand unified theories (GUTS) of strong and electroweak interactions. In SU(5) the simplest assumption on the Higgs sector requires[40]

$$M_b = M_\tau \qquad (3.7)$$

in the symmetry limit. Symmetry breaking corrections[41,8,42] to (3.7) reproduce the "observed" mass of about 5 GeV if there are only three, or possibly four generations. This is obviously model dependent; with a different Higgs multiplet one can impose

$$M_b = 1/3 \, M_\tau \qquad (3.9)$$

instead of (3.7); then five or six genenerations[43] are required to get the desired value after renormalization. Another argument[44] is based on the stability of the proton. The number of fermion generations does not affect significantly the rate at which the strong and electroweak coupling constants come together, but it does affect their common value at the unification energy. With more than 8 generations the coupling is sufficiently large that the proton decay is expected to exceed the experimental limit[45] of about 10^{-30}/yr. Still in the context of SU(5), one can demand[46] that the running coupling constants for both the Yukawa $H f \bar{f}$ couplings and the scalar self-coupling remain sufficiently small (and $\lambda(m) > 0$) that the lowest order evolution equations remain a good approximation at
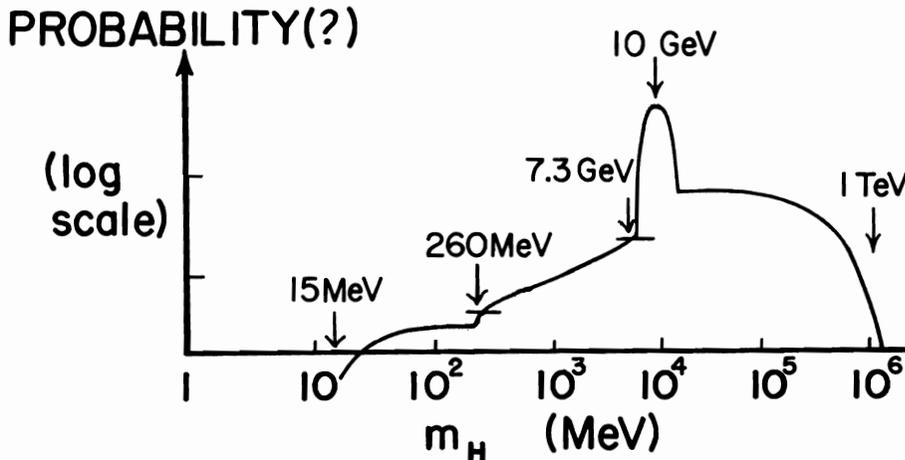

Fig. 8 Theoretically inspired probability distribution for the Higgs mass[25]
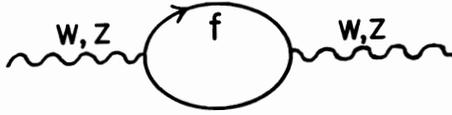
Fig. 9 Potentially large radiative
corrections to fermi coupling strength.

energy scales up to the grand unification mass; this
restriction gives a limit of about 200 GeV for both
the Higgs and the top quark masses.

While such arguments do not provide rigorous
limits, they strongly suggest that fermion generations
should not keep duplicating themselves with the
"canonical" mass ratio of about three for each succes-
sive generation. For example, if we believe that
radiative corrections govern the spontaneous breakdown
of $SU(2)_L \times U(1)$ as discussed above, $\mu \simeq \lambda(v) \simeq 0$, then
$m_H = 0$ in Eq. (3.4) and the stability condition $A > 0$
gives

$$M_t < 80 \text{ GeV}; \qquad (3.10)$$

and together with the experimental limit $M_t \gtrsim 15$ GeV
implies that a new generation would have to satisfy

$$15 \text{ GeV} \leq M_{t',b',L} \leq 70 \text{ GeV} \qquad (3.11)$$

so that the empirical factor of three rule would not
allow it.

A final argument from astrophysics[47] is based on
the helium abundance of the universe and limits the
number of quasi-massless, weakly coupled neutrino
helicity states (not including their anti-particles) to
three or possibly four, as long as there is not a large
ν-$\bar{\nu}$ asymmetry. This provides a meaningful limit in the
context of SU(5), for example, which predicts only
massless neutrinos and a ν-$\bar{\nu}$ asymmetry of the same
order as the observed baryon number asymmetry of
$(10^{-10} - 10^{-8})$ per photon.

3.2 Can we calculate fermion mass matrices?
Unified theories of weak and strong interactions group
quarks with leptons into larger multiplets and for
simple Yukawa couplings their masses may be related in
the symmetry limit. As mentioned above, the simplest
Higgs structure in SU(5) leads to the relation

$$M_{q(-1/3)} = M_{\ell(-1)} \qquad (3.12)$$

for each generation. If there are exactly three
generations symmetry breaking effects modify (3.12) to
give[41,8,42]

$$M_b \simeq (4.8 - 5.6) \text{ GeV} \qquad (3.13a)$$

$$M_s \simeq (0.4 - 0.5) \text{ GeV} \qquad (3.13b)$$

$$(M_d/M_s)_{bare} = M_e/M_\mu \qquad (3.13c)$$

Eq. (3.13a) is generally regarded as a success, the
success of (3.13b) is controversial, and there is a
definite problem with (3.13c) since the "bare" or
"current" quark ratio is believed to be 1/20 based on
PCAC analyses[48]. On the other hand, no one believes
that SU(5) is a complete theory and any small external
source of mass could contribute to the tiny e and d
masses so as to change (3.13c) without significantly
modifying the other results.

What about the other elements in the fermion mass
matrices? SU(5) (nor, to my knowledge, any proposed
phenomenologically acceptable embedding of SU(5) in a
larger gauge group) has nothing to say about the mass
ratios of charge +2/3 to charge –1/3 quarks, mass
ratios between generations, nor generalized (complex)
Cabibbo angles. These parameters are not determined
by grand unification alone.

It has been observed by several authors[49,50,51]
that a particular form of the mass matrix, e.g.

$$M = \begin{pmatrix} o & a & o \\ a & o & b \\ o & b & c \end{pmatrix} \qquad (3.14)$$

in the three generation case, yields after diagonal-
ization a well known[52] phenomenological relation for
the Cabibbo angle:

$$\theta_c^2 \simeq (M_d/M_s)_{current}. \qquad (3.15)$$

Using (3.14) one gets for the top quark mass

$$M_t \simeq M_b (M_u M_c / M_d M_s)^{1/2}. \qquad (3.16)$$

There is considerable uncertainty as to what values
should be used for the masses on the right hand side
of (3.16), but a judicious choice gives

$$M_t \simeq 3M_b$$

a result which was also the popular guess based on
numerology. It is generally speculated that a form
like (3.14) might arise from some discreet or possibly
continuous symmetry which could be imposed on the
Yukawa couplings of Higgs isodoublets to fermions.
However, if the electroweak gauge theory is $SU(2)_L \times$
$U(1)$, there is a general theorem[53] which states that
there is no discreet or continuous symmetry which
allows a non-trivial prediction ($\theta_c \neq 0, \pi$) for the
Cabibbo angle if flavor changing neutral current
couplings are naturally suppressed. The difficulty in
imposing the latter criterion comes from the appearance
of flavor changing neutral Higgs couplings[54], inducing
for example $K_L \to \mu^+\mu^-$ via the diagram of Fig. 10.
These effects can be made arbitrarily small, however,
by letting the relevant Higgs masses get arbitrarily
large; in a multi Higgs model this is possible[55] with-
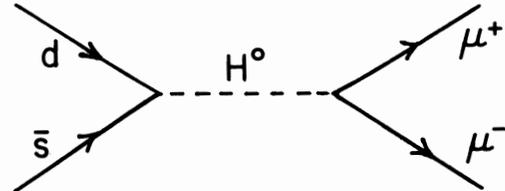out encountering the strong coupling disease discussed
in the previous section.

Fig. 10 Mechanism for $K_L \to \mu\mu$
in a multi-Higgs model.

Finally, one can impose discreet symmetries on the
Yukawa couplings of grand unified theories. Since
this necessarily implies a Higgs sector more complex
than the simplest possible choice, one can arrange
couplings so that they not only produce "new" relations
like (3.15) and (3.16) but also "improve" the "old"
ones, Eqs. (3.13). In particular[56] couplings can be
arranged so that (3.13a) remains unchanged, the right-
hand side of (3.13b) is divided by 3, and the right-
hand side of (3.13c) is multiplied by 9, modifications
which give better accord with theoretical prejudice.

To illustrate the general theoretical ignorance concerning the fermion mass matrix, Fig. 11 shows a histogram of predictions[50],[51],[57] which have appeared in the literature for the mass region of $t\bar{t}$ onia and naked top threshold. Results[14] from the four high energy data points at PETRA exclude the region below about 30 GeV, but it is still possible that some data points lie between the toponium 1^- ground state and bare top threshold. To give an idea of the mass scale which remains to be scanned, the lowest upper bound on the top mass comes from the assumption that $SU(2)_L$ x $U(1)$ breaking arises from radiative corrections ($\lambda \simeq 0$) to the Higgs potential, giving eq. (3.10).

4. CP Violation

The origin of CP violation is not yet understood and the mystery has only deepened with the discovery[58] that QCD contains a potential source of strong CP violation. I shall briefly review theoretical and phenomenological aspects of the problem.

4.1 Why is CP violation weak? Non-perturbative phenomena contribute an effective term to the QCD Lagrangian:

$$L_{QCD} \ni \theta \frac{\alpha_s}{8\pi} \tilde{F}^i_{\mu\nu} F^i_{\mu\nu} \qquad (4.1)$$

which is odd under both parity and CP, where $F^i_{\mu\nu}$ is the gluon field strength tensor, $\tilde{F}$ its dual and θ is an a priori arbitrary parameter. Present limits[59] on the neutron electric dipole moment restrict the parameter θ to be very small[60],

$$\theta \lesssim \text{a few x } 10^{-9}, \qquad (4.2)$$

and the puzzle[61] is: what makes it small? There are several alternative viewpoints.

a) θ is identically zero. This can be assured by imposing an extra chiral (i.e. helicity-dependent) global symmetry[62] on fermion couplings in addition to the local gauge symmetry. Within the context of the standard model this requires either the existence of the axion[63] or one massless quark, the most plausible possibility being $M_u = 0$. Both these possibilities are disfavored phenomenologically[64] unless the Higgs sector is contrived so that the axion mass gets large with its couplings to ordinary particles remaining samll[65].

b) θ is small and finite (i.e. calculable). The problem is that even if one sets $\theta = 0$ in the QCD Lagrangian, the weak source of CP violation, which necessary exists to account for the observed CP violation in the neutral kaon system, will in general generate a non-zero θ via radiative corrections which are infinite unless the source[66] of CP violation is "soft", which means that the original CP violating term in the Lagrangian has dimension ≤ 2 as does a scalar mass term:

$$L_{CP} = \phi_i^* \mu_{ij}^2 \phi_j \qquad (4.3)$$

A possible objection to this possibility is that soft CP violation disappears at high energies, thus invalidating recent conjectures[67] that the combined features of CP violation and baryon number violation in unified theories allow an understanding of the observed baryon number asymmetry of the universe; this mechanism requires CP violating forces to play a role at superhigh temperatures after the big bang when baryon number violating forces were important. However, it

has recently been shown[68] that a certain class of soft CP violating models allows a choice of couplings such that CP violation does remain important at high temperatures. These models have a definite prediction: there must be several Higgs multiplets with mass $M_H \simeq M_W$.

c) θ is small and infinite. The point is that just as coupling constants run, the paramater θ also runs. Any unspecified parameter in the theory has infinite radiative corrections which are absorbed by defining the parameter at some renormalization point specified in terms of external momenta. If the theory is renormalizable the value at any other point is finite and calculable in terms of the first. In the standard Kobayaski-Maskawa(KM) model[69] of CP violation, if θ is specified at Λ, then for a momentum relevant to the neutron one finds[70]

$$\theta(M_n) - \theta(\Lambda) \sim (\frac{M_q}{M_W})^4 (\frac{\alpha}{\pi})^2 + (\frac{M_q}{M_W})^{12} (\frac{\alpha}{\pi})^7 \ln (\Lambda/M_n) = 0(10^{-16})$$
$$\qquad (4.4)$$

as long as $\Lambda < e^{10^{18}}$ GeV. Few theorists believe that the presently known interactions describe all of particle physics. It may well be that the symmetry principle which sets $\theta = 0$ is spontaneously broken at some superhigh energy, plausibly the Planck mass. If we set $\theta(M_p) = 0$, there is no strong CP violation problem in the standard model which is in fact better off than some soft CP violation models which give $\theta = 0(10^{-7})$.

4.2 The Kobayashi-Maskawa model. In the standard model CP violation can appear in the Yukawa couplings of the $SU(2)_L$ x $U(1)$ symmetric Lagrangian; after symmetry breaking and diagonalization of the fermion mass matrix the CP violating term is transferred to the gauge couplings of fermions to the charged vector bosons and appears via complex generalized Cabibbo angles. In the six flavor model the Cabibbo matrix can be expressed in terms of four observable parameters:

$$U_c = \begin{pmatrix} c_1 & s_1 c_3 & s_1 s_3 \\ -s_1 c_2 & c_1 c_2 c_3 - s_2 s_3 e^{i\delta} & c_1 c_2 s_3 + s_2 c_3 e^{i\delta} \\ s_1 s_2 & -c_1 s_2 c_3 - c_2 s_3 e^{i\delta} & -c_1 s_2 s_3 + c_2 c_3 e^{i\delta} \end{pmatrix} \qquad (4.5)$$

where $c_i \equiv \cos\theta_i$, $s_i \equiv \sin\theta_i$. In the limit of small mixing angles (4.5) can be approximated by

$$U_c \simeq \begin{pmatrix} 1 & \theta_{us} & \theta_{ub} \\ \theta_{cd} & 1 & \theta_{cb} \\ \theta_{td} & \theta_{ts} & 1 \end{pmatrix} ; \theta_{us} \equiv \theta_c, \qquad (4.6)$$

where θ_c is the Cabibbo angle and the $\theta_{qq'}$ are related by unitarity constraints as implied by the form (4.5). Both intuition and the data suggest that quarks couple preferentially to "nearest neighbors" in mass:

$$|\phi_{ub}|, |\phi_{td}| \lesssim \theta_c^2$$
$$|\phi_{cb}| \approx |\phi_{ts}| \lesssim \theta_c \qquad (4.7)$$

For example, the experimental success of Cabibbo universality

$$(U_c)_{ud}^2 + (U_c)_{us}^2 \simeq 1 \qquad (4.8)$$

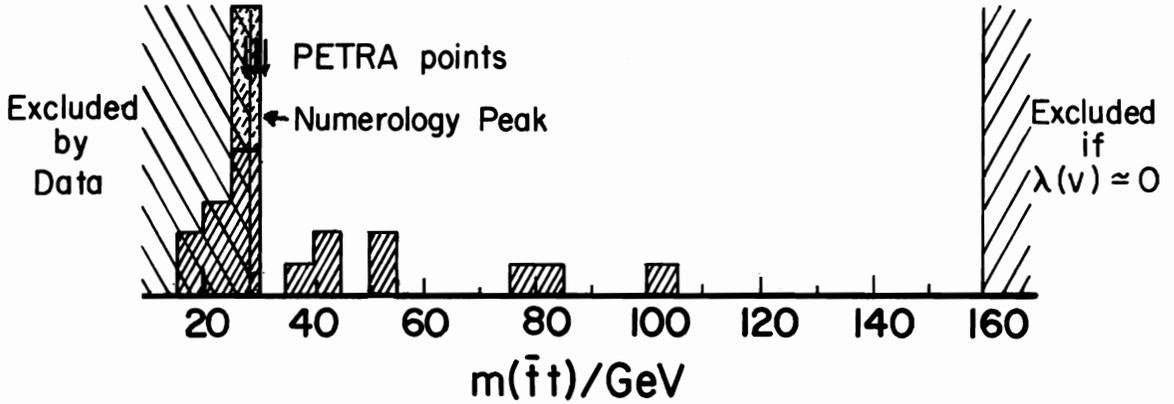limits[71] the allowed value of θ_{ub}. The most recent analysis[72] gives

Fig. 11 Predictions for Toponium Mass

$$\left|\theta_{ub}/\theta_{us}\right| = 0.28\ ^{+0.21}_{-0.28} \qquad (4.9)$$

The t quark couplings enter in low energy phenomenology only through virtual effects like the K_L-K_S mass difference, Fig. 12.
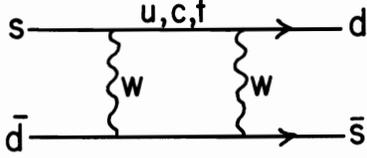

Fig. 12 Contributions to the K_L-K_S mass difference.

Since the original estimates[73] of the charmed quark mass based on Fig. 12 without the t contribution turned out to be in the right ball park, one doesn't expect[74] that the t-quark contribution can be too important. Several analyses[75] have recently been performed which attempt to extract the quantity $|\theta_{ts}\ \theta^*_{td}|$ as a function of the top quark mass. However, these are fraught with uncertainties in the matrix element of the effective four quark operator obtained from Fig. 12, the strong interaction corrections to the free quark diagram in the low momentum region of integration, and the value of the charmed quark mass which should be used in the Feynman integral. A new limit on the top quark couplings has recently been obtained[76] from the decay rate for $K_L \rightarrow \mu^+\mu^-$, Fig. 13. While this process was largely at the origin of the GIM mechanism, cancellation between Figs. 13a and 13b resulted in only a poor limit[73], $M_c \lesssim 9$ GeV, on the charmed quark mass and so it was forgotten. However, we already know[14] that $M_t > 9$ GeV, and precisely because the c contribution is unimportant the calculation is rather insensitive to uncertainties related to M_c and low momentum contributions, in addition to the fact that there is no matrix element uncertainty in the short distance approximation. Shrock and Voloshin find:

$$|\phi_{ts}\ \theta^*_{td}| \lesssim \theta_{us}\ \frac{57\ (\text{GeV})^2}{M_t^2} \lesssim 0.06 \simeq (0.38)^3 \qquad (4.10)$$

where the last inequality corresponds to $M_t > 15$ GeV.

4.3 Phenomenology. The only measured CP violating effect is in the neutral kaon system, and the present data are compatible with the "super weak" model[77] which means that CP violation is confined to K^0-$\bar{K}^0$ mixing. In the K-M model the K^0-$\bar{K}^0$ mixing diagram of Fig. 12 is complex if $\delta \neq 0, \pi$ in the Cabibbo matrix (4.5). Gilman and Wise[78] recently pointed out that if penguin diagrams are important in K-decay, CP violation in the $K \rightarrow 2\pi$ decay amplitude may give an observable deviation from the super weak model via the diagram of Fig. 14. There are three amplitudes relevant to the analysis of CP violation in $K \rightarrow 2\pi$:

$$A_m=A(K^0\leftrightarrow\bar{K}^0),\ \ A_o=A[K\rightarrow(2\pi)_{I=0}],\ \ A_2=A[K\rightarrow(2\pi)_{I=2}] \quad (4.11)$$

One can always choose a phase covention so that one of these amplitudes is real, and the standard choice is the Wu-Yang (WY) convention[79] which defines A_o as real:

$$(\text{Im}A_o)_{WY} \equiv 0 \qquad (4.12)$$

The parameter ε_m is related to the phase of A_m:

$$\varepsilon_m \equiv \frac{(\text{Im}\ A_m)_{WY}}{\Delta m} \qquad (4.13)$$

and ε' is related to the phase of A_2:

$$|\varepsilon'| \equiv \frac{(\text{Im}\ A_2)_{WY}}{\sqrt{2}\ A_o} \qquad (4.14)$$

In the super weak model, $\varepsilon' = 0$, and the CP violating parameter measured in K-decay is $\eta_{+-} = \eta_{oo} = \varepsilon \simeq (\varepsilon_m/\sqrt{2})\ e^{i\ \pi/4}$. In the KM model with the phase convention used in Eq. (4.5), CP violation occurs in the $\Delta I = 1/2$ amplitude A_o: $(\text{Im}\ A_2)_{KM} = 0$.
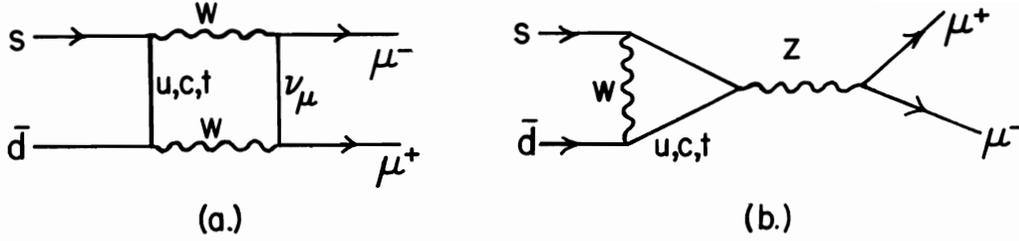
Redefining the amplitudes to match the Wu-Yang convention:
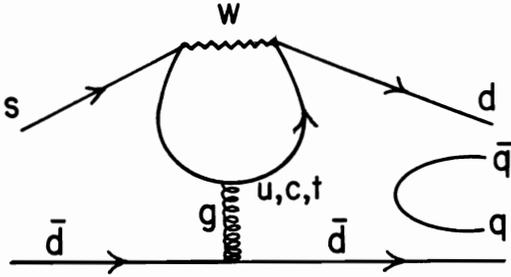
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Fig. 13   Contributions to $K_L \to \mu^+\mu^-$

Fig. 14   Diagram contributing to $\varepsilon' \neq 0$

$$(\text{Arg } A_2)_{WY}=-(\text{Arg } A_o)_{KM}, \ (\text{Arg } A_m)_{WY}=(\text{Arg } A_m+2\text{Arg } A_o)_{KM} \tag{4.16}$$

we get (for small CP violation)

$$\left|\frac{\varepsilon'}{\varepsilon}\right| \simeq \left|\frac{A_2}{A_o}\frac{\text{Im } A_o}{\text{Re } A_o}/(\frac{\text{Im } A_m}{\text{Re } A_m}+2\frac{\text{Im } A_o}{\text{Re } A_o})\right| = \begin{cases} 1/50\text{-}1/150 \ (\text{GW})^{80} \\ 1/250\text{-}1/500(\text{GP})^{81} \end{cases}$$

where the two groups have calculated the effective Fermi coupling of the penguin operator of Fig. 14 to all orders in QCD leading logs, and differ only in their evaluation of its importance in the $K \to 2\pi$ transition. The GW evaluation avoids an estimate of the operator matrix element but requires trusting the leading log approximation at low momentum transfer and needs an a priori estimate of the importance of Fig. 14 in the $K \to 2\pi$ decay rates. The GP estimate avoids the last two problems, using only the experimental decay rates and an evaluation of high momentum gluon corrections to Fig. 14, but requires the knowledge of the matrix element. My own prejudice is that the GP evaluation is more realistic. If the result is anywhere in the ranges given in Eq. (4.17) the next experiment[82] should be able to detect a deviation from super weak theory; the present data constrain[83] $|\varepsilon'/\varepsilon| \lesssim 1/50$.

Another observable source of CP violation may be the neutral B-meson system. B^o-$\bar{B}^o$ mixing may be appreciable because, as for kaons and unlike D's, the Cabibbo allowed transition $b \to t$ is kinematically inaccessible and the GIM mechanism, exact if $M_t = M_c = M_u$, is badly broken: $M_t \gg M_b \gg M_{c,u}$. Then same sign dileptons may not be infrequent above $B\bar{B}$ threshold in e^+e^- annihilation:

$$e^+e^- \to B^o\bar{B}^o + X \to \ell^-\ell^- + X, \tag{4.18}$$
$$\quad\quad\quad \hookrightarrow B^o$$

and a lepton charge asymmetry would signal CP violation[74,84].   A recent analysis within the K-M model gives[85]

$$\frac{|N^{++}-N^{--}|}{N^{+-}} \lesssim \begin{cases} 10^{-4}, \ B^o= (b\bar{d}) \\ 10^{-1}, \ B^o= (b\bar{s}) \end{cases} \tag{4.19}$$

assuming $M_t = (15\text{-}30)$ GeV, $\theta_{bu}/\theta_{su} = (0\text{-}0.5)$ and using further constraints on the θ_{ij} extracted[75] from ΔM_K. Generally mixing effects increase as the t mass is increased but CP violating effects decrease (because for large M_t the CP violation in the mass matrix dominates and the B^o-$\bar{B}^o$ system becomes effectively super weak). In addition the effects may be larger than (4.19) if the phase δ in (4.5) is larger than values suggested by the uncertainly fraught analyses of ΔM_K. It has also been pointed out[86] that diagrams analogous to Fig. 14 can interfere with the usual V-A four fermi effective operator to give CP violating effects in charged B-decays; then one could get hadron-anti-hadron symmetries in one-particle-inclusive measurements:

$$e^+e^- \to B^+B^- + X \to h + X \tag{4.20}$$

Clearly these experiments will be difficult. In particular the process (4.18) has an important background from cascade decays[87]

$$e^+e^- \to B^o\bar{B}^o + X \to \ell^-\ell^- + X \tag{4.21}$$
$$\quad\quad\quad \hookrightarrow \bar{D} + X$$

without $B\bar{B}$ mixing. However, I think it is important to look for any asymmetries in particle-antiparticle spectra above $B\bar{B}$ threshold, since this might be the first observed effect of CP violation outside the neutral kaon system.

5. Decay Dynamics

The parameters of the K-M matrix (4.5) could be further pinned down by a measurement of the B lifetime. Neglecting strong interaction effects gives[74,75]:

$$\tau_B^{-1} \simeq 5\Gamma_\mu \left(\frac{m_b}{m_\mu}\right)^5 [\theta_{bc}^2 \ F(m_c/m_b) + \theta_{bu}^2]$$
$$\begin{cases} \approx 10^{13} \ \text{sec}^{-1}; \ \theta_{bc} \simeq \theta_c, \ \theta_{bu} \simeq \theta_c^2 \\ \lesssim 10^{15} \ \text{sec}^{-1}; \ \theta_{bu}, \ \theta_{bc} < 1 \end{cases} \tag{5.1}$$

where the V-A phase space factor is

$$F(m_c/m_b) \simeq 0.3 - 0.5$$

for reasonable values of m_c and m_b. Published limits[88] give only

$$\tau_B^{-1} \gtrsim 10^9 \ \text{sec}^{-1} \qquad (5.2)$$

but better limits should soon be forthcoming from PETRA.

In order to take the estimate (5.1) seriously, we have to believe that the approximation of a freely decaying quark is a good one. It is, therefore, appropriate to ask how well the same approximation works for charmed particle decays, and more generally how well non-leptonic decay dynamics is understood.

5.1 Inclusive D-decay. In the free quark model, Fig. 10a, the inclusive-semi-leptonic branching ratios are predicted to be

$$B(e) = B(\mu) = 20\%. \qquad (5.3)$$

Hard gluon corrections to the non-leptonic weak vertex, Fig. 10b, modify this result[89,90]

$$B(e) = B(\mu) \simeq 10\%. \qquad (5.4)$$

The same model gives predictions for the total decay width and the inclusive semileptonic decay spectrum, but these are sensitive to the effects of hard gluon bremsstrahlung[91], Fig. 16, and uncertainties in the quark masses M_c and M_s, which largely cancel out in the branching ratio (5.4). One approach[92] is to fit the predicted lepton decay spectrum to the observed one in order to determine M_c and M_s. With hard gluon radiation effects included a good fit is obtained for

$$M_c \simeq 1.75 \ \text{GeV}, \ M_s \simeq 0.5 \ \text{GeV}. \qquad (5.5)$$

These values allow a prediction of the total lifetime[92]

$$\tau_D \simeq 5 \times 10^{13} \ \text{sec}.$$

While isotopic spin selection rules require

$$\Gamma_{D^o}(\ell) = \Gamma_{D^+}(\ell) \qquad (5.6)$$

the prediction

$$\Gamma_{D^o}(\text{all}) = \Gamma_{D^+}(\text{all}) \qquad (5.7)$$

is specific to the above model, and does not seem to be supported by data[93] presented at the conference[94] Corrections could arise from final state interactions among the final state quarks in Fig. 15, but the approximate scaling observed in Gargamelle neutrino events and e^+e^- annihilations at momentum transfers as low as those relevant to charm decay suggest that final state interactions should not drastically modify the estimate of inclusive decay rates. A different decay mechanism, which contributes only to D^o decay for Cabibbo favored modes is the $q\bar{q}$ "annihilation" process of Fig. 17a. Via a Fierz transformation it is equiva-
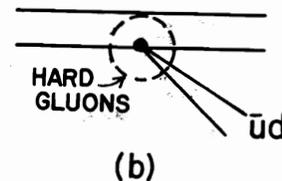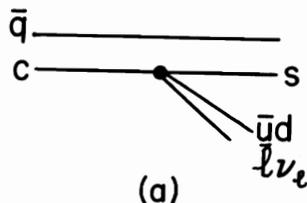
lent to the diagram of Fig. 17b, and a "parton model" estimate with hard gluon vertex corrections included gives:

$$\frac{\Gamma_{\text{annihilation}}}{\Gamma_{c \to su\bar{d}}} \simeq \frac{8\pi^2}{3} \frac{f_D^2 m_s^2}{m_c^4} \simeq 10^{-2} \left(\frac{f_D}{f_\pi}\right)^2 \qquad (5.8)$$

where the evaluation on the right hand uses the mass values (5.5) which some theorists would consider optimistic, and f_D, f_π are the D- and pion-axial current coupling constants: the ratio $(f_D/f_\pi)^2$ has been estimated[95] to be of order 10. The suppression of the annihilation mechanism is due to the same helicity conservation effect which suppresses the $\ell\nu$ modes of pseudoscalar decays. However, this suppression need not be operative if the final state quark pair is emitted together with a hard gluon via Fig. 18 because in this case[96] the final state ($s\bar{d}$) system can have J=1. Another possible mechanism for enhancement of D^o non-leptonic decays is resonance dominance[97]

$$D^o \xrightarrow{\text{weak}} (0^+, 0^-)_{S = -1} \xrightarrow{\text{strong}} \text{hadrons}. \qquad (5.9)$$

Since final states from D^+ decay are exotic, they would not be enhanced by this mechanism. However, the weak transition in (5.9) is again forbidden in the chiral symmetric quark model. In addition, data[93] presented here suggest that F^+ and D^+ lifetimes are similar, whereas the above mechanisms should also enhance F^+ decays - one reservation being that color factors are such that the mechanism analogous to Fig. 18 for F^+ decay requires emission of two gluons.

5.2 D decays: two body. The effective non-leptonic charm changing interaction in the GIM-KM model is (dropping Dirac matrices)

$$H_{\text{GIM}} \ni (\bar{c}s)(\bar{d}u) + \theta_{cd}(\bar{c}d)(\bar{d}u) + \theta_{us}(\bar{c}s)(\bar{s}u) + h.c.$$
$$+ \ldots.. \qquad (5.10)$$

The Cabibbo allowed $|\Delta S|=1$ piece has U=1, $|U_3|=1$, where U-spin is the SU(2) subgroup of flavor $SU(N_F)$ which mixes s and d. The "first Cabibbo forbidden" $\Delta S=0$ piece has $U_3=0$ and in the limit of the GIM four flavor model: $\theta_{cd} = -\theta_{us}$ has U=1. Since D^o is a U-spin singlet and $(K^\pm, \pi^\pm)$ form U-spin doublets one gets a simple SU(3) relation[98]

$$\Gamma(D^o \to K^+K^-) = \Gamma(D^o \to \pi^+\pi^-) = \tan^2\theta_c \ \Gamma(D^o \to K^-\pi^+) \qquad (5.11)$$

Present data[99] suggest that this relation may be rather strongly violated. While the experimental errors are still too large to confirm a discrepancy, theorists have offered various suggestions to explain a deviation from (5.11).

Fig. 15 Parton model for charm decay without (a) and with (b) hard gluon radiative corrections

a) SU(3) breaking effects generally tend to favor the KK Cabibbo suppressed mode over the $\pi\pi$ mode, in accordance with observation. In a simple quark model, the decays (5.11) occur via the diagram of Fig. 19 (similar diagrams describe reasonably well[89,100] $\Delta I=3/2$ transitions in strangeness changing decays) with matrix elements proportional to:

$$\Gamma(D \to PP') \propto f_{P'} (M_D^2 - M_P^2) < P \, |q\dagger c| \, D >$$
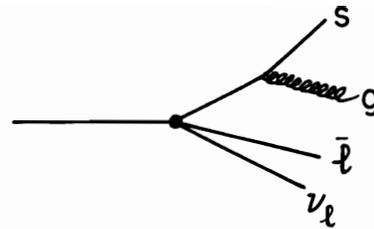
$$(5.12)$$

Fig. 16 Hard gluon bremsstrahlun mechanism which softens lepton spectrum in semi-leptonic D-decay
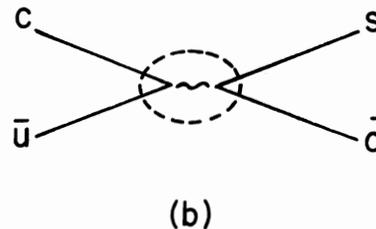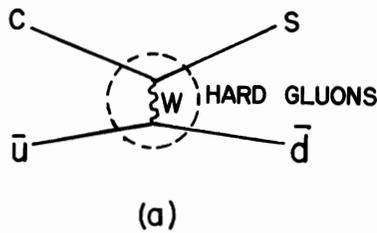

(a)

(b)

Fig. 17 "Annihilation" diagram (a) for $D^O \to$ hadrons and (b) its Fierz transform
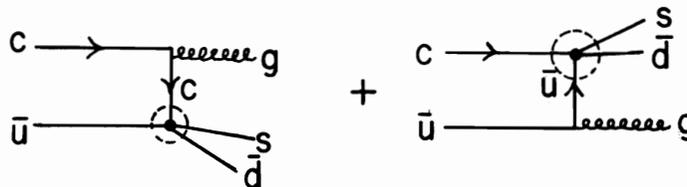

Fig. 18 Possible mechanism for D^O decay enhancement

In the Orsay non-relativistic quark model wave function overlap integrals give[101]

$$<K \, |s\dagger c| \, D > \ > \ <\pi \, |d\dagger c| \, D > \ ,$$

and from experiment we know[102,103] that $f_K > f_\pi$. Another SU(3) breaking mechanism could[104] arise from the penguin diagram of Fig. 20 which has $|\Delta U|=0$ and vanishes in the limit of s, d mass degeneracy; similar diagrams are thought[100]
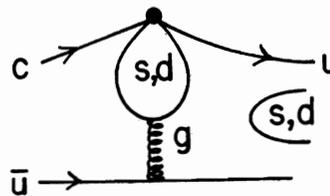
Fig. 20 Penguin diagram for 2-body D-decay

to be important in enhancing the $\Delta I=1/2$ amplitudes in K-decay. However, the Feynman integral which determines the Fermi coupling constant of the effective four-quark operator gives[73] (in lowest order and neglecting color-factors; the result of a more correct[100] treatment is similar):
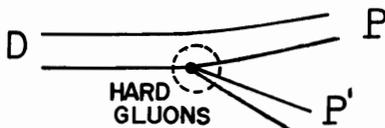
Fig. 19 Quark model diagram for 2-body D-decay

$$G_P^c = \frac{8\alpha_s}{\pi} \frac{G_F}{\sqrt{2}} \int_0^1 dz \, z(1-z) \, \ln \frac{M_s^2 - k^2 z(1-z)}{M_d^2 - k^2 z(1-z)} - \qquad (5.13)$$

Since the average momentum k^2 transmitted by the gluon in charm decay is characterized by the c mass:

$$k^2 \sim M_c^2 \gg M_s^2, M_d^2$$

we get $G_P^c \propto (M_s^2 - M_d^2)/M_c^2$ instead of $G_P^s \propto \frac{1}{6} \ln (M_c^2/M_K^2)$ for K-decay[105].

b) Final state interactions[103] could modify the simple picture of Fig. 14 and could enhance the K^+K^- final state if for example there were[97] a nearby spin zero resonance which is mostly $\bar{s}s$. In the absence of final state interactions, the model of Fig. 15 predicts[89,90]

$$\Gamma(D^0 \to K^- \pi^+) \simeq 40\Gamma(D^0 \to \bar{K}^0 \pi^0) \simeq 1.6\Gamma(D^+ \to \bar{K}^0 \pi^+). \tag{5.14}$$

A test of the last relation depends on whether the reported[93] lifetime difference turns out to be real, but the relation between the charged and neutral final states in D^0 decay is apparently not satisfied[106] suggesting that final state interactions (or some other mechanism) is indeed at work.

c) Deviations from the GIM current[102] are certainly expected through mixing with top and bottom quarks, but the limits on mixing angles discussed above suggest that these effects should be small. Turning the problem around, if one neglects SU(3) breaking effects, SU(3) sum rules can be used with the data to measure the heavy quark mixing angles. For example one gets[107]:

$$\frac{\Gamma(D^+ \to \pi^+ \pi^0)}{\Gamma(D^+ \to \bar{K}^0 \pi^+)} = |\theta_{cd}|^2, \quad \Gamma(D^0 \to K^0 \bar{K}^0) \propto |\theta_{cd} + \theta_{us}|^2 \tag{5.15}$$

The presently measured decay modes do not allow a direct extraction of these quantities because they involve a superposition of $|\Delta U|=1$ and $|\Delta U|=0$ amplitudes. Instead one can exploit triangle inequalities to obtain[108]:

$$0.20 \pm 0.06 \le |\theta_{cd} - \theta_{us}| \le 0.54 \pm 0.06 \tag{5.16}$$

for a quantity which reduces $2|\theta_{us}| = 0.46$ in the GIM four flavor limit. This result shows that the measured deviation from (5.11) does not require large t,b mixing, but in order to understand a ratio as large as

$$\frac{\Gamma(D \to K^+K^-)}{\Gamma(D \to \pi^+\pi^-)} \simeq 3$$

without invoking SU(3) breaking nor large mixing would require a rather large enhancement of the $|\Delta U| = 0$ amplitude relative to $|\Delta U| = 1$. A possible mechanism is via the penguin diagram of Fig. 20 with s,d replaced by b, giving

$$\frac{G_{Fig\ 20}}{G_{Fig}} \sim \ln(M_t^2/M_c^2) \left| \frac{\theta_{cb} \theta_{bu}}{\theta_c} \right| \simeq (0.28 \, {}^{+.21}_{-.21}) \, |\theta_{bc}| \ln (M_t^2/M_c^2). \tag{5.17}$$

While the logarithmic enhancement factor in (5.17) may be somewhat more important than the analogous factor in $K \to \pi\pi$, there is a double Cabibbo suppression and in addition the matrix element enhancement of the penguin operator is probably less important than for K-decay. The point is that matrix elements of the four fermion V-A operator are suppressed by approximate chiral symmetry, while the penguin operator, which can have a right handed quark coupled to the gluon, is

not[100]. Chiral symmetry is less relevant for charm decay; using a valence quark model for the matrix elements one loses a factor M_s/M_c in going from K-decay to charm decay.

It may well be that some combination of the[109] effects discussed above contrive to suppress the $\pi\pi$ mode relative to the prediction (5.11). At any rate the present state of the data and of theoretical understanding do not yet allow the conclusion that a more interesting mechanism like Higgs meson exchange need be invoked.

5.3 Strangeness changing decays. The long standing issues in kaon and hyperon decays are why $|\Delta I| = 1/2$ amplitudes are enhanced by typically a factor 20 in amplitude relative to $|\Delta I| = 3/2$ amplitudes, and why decay rates are enhanced by roughly the same factor relative to non-leptonic rates. There are now high statistics data on Ω^- decay, and in particular the result[110]

$$\frac{\Gamma(\Omega^- \to \Xi^0 \pi^-)}{\Gamma(\Omega^- \to \Xi^- \pi^0)} = 2.94 \pm 0.35 \tag{5.18}$$

to be compared with the $\Delta I = 1/2$ prediction of 2.0. The result (5.18) represents a deviation from the $\Delta I = 1/2$ rule of about 20%, considerably larger than previously measured deviations of 5%. This result was actually predicted[111] using the various tools of QCD and QAD phenomenology which have been developed through attempts to understand kaon and hyperon decay. Several effects have been found which favor $|\Delta I| = 1/2$ amplitudes[84]:

a) Hard gluon corrections to the V-A four fermion coupling enhance[112] $|\Delta I| = 1/2$ amplitudes relative to $|\Delta I| = 3/2$ by about a factor of 4.

b) In the non-relativistic quark model, matrix elements of the $|\Delta I| = 3/2$ part of the effective weak hamiltonian vanish[113] between baryon states; these matrix elements give the dominant contribution in the chiral limit.

c) Penguin diagrams, which contribute only to $\Delta I = 1/2$ transitions[112,89,100] have enhanced matrix elements[100] if the four fermion operator acts on both valence quarks in the pseudoscalar wave function, because of the chiral properties discussed above.

Recently two groups[114] have analyzed kaon and hyperon decay using the above ingredients with standard PCAC techniques[115] and the MIT bag model[116] to estimate matrix elements. With no parameters to be fitted they find a satisfactory description of kaon and s-wave baryon decay amplitudes, but p-wave amplitudes are generally too small by a factor of about one half. Their general conclusion is that all of the above effects play a role in the observed $|\Delta I| = 1/2$ enhancement, and there is no single predominant effect. In addition, the Ω^- decay[110,117] rates are adequately described[11]; they turn out to be predominantly p-wave with small decay asymmetries as confirmed[110] for the ΛK final state.

My own conclusion is that the $\Delta I = 1/2$ rule is understood within the current theoretical framework and the overall strength of non-leptonic amplitudes are also understood, although the details of their relative strengths is not yet fully accounted for by the theory. In particular, the failure to describe adequately both s- and p-waves in hyperon decay is an old problem[118] which emerged from PCAC analysis as a consequence only of the V-A nature of the primary interaction and approximate chiral symmetry, properties

which remain unaffected by the more recent developments.

6. Conclusions

My principal conclusion is that there is an honest theory of weak interactions, which I consider to be a major accomplishment of the past decade. There is still much work to be done in pinning down the more elusive aspects of the theory and understanding better the dynamics.

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Discussion

Q. (Roos, Helsinki) You state at the beginning quite correctly that $SU(2)_{left} \times SU(2)_{right} \times U(1)$ was not relevant at the present energies. Just to give you a quantitative figure if one takes any one of those models they have two Z bosons. One of them comes out to be about the same as the Weinberg-Salam model and the other one now comes out to be with 95% confidence heavier than 200 GeV.

A. Oh, that's a nice result.

Q. (Sanda, Rockefeller) I have a comment on D decays to KK and $\pi\pi$. My comment is on the relevance of so-called penguin graphs to this particular process. You have indicated that the momentum transferred through the gluon is small. Therefore, the factor $\ln [(ak^2 + m_s^2)/(ak^2 + m_d^2)]$ is near zero. It seems that for the 2 body decay case like this is very hard to rigorously argue the case for a zero. In fact, if you consider the fact that this is proportional to the running coupling constant $\alpha(k^2)$, it might perhaps favor small k^2. So that my statement is that one should not naively dismiss

this kind of a possibility and whether the diagram is relevant or not should be up to further experimental studies.

A. Okay, my own feeling is that on the average it should be something large but, in any case, you can't calculate reliably that contribution because you will get zero in the usual approximation.

Q. So it should be left up to the experimentalist to verify whether it's zero or not.

Q. (Rosen, Los Alamos) I'd like to point out that there's one important aspect in which the standard Salam-Weinberg model has not yet been subjected to a really severe experimental test. And that is in the absolute sign of the amplitudes. From ν-quark and e-quark neutral-current scattering, one can determine the coupling constants up to an overall sign. This sign can be measured by studying the interference between charged- and neutral-currents. The only experiment that speaks to that issue at all is the reactor experiment of Reines and company and if you examine the data very carefully you find that the errors on the experiment at this time are really much too large to decide the issue one way or the other.

Q. (Schopper, DESY) You didn't mention any models where you have a second 1/3 quark instead of a t-quark. Could you comment on that?

A. No. I have little to say. My prejudice is strongly in favor of the doublet structure.

Q. (Rosner, Minnesota) Can you comment on experimental limits e^+e^- going to H^+H^- and what is the lower limit of the Higgs mass that such experiments exclude.

A. I'm not sure. The contribution to R is about 1/4. Is there anybody here from DESY who could comment on this? It would be nice to set a limit, I agree with you.