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Electroweak Interactions

J.D. BJORKEN
Fermi National Accelerator Laboratory, Batavia, Illinois 60510

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Introduction

These lectures are not intended to be a systematic exposition or comprehensive review of the electroweak interaction. Instead we shall present a point of view which begins phenomenologically and moves in stages toward the conventional gauge theory formalism containing elementary scalar Higgs-fields and then beyond.

Our purpose in so doing is that the success of the standard $SU(2) \times U(1)$ theory in accounting for low energy phenomena need not automatically imply success at high energies. It is deemed unlikely by most theorists--and I would not disagree--that the predicted $W^{\pm}$ or Z^{0} does not exist or does not have the mass and/or couplings anticipated in the standard model. However, the odds that the standard predictions will work are not 100%. Therefore there is some reason to look at the subject as one would were he forced by a "wrong" experimental outcome--to go back to fundamentals and ascertain what is the minimal amount of theory necessary to account for the data.

Another reason for a conservative approach is that even upon accepting the gauge-theory ideology--which is a priori a nearly irresistible thing to do--one is left with an awkward system of spinless Higgs-particles. Some of the awkwardness can be seen by setting all gauge couplings in the Lagrangian to zero. The interactions left behind in the standard model are an ad hoc array of nonlinear Higgs-boson self-couplings and Yukawa couplings to fermions. The oft-stated refrain that "all known interactions appear to be described by gauge couplings" is simply not true in the orthodox theory.

It would be nice to replace the Higgs system by one also based on gauge couplings. This is an active field of study these days. Schemes in which the Higgs bosons are composites, with new gauge-type couplings ("technicolor" or "hypercolor") providing the binding, contain underlying ideas which are very attractive. But as yet no realistic model has been produced.

These notes are organized as follows: in Chapter 1, we discuss and motivate the phenomenological weak Lagrangian needed to correctly describe low energy weak-interaction phenomena. This much does not even require intermediate-vector bosons. We then introduce the intermediate-boson hypothesis, and finally the additional assumptions needed ("unification hypothesis") to obtain the predictions for m_W and m_Z.

In Chapter 2, we adopt the gauge theory ideology and develop the standard model. The Higgs-boson system is viewed in analogy to the linear σ-model sometimes used to describe low energy π - π interactions. This leads to a formalism which appears a little different from that usually found in the literature. While it is in fact equivalent, there may be some advantage in looking at the same phenomenon from a slightly different viewpoint. This viewpoint also offers an easy point of departure for the technicolor (hypercolor) models. These are also described.

Chapter 3 is devoted to a very brief sketch of experimental implications. Chapter 4 is devoted to a description of the rich phenomenology of technicolor models, while Chapter 5 explores some of the impressive and promising ideas and concepts underlying the technicolor models.

1. The Phenomenological Lagrangian

1.1. The Generalized Fermi Interaction

1.1a. Structure. The effective Lagrangian for the charged-current weak interactions, valid at low energies, has conventionally been taken to be [1]

$$\mathscr{L}_{eff}(x) = \frac{G_F}{\sqrt{2}} J_\mu(x)^\dagger J^\mu(x) \tag{1.1}$$

where the charged current J_μ is composed of several pieces, each with V-A structure

$$\begin{aligned} J_\mu = {} & \bar{e}\gamma_\mu(1-\gamma_5)\nu_e + \bar{\mu}\gamma_\mu(1-\gamma_5)\nu_\mu + \bar{\tau}\gamma_\mu(1-\gamma_5)\nu_\tau \\ & + \bar{d}'\gamma_\mu(1-\gamma_5)u + \bar{s}'\gamma_\mu(1-\gamma_5)c + \ldots \quad . \end{aligned} \tag{1.2}$$

The primed superscripts indicate the existence of mixing of the quark species

$$\begin{aligned} d' &= d\cos\theta_c + s\sin\theta_c \\ s' &= s\cos\theta_c - d\sin\theta_c \end{aligned} \tag{1.3}$$

The Cabibbo angle θ_c is (consistently) measured to be [2]

$$|\sin\theta_c| = 0.219 \pm .002 \pm .011 \text{ (systematic)} \approx \sin^2\theta_W \quad . \tag{1.4}$$

This "classical" view of charged current weak interactions--with which we assume that the reader is familiar--is nowadays expected to be slightly modified, owing to the existence of the bottom quark b and to its possible mixings with the s and d quarks. The simplest generalization supposes existence of a charge -2/3 top-quark as partner to the bottom [3]. If this hypothesis is adopted, the natural generalization of Eq. (1.2) adds one more term $\bar{b}'\gamma_\mu(1-\gamma_5)t$ to the current. More succinctly,

$$\mathcal{J}_\mu = \bar{\ell}\gamma_\mu(1-\gamma_5)\nu + \bar{U}\gamma_\mu(1-\gamma_5)D' \tag{1.5}$$

with

$$\ell = \begin{pmatrix} e \\ \mu \\ \tau \end{pmatrix} \qquad \nu = \begin{pmatrix} \nu_e \\ \nu_\mu \\ \nu_\tau \end{pmatrix}$$

$$U = \begin{pmatrix} u \\ c \\ t \end{pmatrix} \quad D' = \begin{pmatrix} d' \\ s' \\ b' \end{pmatrix} = \begin{pmatrix} V_{dd} & V_{ds} & V_{db} \\ V_{sd} & V_{ss} & V_{sb} \\ V_{bd} & V_{bs} & V_{bb} \end{pmatrix} \begin{pmatrix} d \\ s \\ b \end{pmatrix} . \tag{1.6}$$

The 3×3 unitary matrix $\mathcal{V}$, known as the Kobayashi-Maskawa matrix [4], now describes the mixings of d, s, and b. It contains 4 independent parameters, including a possible CP-violating phase angle.

This picture of weak interactions is incomplete. Neutral currents exist as well as charged currents [5]. A most natural hypothesis for the structure of the natural currents is simply to

(i) consider each pair of left-handed fermions as a weak isotopic-spin doublet:

$$f_i = \begin{pmatrix} \nu_e \\ e^- \end{pmatrix}, \begin{pmatrix} \nu_\mu \\ \mu^- \end{pmatrix}, \begin{pmatrix} \nu_\tau \\ \tau^- \end{pmatrix}, \begin{pmatrix} u \\ d' \end{pmatrix}, \begin{pmatrix} c \\ s' \end{pmatrix}, \begin{pmatrix} t \\ b' \end{pmatrix} . \tag{1.7}$$

(ii) Construct the full isotopic-spin current

$$\vec{J}_\mu = \sum_{i=1}^{6} \bar{f}_i \gamma_\mu \left(\frac{1-\gamma_5}{2} \right) \frac{\vec{\tau}}{2} f_i \tag{1.8}$$

and write down an invariant current-current interaction which generalizes Eq. (1.1):

$$\mathscr{L}_{weak} = 2\sqrt{2} G_F \vec{J}_\mu \cdot \vec{J}^\mu \quad . \tag{1.9}$$

This form, which was proposed long, long ago [6], does not work. It predicts a pure left-handed V-A structure for neutral-current processes, while the measurement of neutral currents in neutrino-induced reactions gives clear evidence for a right-handed component. However, one does not have to go too far to find a modification which does work. It was also proposed long ago [7], and is simply an electromagnetic correction coming from photon exchange, as shown in Fig. 1. The first-order weak correction to the electromagnetic vertex of each fermion will (in the limit of negligible fermion mass and reasonably low energy) be proportional to q^2, thereby cancelling the q^{-2} of the photon propagator. This term also will necessarily possess the V-A structure at the weak vertex. Thus such terms will (provided they are universal) lead to an additional effective Lagrangian of the form

$$\mathscr{L}'(x) = -4\sqrt{2} G_F J_\mu^{(3)}(x) J_{em}^\mu(x) x_W \quad . \tag{1.10}$$

The coefficient x_W is here chosen by convention to reduce to $\sin^2\theta_W$ if the Lagrangian of the standard SU(2) × U(1) electroweak gauge theory is adopted. The full neutral-current Lagrangian we have discussed is then

$$\mathscr{L}^{NC} = 2G_F\sqrt{2}\left(J_\mu^{(3)} J^{\mu(3)} - 2x_W J_\mu^3 J^\mu_{em} \right) \tag{1.11}$$

while the Weinberg-Salam SU(2) × U(1) effective Lagrangian for neutral-current processes is [8]

$$\mathscr{L}^{WS} = 2G_F\sqrt{2}\left(J_\mu^{(2)} - x_W J_\mu^{em} \right)^2 \quad . \tag{1.12}$$

The difference between the two is a purely electromagnetic contribution, proportional to $J_\mu^{em} J^\mu_{em}$. Such a contribution is provided by conventional vacuum-polarization insertions (cf. Fig. 2) to the photon propagator, where the virtual states have intrinsic mass large compared to q^2. [The states such as e^+e^-, $u\bar{u}$, etc. with intrinsic mass small compared to q^2 contribute to the running electrical charge.] The existing body of data is not sensitive to this latter term; hence Eq. (1.11), unmotivated by gauge-theories, provides the same description of the low-energy phenomena thus far measured as does the standard SU(2) × U(1) model [9]. However, it is possible to measure (or bound) the extra electromagnetic term via parity conserving electroweak interference measurements such as in $e^+e^- \to \mu^+\mu^-$; this will be mentioned again in Section 1.2.

A most important feature of the neutral-current Lagrangian, Eq. (1.11), is that all terms are flavor-diagonal, i.e. the (generalized) GIM mechanism [10] applies. That is, only the combination

$$\bar{d}'\gamma_\mu\left(\frac{1-\gamma_5}{2}\right)d' + \bar{s}'\gamma_\mu\left(\frac{1-\gamma_5}{2}\right)s' + \bar{b}'\gamma_\mu\left(\frac{1-\gamma_5}{2}\right)b' =$$

$$= \bar{d}\gamma_\mu\left(\frac{1-\gamma_5}{2}\right)d + \bar{s}\gamma_\mu\left(\frac{1-\gamma_5}{2}\right)s + \bar{b}\gamma_\mu\left(\frac{1-\gamma_5}{2}\right)b \tag{1.13}$$

or the electromagnetic current occurs in the effective Lagrangian. The removal of primes in Eq. (1.13) is made legitimate by the unitarity of the K-M matrix defined in Eq. (1.6).

1.1b. Evidence. The conventional body of textbook information on charged-current weak interactions, together with the more recent measurements of neutral-current phenomena which give such impressive agreement with the standard model, provides more than enough evidence to motivate the hypothesis of the current-current Lagrangians embodied in Eqs. (1.1) and (1.11). Nevertheless, it is clearly desirable to check wherever possible the consistency of that Lagrangian with data, as well as to explore for additional unanticipated terms. Just in the charged-current Lagrangian there are 45 distinct terms (omitting those involving the hypothetical top quark). These terms are classified in Table I. We have used a very subjective Michelin star system to rank the experimental status of the entries. According to that ranking, about 16 out of the 45 candidates have been checked to some extent for existence, V-A structure, and/or reasonable normalization of strength. Of the remainder, about 13 more can be checked in the foreseeable future. Among the candidates are:

1) $\tau \to \nu_\tau s\bar{u}$: This Cabibbo-suppressed term should be measurable [11] in $\tau \to K\nu_\tau$, or $\tau \to K^* \nu_\tau$. The branching ratios should be reasonably calculated, but a check of V-A structure may be rather remote.

2) $\tau \to \nu_\tau s\bar{c}$: The decay $F \to \tau\nu_\tau$, important for beam-dump experiments designed to search for existence of the ν_τ, is dependent upon the presence of this term.

3) $b \to ue\bar{\nu}_e$, $b \to u\mu\bar{\nu}_\mu$, $b \to u\tau\bar{\nu}_\tau$ (??), $b \to ce\bar{\nu}_e$, $b \to c\mu\bar{\nu}_\mu$, $b \to c\tau\bar{\nu}_\tau$ (???): These processes are prime material to be studied at e^+e^- "bottom-factories" such

as CESR [12]. The information is likely to be rather indirect for some time to come. The partition between $b \to u$ and $b \to c$ transitions is an especially important issue.

4) $c \to de^+\nu_e$, $c \to d\mu^+\nu_\mu$: These Cabibbo forbidden processes may eventually be measured in charm-factories such as SPEAR or DORIS. Already dilepton production by $\bar{\nu}_\mu$ provide some information on the existence of the second of these processes [13].

5) $b \to c\bar{u}d$, $b \to u\bar{u}d$: These terms should dominate the nonleptonic decay of the b; again the fraction of events containing charm is important. Nevertheless the fraction of nonleptonic decays containing charm need not be the same as the fraction of semileptonic decays containing charm. Strong-interaction corrections can mess things up [3].

6) $b \to c\bar{c}s$: This channel, Cabibbo-allowed but probably somewhat phase-space-suppressed, has some fascinating decay channels [14], in particular $b \to s\psi$.

There are 66 distinct (flavor-diagonal) neutral current processes, even omitting the top-quarks. The vast majority of these are quite inaccessible at low energies. The neutral-current terms are given their Michelin-tabulation in Table II. One sees that up to now only 8 out of 66 have been checked to some extent. The best information is on $\nu_\mu u$ and $\nu_\mu d$ couplings, which have been essentially uniquely determined from the data and which are in accord with the Weinberg-Salam effective Lagrangian, Eq. (1.11). This has been repeatedly reviewed [15], and we shall not digress here to do it again. The eu and ed parity-violating neutral current couplings have been detected in the beautiful SLAC-Yale experiment on electro-weak interference in electron-nucleon inelastic scattering [16], although a complete determination of eu and ed neutral current couplings remains for the future. Perhaps atomic parity-violation experiments in H and D will be needed to complete that program [17]. The $\nu_\mu e$ coupling has been seen, but not too accurately determined in elastic $\nu_\mu e$ scattering [18]. Likewise, $\nu_e e$ scattering has been observed via reactor antineutrinos [19].

Opportunities do exist for adding to this list of observed neutral current couplings; in Table II there exist 9 such entries. Electroweak interference experiments with incident muons should give information on μu and μd neutral current couplings. Such information may also come from parity violating asymmetries in Drell-Yan dilepton production by hadron beams. Measurement of parity violating asymmetries in Møller scattering of polarized electrons has been discussed as a SLAC fixed-target experiment and might eventually be done there. Measurements of e^+e^-, $\mu^+\mu^-$, and $\tau^+\tau^-$ production in e^+e^- annihilation (at higher energies) should show evidence of ee, $e\mu$, and $e\tau$ neutral currents. And diffractive production (elastic or inelastic) of ϕ, ψ, and Υ by ν_μ might ultimately be an observable dilepton signal in neutrino experiments, providing some evidence for $\nu_\mu s$, $\nu_\mu c$, and $\nu_\mu b$ neutral current couplings.

How important is it to fill out this list of observation of dozens of neutral or charged current couplings? The answer depends a great deal on the existence of the intermediate bosons $W^\pm$ and Z^0 conjectured (better, expected) to mediate these weak interactions. If they do exist with the anticipated properties, then it suffices to measure the n ($n \sim 10$) couplings of the bosons to these fermion channels rather than the $\sim n^2/2$ four-fermion couplings consequent from the virtual W or Z exchange. On the other hand, if something goes wrong with the standard picture, the catalogue of measurements we have sketched would take on considerably greater importance.

It is perhaps worth emphasizing that accurate measurments of x_W (i.e. $\sin^2\theta_W$) in various processes may be quite important [20]. From the point of view expressed above, namely that x_W is related to the electromagnetic charge radii of the participating fermions, it might not be unreasonable to expect different fermions to be characterized by different charge-radii, and hence by different values of x_W. [In the standard model the charge-radius is dominated by the Z^0, and is universal. This will be discussed later on.] At present, accurate measurements of x_W for u, d, and ν_μ exist.

1.1c. Possible deviations from the standard Lagrangian. Thus far we have mainly discussed the scope of the (implicitly favorable) evidence regarding the structure of the conventional effective Lagrangian. But is it true that existing data all supports the standard picture? Is it all well-understood? The answer is not completely positive, and some of the troublesome areas are listed below:

1) Atomic parity violation: Two out of four experiments are in agreement with the standard model, while the remaining two still do not see a large enough effect. Progress has been made in the experiment on Cs, and someday we may hope to see data from H and D. The measurements complement the SLAC-Yale measurements, and are therefore of interest in their own right.

2) Nuclear parity violation: A recent measurement of Ramsey and collaborators [21] of the parity violating "Faraday" rotation of the transverse polarization of a slow neutron in passage through various isotopes of tin has shown a very large effect for Sn^{117} but not for Sn^{124}. The effect seems two or three orders of magnitude larger than nominally expected from the parity-violating $\Delta S = 0$ nonleptonic weak interaction. It is premature to say whether this is some artifact of the nuclear physics or something more fundamental. Nevertheless, it is a topic worth watching. In this connection, one should recall also the large γ-ray circular polarization measured by Lobashov [22] in the process $n + p \rightarrow d + \gamma$.

3) <u>$\Delta S = 1$ nonleptonic decays</u>: This subject has never really been satisfactorily understood, especially the success of the $\Delta I = ½$ rule. There has always been the issue of dynamical enhancement versus an intrinsic selection-rule, with the standard effective Lagrangian, Eq. (1.1), arguing for dynamical enhancement [3]. Nevertheless, especially for p-wave hyperon decays, it has been difficult to understand why the $\Delta I = ½$ rule should be so accurate. A new argument in favor of dynamical enhancement comes from the recent measurements of nonleptonic Ω^- decays. The branching ratio [23]

$$\frac{\Gamma(\Omega^- \to \Xi^0 \pi^-)}{\Gamma(\Omega^- \to \Xi^- \pi^0)} = 2.94 \pm 0.35 \tag{1.14}$$

is in disagreement with the $\Delta I = ½$ rule prediction of 2.03 and gives the most substantial violation so far measured.

QCD has provided some rationales for $\Delta I = ½$ enhancement. First of all, gluonic radiative corrections enhance the $\Delta I = ½$ and suppress the $\Delta I = 3/2$ part of the effective Lagrangian but only by a relative factor $\sim$ 2-3. But in addition, the so-called "penguin" diagrams [24] (Fig. 3) which are pure $\Delta I = ½$ have been argued to provide an additional, possibly dominant contribution [3]. However, this contribution strains the limits of applicability of perturbative QCD, and it is not clear to me that it is the ultimate answer to the $\Delta I = ½$ question.

4) <u>$\Delta S = 1\ \Sigma^- \beta$ decay</u>: Existing measurements of Σ^- β-decay do not support the expected Cabibbo structure of this matrix element of the weak current, in particular the relative sign of the vector and axial contributions [25]. Improved measurements can be expected eventually, and it again may be a subject worth watching.

In addition to checking that expected terms in the weak Lagrangian are present, it is also of interest to check that unexpected terms are indeed absent. With the number of allowed terms in $\mathscr{L}$ of order 100, the number of absent, or disallowed terms is much greater. Without some theoretical guidance, it is hard to discuss this question in a very meaningful way. Nevertheless, some broad classifications can be discerned even in the absence of any theoretical motivation.

1) Right-handed currents: Various experiments (e.g. positron or muon helicity measurement; asymmetry parameters in μ decay; violations of universality) limit the amount of V + A admixture to the ordinary V - A weak interaction to less than a few percent (in rate). There is still considerable room for improvement. For example the completeness of the muon polarization in $K_{\mu 2}$ and $K_{\mu 3}$ decays has only been determined to the 10-20% level [26] . Were a right-handed $\Delta S = 1$ current not Cabibbo-suppressed, these would be rather sensitive channels for a search.

The recent revival of interest in ep colliding beams has focussed attention on this process as a sensitive way of searching for right-handed weak currents. One looks for $e^- + p \rightarrow \nu$ + hadrons with a right-handed electron incident. But again it may be hard to go beyond a level of a few percent of the left-handed weak process [27].

2) Helicity-flip currents (S, P, T): Even within a standard gauge-theory framework, exchange of Higgs-bosons can provide additional effective four-fermion interactions [28]. The strength and structure of such terms is very uncertain. Some focus has been provided by the technicolor (or hypercolor) models, to which we return later in these lectures.

Probably the best experimental limits on such terms come from the decays $\pi \rightarrow e\nu_e$, $K \rightarrow e\nu_e$, and $K_L \rightarrow \mu^+\mu^-$ (or e^+e^-), as well as the $K_L - K_S$ mass difference, all of which are very sensitive to admixtures of S and P currents.

3) Selection-rule violation: This can include decays such as $K_L \rightarrow \mu e$, $K^{\pm} \rightarrow \pi^{\pm} \nu\bar{\nu}$, $\mu \rightarrow e\gamma$, which might be mediated either by Higgs quanta or by massive vector bosons. Again, this issue will be addressed later in these lectures. It is clearly of importance to push limits on such flavor violating decays as far as possible.

There are other selection rules that would be nice to test. For example, while $\Delta S = 1$ weak interactions occur, and $\Delta S = 2$ transitions are second-order weak at best (from the small size of $K^0 - \bar{K}^0$ mixing), not much is known about $\Delta S = 3$. For example, does $\Omega^- \rightarrow p\pi^-\pi^-$ decay occur? With the advent of intense charged hyperon beams, it should be possible to put a pretty good limit on $\Delta S = 3$ transitions.

The heavy quarks c and b and the heavy lepton τ may also be good candidates for unexpected decay modes. If such processes are correlated with the Higgs sector or whatever mechanism exists for fermion mass-generation, the relevant amplitudes may have a strong mass-dependence and be more dominant. Decays such as $\tau \rightarrow$ eee, μee, etc. or $b \rightarrow s\tau\mu$, $s\mu e$, etc. should be observable experimentally if a reasonable coupling strength exists [29]. Two other selection rules, $\Delta C = 2$, and $\Delta B = 2$ (B here means bottom or beauty) are analogous to the $\Delta S = 2$ selection rule tested in the $K_L - K_S$ mass mixing. Again, sensitive tests are provided by the absence (or presence) of $D^0 - \bar{D}^0$ or $B^0 - \bar{B}^0$ mixing. Threshold production in e^+e^- annihilation is probably the technique of choice here.

A non-Abelian selection rule which deserves a test is the $\Delta I = 1$ rule for Cabibbo-allowed charm decays. For example, ratios such as $\Gamma(D^+ \rightarrow \bar{K}^0\pi^+)$: $\Gamma(D^0 \rightarrow K^-\pi^+)$: $\Gamma(D^0 \rightarrow \bar{K}^0\pi^0)$ must satisfy triangle inequalities and should eventually be directly measurable.

1.2. The Intermediate Boson Hypothesis

On the whole, the evidence for the Weinberg-Salam effective Lagrangian is very good. Let us here set aside any residual conservatism and skepticism, and accept the validity of that Lagrangian. What are the implications? First of all, the current-current form and universal (?) structure strongly suggests that the weak force is mediated by a triplet $W^{\pm}$, Z^{o} of intermediate $J = 1$ bosons coupled universally. Furthermore, if one were to set x_W equal to zero one would obtain a weak SU(2) isospin symmetry, implying that in that limit M_W and M_Z have the same mass. The couplings of $\vec{W}$ to the fermions may be written

$$\mathscr{L} = g \sum_i \bar{f}_i \gamma_\mu \left(\frac{1 - \gamma_5}{2} \right) \frac{\vec{\tau}}{2} f_i \cdot \vec{W} \qquad . \qquad (1.15)$$

Applying this form to, say, muon decay leads directly to the identification

$$\frac{g^2}{m_W^{\ 2}} = 4\sqrt{2} G_F \qquad . \qquad (1.16)$$

In the previous section, we indicated that the terms in the neutral current proportional to x_W can be considered as charge-radius contributions. We now may regard this charge-radius as vector-dominated by the W_3. That is, if W_3 and the photon mix, with mixing amplitude eF, then the diagram in Fig. 4 leads to the weak amplitude [30]

$$\mathscr{M} = \bar{f}_i \gamma_\mu \left(\frac{1 - \gamma_5}{2} \right) \frac{g\tau_3}{2} f_i \cdot \frac{1}{m_W^{\ 2} - q^2} \cdot eF \cdot \frac{(-1)}{q^2} \cdot eJ^{\mu}_{em} \qquad . \qquad (1.17)$$

Expanding out the propagator to obtain the contact term, and then relating this to the neutral current Lagrangian in Eq. (1.11) gives

$$\frac{ge^2F}{m_W^4} = 4\sqrt{2}G_F x_W \quad . \qquad (1.18)$$

There is one additional effect of importance. The mixing of W_3 and photon produce a charge renormalization, as well as a mass-shift of the W_3. This comes about by the geometric summation of W_3 insertions in the photon propagator. The renormalized propagator is

$$D(q^2) = \frac{e_o^2}{q^2\left[1 - \frac{e_o^2F^2}{m_W^2(q^2 - m_W^2)}\right]} = \frac{e_o^2(q^2 - m_W^2)}{q^2(q^2 - m_Z^2)} \qquad (1.19)$$

with

$$m_Z^2 = m_W^2 + \frac{e_o^2F^2}{m_W^2} \quad . \qquad (1.20)$$

Elimination of $e_o \equiv Z_3^{-1/2}e$, g and F from the above equations leads to the relations [31]

$$m_W = \frac{37 \text{ GeV}}{\sin^2\theta_W}\sqrt{1 - Z_3} \leq \frac{37 \text{ GeV}}{x_W}$$

$$\frac{m_W}{m_Z} = \sqrt{Z_3} \quad . \qquad (1.21)$$

However, Z_3 is not yet determined.

What are the consequences of this IVB hypothesis? One is that the sign of G_F is determined. This has in fact been measured in the SLAC polarized electroproduction experiment. There the sign of the intrinsic weak interaction, as well as x_W, is determined; it indeed agrees with the IVB hypothesis.

Another consequence of this analysis is that the mass of the charged boson $W^{\pm}$ must be ≤ 150 GeV. This in fact can be considerably generalized [9]. No matter what set of weak quanta are exchanged in order to build the four-fermion effective Lagrangian, the range of the charged-current weak interactions must be greater than $(150 \text{ GeV})^{-1}$. This conclusion is rather academic within the context of the gauge theories. However, if they are not correct and W's and Z's are not found in the $p\bar{p}$ colliders or LEP, then such a bound could become very relevant. The charged current weak interaction measured in electron-proton colliding beams would then necessarily have to show manifestations of the nonvanishing range of the weak force, no matter what the weak quanta mediating the force eventually turned out to be.

Unfortunately, I know of no similar statement for the neutral Z^{o} boson. There do exist generalizations [32] of the standard $SU(2) \otimes U(1)$ gauge-theory model which utilizes the group $SU(2) \otimes U(1) \times G$. In these models, which contain one $W^{\pm}$ but many Z's, there is a center-of-gravity theorem which requires at least one of the Z's to have a mass no larger than the standard mass of 85-90 GeV. However, without using the gauge theory formalism, no similar statement appears.

Many people have studied possible models within the $SU(2) \otimes U(1) \times G$ framework. A nice, especially simple example has been given by de Groot and Schildknecht [33], who choose $G = U(1)$. The only effect is the mixing of the extra boson with W_3 and photon; it is chosen to couple not at all to the known fermions. From the existing PETRA QED tests, they find significant constraints on the masses of the two physical Z's. These are shown in Fig. 5.

Within this model it appears that it is difficult to have a Z^{o} with mass much less than 60-70 GeV. However, this is a model-dependent consequence, and one should not conclude from this that a search for lighter Z^{o}'s is fruitless.

Another conclusion of de Groot and Schildknecht is more model-independent. Within the general framework we have discussed (not even assuming the IVB hypothesis), we argued that the general structure of the neutral current Lagrangian is

$$\mathscr{L}_{NC} = 2\sqrt{2}G_F\left[\left(J_\mu^{(3)} - x_W J_\mu^{em}\right)^2 + \lambda x_W^2 \left(J_\mu^{em}\right)^2\right] \tag{1.22}$$

where the last term simply comes from extra contributions of weak quanta to the photon vacuum polarization (see the discussion above Eq. (1.21)). In the single IVB hypothesis (but not assuming gauge theories), it is not hard to show [31] that $\lambda = 0$. In more general circumstances, it is possible to show [9], via Schwartz-inequalities, that $\lambda \geq 0$; i.e. that the $x_W^2 J_{em}^2$ term should be enhanced by the factor $(1 + \lambda)$. This parity-conserving term interferes destructively with the standard photon-exchange term, and modifies $\sigma(e^+e^- \to \mu^+\mu^-)$ from its expected QED value. The conclusion from the analysis of de Groot and Schildknecht (using MARK J PETRA data) is that $\lambda < 3$. This bound should be considerably refined as more e^+e^- data is accumulated. It clearly implies a restriction on how much extra contributions to vacuum polarization--beyond the Z^o--can exist.

1.3. Unification Hypothesis

We have seen that the Weinberg-Salam effective Lagrangian plus the intermediate-boson hypothesis does not produce all the results of the gauge theory. What is missing? It is the statement of exact SU(2) symmetry at short distances, as well as the approximate SU(2) symmetry at large distances. Evidently the intrinsic weak interaction satisfies this symmetry condition. Photon exchange does not; in particular the charge-radius terms must be examined with care. Let us write out the electromagnetic vertex of one of our fermions in detail:

$$J_\mu^{em} = \bar{f}_i \left\{ \gamma_\mu \left(\frac{1-\gamma_5}{2} \right) \left[Q + \frac{\tau_3}{2} \frac{q^2}{m_W^2} \frac{gF}{\left(m_W^2 - q^2 \right)} \right] + \gamma_\mu \left(\frac{1+\gamma_5}{2} \right) Q \right\} f_i . \quad (1.23)$$

Here $Q = (\tau_3/2) + Y$ is the charge-operator for the fermion in question, and Y is the weak hypercharge (i.e. the mean charge of the weak doublet). The form for the left-handed current written above is determined by the requirements that (1) the current reduce to the correct limit as $q^2 \to 0$, and that (2) the residue of the pole at $q^2 = m_W^2$ (this pole gets cancelled by the zero in the photon propagator [34]) be normalized according to the diagram in Fig. 4.

A sufficient condition for recovering the high-energy predictions of the Weinberg-Salam gauge theory is that the coefficient of τ_3 vanish as $q^2 \to \infty$. This gives

$$1 = \frac{gF}{m_W^2} \quad . \quad (1.24)$$

When combined with Eq. (1.18), this implies

$$m_W^2 = \frac{(37\ \text{GeV})^2}{x_W} \quad \text{and} \quad \frac{m_W^2}{m_Z^2} = 1 - x_W \quad (1.25)$$

along with determining the fact that

$$0 \leq x_W = \sin^2\theta_W \leq 1 \quad (1.26)$$

a result which, of course, is in accord with experiment, and provides some evidence in favor of the gauge theory per se.

The above condition, (Eq. 1.24), dubbed the unification hypothesis by Hung and Sakurai [31] , is a strong one and implies pole-dominance at all energies. This condition may be tantamount to assuming the gauge theory structure. Cornwall, Levin, and Tiktopoulos [35], among others [36], have shown that if one starts with the IVB hypothesis and demands a consistent S-matrix for the lowest-order diagrams (tree-unitarity) one is led to the gauge theories as an essentially unique solution. I have discussed this approach elsewhere [37], and will not repeat this here. However, it indicates that the addition of the unification hypothesis, at the least, comes close to adoption of the gauge theory in toto--although it has been done via a phenomenological route, as opposed to a construction from first principles, as usually done. It is possible that there is an advantage in doing things via the phenomenology. If for some reason the W and Z are composites, we may expect that pole dominance is at best a low-energy approximation. An instructive analogy is the low energy neutron-electron interaction, which has a four-fermion structure similar to the weak interaction. One normally views this from the point of view of photon exchange. However the neutron electromagnetic form factor is, to a fair approximation, dominated by the ρ^0 and ω contributions, so that one may also view the low energy neutron-electron interaction as mediated by ρ^0 and ω exchange. Furthermore, in the absence of any information at high energies, it would be (and in fact was) an attractive idea to consider the ρ^0 and ω as non-Abelian gauge fields [38]. Only upon going to energies large compared to m_ρ and m_ω did it become clear that that picture was wrong. It is also instructive to recall that at that time a model of ρ^0 and ω as bound states of nearly massless, fractionally charged, confined quarks looked like madness.

Is it possible that W and Z are composites, and that history can repeat itself? Certainly the possibility is there, but there are some obstacles to this idea. First of all, the intrinsic weak interaction is pure $I = 1$ exchange, i.e. ρ-like. If the $\vec{W}$ is made of constituents, say, having $I = \frac{1}{2}$, we may expect ω-like $I = 0$ exchange as well as $I = 1$ exchange. There is no evidence for such a contribution. Probably the best test comes from the deep-inelastic neutral-current data. Writing the ν_μ - u, ν_μ - d effective Lagrangian in the form

$$\mathscr{L} = \frac{G}{2\sqrt{2}} \bar{\nu}\gamma^\mu(1-\gamma_5)\nu \Big\{ \bar{u}\gamma_\mu(1-\gamma_5)u - \bar{d}\gamma_\mu(1-\gamma_5)d - 4x_W(2/3\,\bar{u}\gamma_\mu u - 1/3\,\bar{d}\gamma_\mu d)$$

$$+ \xi[\bar{u}\gamma_\mu(1-\gamma_5)u + \bar{d}\gamma_\mu(1-\gamma_5)d] \Big\} \tag{1.27}$$

and looking at the model-independent analyses [15], I conclude that the bound on $|\xi|$ from experiment is

$$|\xi| \lesssim 0.05 - 0.10 \quad . \tag{1.28}$$

That is a quite nontrivial bound on an additional I = 0 exchange.

A second obstacle comes from the large value of x_W or $\sin^2\theta_W$. We have viewed contributions proportional to x_W as electromagnetic in origin; hence proportional to α. It is therefore somewhat awkward that $x_W \gtrsim 0.2$ rather than $\lesssim 0.01$. This turns out to have a quite tangible consequence [9]. Because x_W is large, the weak quanta are strongly mixed with the photon. This implies that annihilation of a right-handed electron with a positron (a process mediated only by photon exchange) into weak quanta must be large, in proportion to x_W. In the most general framework, this leads to a bound as follows: Defining

$$R_R = \frac{\sigma(e^-_R e^+ \to \text{weak quanta})}{4/3\,\pi\alpha^2 s^{-1}} \tag{1.29}$$

we get [39] (c.f. Fig. 6)

$$\bar{R} = \int \frac{ds}{s} R_R = \frac{3\pi}{\alpha}\left(Z_3^{-1} - 1 \right) \geq \frac{3\pi}{\alpha} \left\{ \left[1 - \frac{m_W^2 \sin^4 \theta_W}{(37\ \mathrm{GeV})^2} \right]^{-1} - 1 \right\} \quad . \quad (1.30)$$

In the standard model, this bound is saturated by the Z^0 resonance. However if the Z^0 is built of constituents, we should expect a good fraction of the sum to be satisfied by the constituents themselves. However, this requires

$$\log \frac{\Lambda^2}{m_Z^2} \left[\sum_{\text{spin } \frac{1}{2}} Q_i^2 + \frac{1}{4} \sum_{\text{spin } 0} Q_i^2 \right] + \begin{matrix}\text{contributions from} \\ \text{spin} \geq 1 \text{ constituents}\end{matrix} \geq \bar{R} \quad . \quad (1.31)$$

[Here Λ is a cutoff in the integral, which could come from extra structure in the fermion vertex functions.] The function $\bar{R}$ is strictly bounded from experiment by 20 and has a nominal value $\geq 10^3 - 10^4$. We conclude that, if W and Z are built of constituents of spin 0 or ½, there probably are many of them, or else m_W is unexpectedly low. It seems that a straightforward repetition of past history holds some difficulty, and that a more imaginative model of compositeness is probably required.

2. The Problem of Mass-Generation

2.1. Intermediate Boson Mass

In the previous section we considered the weak interactions first from a strictly phenomenological point of view (i.e. low-energy Fermi current-current interaction), then from the intermediate-boson hypothesis and finally from the point of view of the "unification hypothesis." Each step was accompanied by less and less empirical evidence, but took us closer and closer to the standard gauge-theory picture. In this section we shall accept those three steps, along with the basic notions of the $SU(2) \times U(1)$ gauge theory. We shall not develop the gauge theory ideas in general, from first principles--that is done in quite a few places [40] , and in any case requires a bit of preparatory formalism. Suffice it to say that from the gauge theory point of view $W^{\pm}$ and Z^{o} are degrees of freedom just as elementary as the photon. There are two distinctions. One is that the electroweak gauge group is non-Abelian, implying trilinear and quartic self-couplings among the W's and Z's. (However, we shall for the most part not go deeply enough into the theory to see their effects). The second distinction is that, unlike the photon and QCD gluon octet, the $W^{\pm}$ and Z^{o} somehow get a mass, and the main problem which we wish to address in this section is understanding how that mass arises. We shall assume that the gauge-symmetry is exact; that at short distances one sees essentially the massless gluons (up to power-law corrections) of the gauge theory. So how does the mass arise? The basic notion goes back to Schwinger [41]. Writing the vacuum polarization tensor of a gauge-boson (say, W_3) as

$$\pi_{\mu\nu}(q) = (q_\mu q_\nu - g_{\mu\nu} q^2)\pi(q^2) \tag{2.1}$$

the gauge-boson propagator would be expected to be

$$D_{\mu\nu}(q) = \frac{(-g_{\mu\nu})}{q^2\,[1 - \pi(q^2)]} + q_\mu q_\nu \text{ terms and other gauge-dependent terms} \quad . \qquad (2.2)$$

If $\pi(q^2)$ has a pole at $q^2 = 0$, the pole at $q^2 = 0$ in D disappears and is replaced by a pole at some finite mass μ^2. How could such a pole in the vacuum polarization arise? An easy and instructive example is provided by the strong interactions themselves. Suppose the chiral $SU(2)_L \times SU(2)_R$ symmetry of the strong interactions were exact. (Then the gauge-invariance of $SU(2)_L$, the piece relevant to the weak interaction, would not be spoiled by the strong interactions.) In that limit the pion would still exist but according to the conventional picture would be massless. Nevertheless the massless pion would still have the coupling $gF_\pi q^\mu$ to the weak current, where F_π is the pion decay-constant. The direct pion-W coupling would give a contribution (cf. Fig. 7) to the vacuum polarization

$$\pi_{\mu\nu} = \left(\frac{g}{2}\right)^2 \frac{F_\pi{}^2 q_\mu q_\nu}{q^2} \qquad (2.3)$$

and by current-conservation would have to be supplemented by a $g_{\mu\nu}$ piece (which however does not have the pole)

$$\pi_{\mu\nu}^{(\text{pion})} = \frac{g^2 F_\pi{}^2}{4} \left(\frac{q_\mu q_\nu}{q^2} - g_{\mu\nu} \right) \quad . \qquad (2.4)$$

This $g_{\mu\nu}$ piece in turn does affect the transverse part of $D_{\mu\nu}$, leading, for small q^2, to the form

$$D_{\mu\nu} = \frac{(-g_{\mu\nu})}{q^2 - \frac{g^2 F_\pi{}^2}{4}} + \dots \quad . \qquad (2.5)$$

Hence

$$\mu_W^2 = \frac{g^2 F_\pi^2}{4} \tag{2.6}$$

and the gauge bosons W have obtained a mass! Numerically, the mass is $\sim$ 40 MeV, off by a factor $\sim 10^3$. Furthermore, the strong-interaction axial current is not by itself exactly conserved, a feature which also vitiates the argument. Nevertheless, the idea can still be used: all that is needed is to suppose there exists some other world not too dissimilar from the pion world, but on a mass-scale sufficiently large to account for an intermediate boson mass $\sim$ 70 GeV instead of $\sim$40 MeV. We shall call this world the Higgs sector. Various views of this world can exist [42], but the irreducible element is evidently that, in the limit of vanishing gauge-coupling constant, there should exist (in addition to the uncoupled Yang-Mills triplet $W^{\pm}$, Z^0) a triplet of Goldstone bosons associated with this world of Higgs degrees of freedom. We emphasize that this world is (at least in part) distinct from the world of electroweak gauge bosons, quarks, and leptons. The mass-scale associated with this world can be naturally estimated in terms of the strong-interaction world by multiplication by a factor $\gtrsim 10^3$; hence somewhere between, say, 100 GeV and 1 TeV.

It is instructive to think about the spectrum of the theory as the gauge coupling is turned off. (To be specific, we keep the mass-scale of the Higgs world fixed, and let g and g', defined at that mass scale, tend to zero.) For a nonvanishing gauge coupling, there are no massless Goldstone bosons, and three polarization-states for each massive gauge boson. As $g \to 0$, the masses of $W^{\pm}$ and Z tend to zero, and they become the transverse gauge quanta (like the photon of QED). The longitudinal polarization states are in one-to-one correspondence with the spinless Goldstone bosons. We may, in fact consider the longitudinal degree of freedom of a massive gauge boson to essentially be the corresponding Goldstone boson. There is some loose talk here, to be sure [43] , but the statement does have some truth to it. It is what is usually called "eating": the massless Goldstone bosons, when coupled to gauge quanta, are "eaten" by the gauge boson. The gauge boson becomes massive, and the Goldstone boson becomes its third polarization state.

This view leads to a remarkable qualitative conclusion. If the Goldstone particles interact strongly with each other--like, say, pions do--then so also will intermediate bosons, despite the smallness of the gauge couplings. That is, at sufficiently high energies, W - W collisions would be dominated by multiple production of W's and Z's!

It is evidently very important to delineate the dynamics which determines the properties of this Higgs world. We may gain some insight by again thinking about ordinary strong interactions. If one only knew about the low energy limit of the strong interactions, namely the pion-pion scattering-lengths (for this part of the story, we needn't worry about fermions and baryons), various field-theory Lagrangians could be written down. The most explicit such Lagrangian is the σ-model [44], and if renormalizability is demanded (most electroweak gauge theorists demand this, but it need not be a sacred principle), the model of choice is the linear σ-model. In that model, the pseudoscalar $\vec{\pi}$ field and scalar σ field are described by a hermitian 2×2 matrix

$$\Phi = \begin{pmatrix} \sigma + i\pi^0 & \sqrt{2}\pi^+ \\ \sqrt{2}\pi^- & \sigma - i\pi^0 \end{pmatrix} = (\sigma + i\vec{\tau}\cdot\vec{\pi}) \qquad (2.7)$$

with $SU(2)_L$ transformations acting on rows and $SU(2)_R$ transformations acting on columns. The gauge gluons $W^{\pm}$, W_3 are associated with the gauge-invariant substitution on the left-handed degrees of freedom

$$i\partial_\mu \Phi \rightarrow \left(i\partial_\mu - \frac{g}{2} W_\mu \right) \Phi \tag{2.8}$$

with

$$W_\mu = \begin{pmatrix} W_3 & \sqrt{2}W^+ \\ \sqrt{2}W^- & -W_3 \end{pmatrix}_\mu = \vec{\tau} \cdot \vec{W}_\mu \tag{2.9}$$

a 2 × 2 matrix acting on the left-handed degrees of freedom.

The Lagrangian for the self-interactions of the Φ-field is

$$\mathcal{L} = \frac{1}{4} \mathrm{Tr}\, \partial_\mu \Phi^+ \partial^\mu \Phi + \frac{\mu^2}{4} \mathrm{Tr}\, \Phi^+\Phi - \lambda (\mathrm{Tr}\, \Phi^+\Phi)^2 \tag{2.10}$$

which is invariant under SO(4) = SU(2) × SU(2) transformations even after introduction of the $SU(2)_L$ gauge degrees of freedom

$$\mathcal{L} = \frac{1}{4} \mathrm{Tr} \left(\partial_\mu \Phi^+ - \frac{ig}{2} \Phi^+ W_\mu \right) \left(\partial^\mu \Phi + \frac{ig}{2} W^\mu \Phi \right) - \lambda [\, \mathrm{Tr}\, (\Phi^+\Phi - F^2)]^2 \quad . \tag{2.11}$$

To this must be added the Yang-Mills $G_{\mu\nu}G^{\mu\nu}$ kinetic term for the gauge bosons. This much of the Lagrangian comprises what was called the intrinsic weak interaction in Section 1.

This way of writing the theory is not very conventional [45]. Usually a single complex spinor

$$\phi = \begin{pmatrix} \phi^o \\ \phi^- \end{pmatrix} \tag{2.12}$$

of $SU(2)_L$ is what is written down, with Lagrangian

$$\mathscr{L} = \frac{1}{2}\left(\partial_\mu \phi^+ - \frac{ig}{2}\phi^+ W_\mu\right)\left(\partial^\mu \phi + \frac{ig}{2} W^\mu \phi\right) + \frac{\mu^2}{2}\phi^+\phi - \lambda(\phi^+\phi)^2 \quad . \tag{2.13}$$

However, they are in fact equivalent. The first column of the 2 × 2 matrix

$$\sqrt{2}\phi \leftrightarrow \begin{pmatrix} \sigma + i\pi_3 \\ \pi_1 - i\pi_2 \end{pmatrix} \tag{2.14}$$

can be identified as the $SU(2)_L$ spinor. The second column

$$\begin{pmatrix} \pi_1 + i\pi_2 \\ \sigma - i\pi_3 \end{pmatrix} = \sqrt{2}\tau_1 \phi^+ \tag{2.15}$$

is, evidently, also an $SU(2)_L$ spinor, which is the "CP-conjugate" of the first complex doublet. Hence there is only one independent $SU(2)_L$ spinor participating in the Lagrangian, Eq. (2.11), and gauge invariance is then enough to imply equivalence of the two formalisms.

The extra U(1) generator of electroweak theory has not yet been included. This is most directly done in the orthodox spinor version, where one couples the generator B_μ to the hypercharge of the doublet, in accordance with the general discussion of the previous section:

$$\partial_\mu \phi \to \mathscr{D}_\mu \phi \equiv \partial_\mu \phi + \frac{ig}{2} W_\mu \phi + ig'(\tfrac{1}{2})B_\mu \phi \quad . \tag{2.16}$$

In our version, the $SU(2)_L \times SU(2)_R$ symmetry of the σ-model is explicitly broken by inclusion of the U(1) hypercharge gauge coupling; the hypercharge of the first column of Φ is (-½), while that of the second column is (+½). Thus if the gauge-invariant substitution is written

$$\partial_\mu \Phi \rightarrow \partial_\mu \Phi + \frac{ig}{2} W_\mu \Phi + ig'(-\frac{1}{2}) \Phi \tau_3 B_\mu \tag{2.17}$$

the desired effect is attained. The presence of the τ_3 implies explicit breaking; the fact that it is placed to the right of Φ implies that it is $SU(2)_R$, not $SU(2)_L$, which has been broken. As far as the Higgs sector is concerned, the U(1) factor is right-handed.

If the sign of the μ^2 term in Eq. (2.10) is appropriately chosen as implicit in our rewrite in Eq. (2.11), Φ will obtain a nonvanishing vacuum expectation value $\langle\sigma\rangle = F$ and one may then read off from the Lagrangian the mass matrix of the $W^\pm$, W_3 and B. In the absence of the g' term, the B is decoupled and the common mass of the W's is immediately read off. In the σ-model formalism one has, upon ignoring all degrees of freedom except the gauge bosons

$$\mathscr{L} \rightarrow \frac{1}{4} \mathrm{Tr}\, \mathscr{D}_\mu \Phi^+ \mathscr{D}^\mu \Phi \rightarrow \frac{g^2}{16} \mathrm{Tr}\, \Phi^+ W_\mu W^\mu \Phi = \frac{g^2F^2}{8} \mathrm{Tr}\, \frac{(\vec{\tau}\cdot\vec{W})^2}{2}$$

$$= \frac{1}{2}\frac{g^2F^2}{4} \sum_{i=1}^{3} W_\mu^{(i)} W^\mu_{(i)} \quad . \tag{2.18}$$

Hence

$$\mu^2 = \frac{g^2F^2}{4} \quad . \tag{2.19}$$

In the general case with $g' \neq 0$, short calculation [46] reproduces the results we found in the previous section, in particular Eq. (1.25) relating m_W, m_Z, and θ_W.

This digression into the algebra of the linear σ-model should not obscure the main point, which we reiterate: in order to give mass to $W^\pm$ and Z, it suffices that there exist, in the absence of gauge couplings, a Higgs sector which contains a triplet of massless, pion-like Goldstone bosons and which has a global $SU(2)_L$ symmetry, in order to preserve the observed SU(2) symmetry of the intrinsic weak interaction. The linear σ-model is one way of describing such a Higgs-world; it is indeed equivalent to the orthodox Weinberg-Salam description. However, we know from experience that this description is not unique; massless pions presumably also emerge from the QCD Lagrangian in the absence of quark mass terms. Thus the possibility exists [47] that the Goldstone bosons are composites of fermions, dubbed techniquarks or hyperquarks. These particles are bound to each other by yet another force, named technicolor or hypercolor, which is mediated by yet another set of gauge gluons. From the previous line of argument, we would again expect the natural (confinement) scale for technicolor to be $\gtrsim 10^3$ that of the strong interaction, namely ~ 500 GeV. There should be in the spectrum of this Higgs world not only the massless Goldstone pions but also techni-mesons of various spins and parities. We shall return to some more details of the phenomenology in the next section. It is, to say the least, very rich--although the energy-scale for many of the phenomena is high.

2.2. Fermion mass

Not only do gauge bosons possess mass, but also the fermions. Indeed the spectrum (Fig. 8) is provocative. It must be plotted on a logarithmic scale; in almost all cases the $SU(2)_L$ electroweak symmetry is badly broken: in almost all cases

$$\left| \frac{m_1 - m_2}{m_1 + m_2} \right| \simeq 1 \qquad (2.20)$$

where m_1 and m_2 are the masses of the members of a weak doublet.

It is, of course, an attractive idea that the mechanism responsible for gauge boson mass should also be responsible for the fermion masses. It in fact turns out to be more than a nice idea--it is close to being a necessity. The reason for this lies in the issue of renormalizability or, equivalently, of decent high-energy behavior. If one puts the fermion mass terms into the Lagrangian by hand, the gauge invariance of the theory is broken. And without the gauge-invariance there is no proof of renormalizability. In fact it is known that addition of an explicit mass term for the gauge bosons--not via the Higgs mechanism--does spoil the renormalizability. While a fermion mass-term is not the same thing, it can influence the vacuum-polarization of the gauge bosons (cf. Fig. 9). Were there no fermion mass-term present, the gauge invariance could be used (as in QED) to argue away any contribution to the gauge-boson mass, leaving only the charge renormalization contribution. But the residual piece does give a contribution. Ignoring, for simplicity, the U(1) contribution and keeping only that of the intrinsic $SU(2)_L$ part, we have

$$\delta_{ij} g_{\mu\nu} \delta\mu_W^2 \sim ig^2 \int \frac{d^4p}{(2\pi)^4} \mathrm{Tr} \frac{\tau_i}{2} \gamma_\mu \left(\frac{1-\gamma_5}{2}\right) \frac{1}{\not{p}-m} \gamma_\nu \frac{\tau_j}{2} \left(\frac{1-\gamma_5}{2}\right) \frac{1}{\not{p}-m} \quad . \quad (2.21)$$

Notice here the mass operator m is a matrix in the internal space

$$m = \bar{m} + \frac{\tau_3}{2} \Delta m \equiv \begin{pmatrix} m_1 & 0 \\ 0 & m_2 \end{pmatrix} \quad . \quad (2.22)$$

Calculation of the integral gives, to logarithmic approximation

$$\delta\mu^2 \sim \frac{\alpha_W}{8\pi} \left(m_1^2 \log \frac{\Lambda^2}{m_1^2} + m_2^2 \log \frac{\Lambda^2}{m_2^2} \right) \quad (2.23)$$

where

$$\alpha_W = \frac{g^2}{4\pi} = \frac{\alpha}{\sin^2\theta_W} \approx \frac{1}{20} \quad . \tag{2.24}$$

It is (logarithmically) divergent, implying--in the context of renormalization theory--the necessity of a mass counterterm for the gauge boson. However, according to the previous argument, this then spoils the renormalizability.

A less field-theoretical argument also leads to a similar conclusion. Consider the amplitude for $e^+e^- \to W^+W^-$, with both W's polarized longitudinally (in the center of mass system) [48]. Again, take the SU(2) intrinsic interaction alone, just to keep algebra simple. Then there are only two diagrams (Fig. 10). The longitudinal polarization-vector ε_L^μ for a W-boson of four-momentum q^μ is

$$\varepsilon^\mu = \frac{1}{m_W}(|\vec{q}\,|, E_1, 0, 0) = \frac{q_\mu}{m_W} + \frac{m_W}{E_1 + |\vec{q}\,|}(1, 1, 0, 0) \quad . \tag{2.25}$$

Because of the factor m_W in the denominator of the first term it is possible to have growth in the amplitude at high energy if it is not compensated. As we mentioned before, gauge couplings provide the cancellation, provided gauge invariance is maintained and provided the Higgs-mechanism is invoked. Let us watch this work in the above example. Keeping only terms of order m_W^{-1} or m_W^{-2} as $s \to \infty$, we get

$$\mathcal{M}_{LL} = \frac{1}{2m_W^2}\bar{v}(p_+)\not{q}_+\left(\frac{1-\gamma_5}{2}\right)\frac{1}{\not{p}_- - \not{q}_-}\not{q}_-\left(\frac{1-\gamma_5}{2}\right)u(p_-)$$

$$- \frac{g^2}{2m_W^2}\frac{\bar{v}(p_+)\gamma_\mu\left(\frac{1-\gamma_5}{2}\right)u(p_-)}{(q_+ + q_-)^2}\Big[(q^+ - q^-)_\mu q^+\cdot q^- +$$

$$+ q_\mu^-(2q_+\cdot q_-) - q_\mu^+(2q_+\cdot q_-)\Big] + O(1) \quad . \tag{2.26}$$

Reduction of the algebra gives

$$\mathcal{M}_{LL} = \frac{g^2 m_e}{4 m_W^2} \bar{v}(p_+) u(p_-) + O\left(\frac{m_e^2}{m_W^2} \right) + O(1) \qquad . \qquad (2.27)$$

The first term has a dimensional coefficient $\sim(\text{mass})^{-1}$; hence the high energy behavior is that of a non-renormalizable coupling. Roughly, the trouble sets in when

$$\sqrt{s} \gg \frac{8}{\alpha_W} \frac{m_W^2}{m_e} \sim 10^9 \text{ GeV} \qquad . \qquad (2.28)$$

However for $\tau^+ + \tau^- \to W^+ + W^-$ it is quite a bit lower; $\sqrt{s} \gg 250$ TeV. For $t + \bar{t} \to W^+ + W^-$, the limit (for $m_t \sim 30$ GeV) is still lower, $\lesssim 15$ TeV.

Thus the presence of fermion mass can interfere with the high-energy, short distance structure of the theory unless something is done about it. Again we may compare the orthodox σ-model approach to that of technicolor.

The solution within the orthodoxy is quite direct. Given the elementary Higgs fields Φ introduced to give $W^{\pm}$ and Z their mass, it is easy and natural to introduce gauge-invariant trilinear Yukawa couplings of Φ to the fermions. It is the most general renormalizable thing that one can do with this raw material. In the σ-model notation we may write for this coupling

$$\mathcal{L}' = \bar{L}\Phi h R + \text{h.c.} \qquad . \qquad (2.29)$$

The notation here is very compact. Schematically L are the left-handed fermion fields, h is a coupling-constant matrix and R are the right-handed fermion fields. Upon replacement of the matrix Φ of higgs fields by their vacuum-expectation value $F = 2m_W/g$ (times the unit matrix) one gets a mass term coupling the various fermions to each other, with the mass-matrix

$$m = Fh = \frac{2m_W}{g}h \tag{2.30}$$

proportional to the input coupling-constant matrix.

We have been overly schematic in the preceding paragraph, and it is necessary to include all the indices to understand better what the implications are. Write

$$L_i^{\alpha} \quad , \quad R_i^{\alpha} \tag{2.31}$$

for left-handed and right-handed fermion fields, where $i = 1,2$ indicates weak-isospin up and down, and where $\alpha = 1,2,\ldots 6$ $(\ldots?)$ labels the type of fermion (i.e. $L_1^{\ 1}, L_1^{\ 2}, \ldots L_1^{\ 6} \equiv \nu_e, \nu_\mu, \nu_\tau, u_L, c_L, t_L$ and $L_2^{\ 1}, L_2^{\ 2}, \ldots L_2^{\ 6} \equiv e_L, \mu_L, \tau_L, d_L, s_L, b_L$). The coupling term is then

$$\mathscr{L}' = \bar{L}_i^{\alpha}\Phi_{ij}h_{jk}^{\alpha\beta} R_k^{\beta} + \text{h.c.} \quad . \tag{2.32}$$

The j and k indices refer to the $SU(2)_R$ symmetry, which therefore is intrinsically broken by this coupling. The $SU(2)_L$ symmetry manifestly is left intact. The weak-hypercharge U(1) symmetry will be preserved provided charge conservation and T_{3R}-conservation are respected; this is assured provided there are no couplings of leptons to quarks and provided the h matrix is diagonal in the $SU(2)_R$ indices:

$$h_{jk}^{\alpha\beta} = h_j^{\alpha\beta}\delta_{jk} \quad . \tag{2.33}$$

Upon replacing the Higgs field with its expectation value

$$\Phi_{ij} \rightarrow \langle\Phi_{ij}\rangle \equiv F\delta_{ij} \tag{2.34}$$

we get the mass term

$$\mathscr{L}' = \sum_{i=1}^{2} F\bar{L}_i^{\alpha} h_i^{\alpha\beta} R_i^{\beta} + \text{h.c.} \tag{2.35}$$

where the 6 × 6 matrices $\underset{\sim}{h}_1$ and $\underset{\sim}{h}_2$ each break down into at most 3 × 3 blocks. If there do not exist right-handed neutrinos, we may set the matrix $h_1^{\alpha\beta}$ ($\alpha,\beta = 1,2,3$) to zero, thereby decoupling those degrees of freedom from the theory (they have vanishing weak hypercharge and isospin; hence do not couple to gauge bosons either).

The resultant mass matrix can be diagonalized by redefining the fermion fields; this in turn leads to the Kobayashi-Maskawa mixing matrix in the weak current. The highly nontrivial structure of that matrix implies highly nontrivial structure in the h_i as well--including quite possibly CP-violating phase factors. In the orthodoxy, all these parameters are put in by hand; one gets out no more or no less than what is put in. Because the $h_i^{\alpha\beta}$ are proportional to the masses of the fermions, the elements of the 3 × 3 matrices are of widely differing magnitude. It is not at all understood why this should be the case.

With this digression we can return to our examples of apparently catastrophic fermion mass-induced effects and see how they are resolved within the orthodox picture. First of all, the divergent mass-insertions into the gauge boson propagator are now (c.f. Fig. 11a) Higgs-induced. The loop is still divergent. Nevertheless it is related, via gauge invariance (recall Eqs. (2.1) and (2.4)) to the diagram of Fig. 11b, and this is just a renormalization of the self-interaction of the Higgs field, and is removed by a counterterm.

In the second example, we have evidently one more diagram to add. The scalar σ degree of freedom not only has a vacuum-expectation value, but will also have a dynamical piece. [This σ is what is generally called the Higgs boson.] It behaves, just like $\langle\sigma\rangle$, as a scalar field with vacuum quantum numbers. The coupling to W arises from the mass term (c.f. Fig. 12) leading to a coupling $gm_W \sigma \vec{W}\cdot\vec{W}$. The coupling to the fermion is evidently

$$h_i^{\alpha\beta}\bar{f}_i^{\alpha} f_i^{\beta}\sigma = \frac{m_i^{\alpha\beta}}{F}\bar{f}_i^{\alpha} f_i^{\beta}\sigma = \frac{gm_f}{2m_W}\bar{f}f\sigma \quad . \tag{2.36}$$

Thus in the process $e^+e^- \to W^+W^-$, the amplitude with s-channel Higgs exchange is

$$\mathcal{M}_{LL} = -\frac{m_e}{F}\bar{v}(p_+)u(p_-)\frac{1}{(q_+ + q_-)^2 - m_\sigma^2}\left(\frac{g^2F}{2}\right)\epsilon_L\cdot\epsilon_L' \underset{s\to\infty}{\longrightarrow} -\frac{g^2m_e}{4m_W^2}\bar{v}(p_+)u(p_-) + O(1) \quad . \tag{2.37}$$

For $s \gg m^2$, this suffices to cancel off the offending term, Eq. (2.27). Inasmuch as $m_\sigma^2 \lesssim 1\ \mathrm{TeV}^2$, this is a satisfactory situation, given that there are no superheavy fermions with weak isospin.

In fact, there is a line of argument against the existence of more flavors of very heavy fermions coupled to W and Z, at least for weak doublets with mass differences comparable to the average mass [49]. This comes from the finite vacuum polarization effect, due to virtual fermion-antifermion pairs. This splits m_W from m_Z--even disregarding the mixing of the U(1) generator with W_3. This splitting, if too large, upsets the strength of the intrinsic neutral-current interaction which is in experimental accord with weak-isospin conservation. The splitting can be directly calculated from Eq. (2.21); the result is

$$\left(m_Z^2 - m_W^2\right)_{\sin^2\theta_W \to 0} = -\frac{\alpha_W}{16\pi}\left\{ m_1^2 + m_2^2 - \frac{2m_1^2 m_2^2}{m_1^2 - m_2^2}\log\frac{m_1^2}{m_2^2}\right\} \quad . \tag{2.38}$$

The SU(2) symmetry of the intrinsic weak interaction is good in amplitudes to $\lesssim 3\%$. Hence for a single extra sequential lepton with an accompanying light neutrino we must have

$$m_l \lesssim \left(\frac{16\pi}{\alpha_W} \times .03\right)^{1/2} m_W \sim 400 \text{ GeV} \quad . \tag{2.39}$$

For a quark pair the limit is even better: $\lesssim$ 230 GeV.

Thus far we have considered the machinery of fermion mass generation within the orthodox "σ-model" picture. In a word, it is almost a corollary of the mechanism used for generating gauge boson mass; no new degrees of freedom had to be introduced. The price paid is that the pattern of Yukawa couplings is complicated and is put into the Lagrangian by hand.

In the technicolor-hypercolor models, the situation is quite different [50]. Again, because we must break the $SU(2)_L$ symmetry, it is attractive to use the same degrees of freedom used to generate the gauge boson masses. However, we now are obliged to couple the technifermions to the ordinary fermions (c.f. Fig. 13). To do this appears to require yet another interaction to effect the transition from ordinary fermions to the technifermions. The technifermions, like quarks at the GeV level, generate a dynamical mass at the TeV level, as shown in Fig. 14. However, at short distances (small compared to 1 GeV^{-1}) this dynamical mass tends to zero rapidly because the Lagrangian, "bare" mass of the technifermions vanishes. Hence the mass insertion should have a dependence on how virtual the techniquark is.

Now go back to Fig. 13 and suppose the four-fermion transition $f\bar{f} \leftrightarrow F\bar{F}$ is effectively local at the sub-TeV level. Then we get, schematically, a fermion mass of order

$$m_f \sim G\langle\bar{\psi}\psi\rangle \sim G\langle\bar{m}\rangle^3 \quad . \tag{2.40}$$

We may guess that $\bar{m}$ is anywhere from 100 GeV to 1 TeV. If we take it to be 300 GeV, we are within a factor 30 of the extreme estimates for m_f. We get, in GeV-units

$$m_f \sim 3 \times 10^{7 \pm 1.5} G \quad . \tag{2.41}$$

Because m_f varies from 5×10^{-4} GeV to > 15 GeV, it is hard to know what to assume for G. In any case it is not much larger than G_F and could be orders of magnitude smaller. Notice that because this interaction should account for the mass matrix, there should at some level be new flavor-changing phenomena as well. However, at this point it is clear that better guiding principles are needed [51].

A natural and attractive starting point is again a gauge principle. The application here is to suppose that the new effective four-fermi interaction is mediated by new very heavy gauge bosons. This is called extended technicolor. But even here there is not very much discipline in building up the theoretical structure. And at present there is no good candidate model which looks realistic. Nevertheless there is much more to this approach than what is apparent from working in the direction from phenomenology toward theory. It is necessary really to start in the other direction, from first principles to fully appreciate this approach. This will be done in a later section.

3. Miscellaneous Experimental Questions

The preceding considerations give on the one hand a rather definite menu of experiments and phenomena associated with the standard $SU(2) \times U(1)$, and on the other a definite suggestion of incompleteness--of the existence of a world of phenomena beyond the standard model. However, the latter, unlike the former, is not at all sharply focussed and therefore requires a very open-minded and broad-based search philosophy. Without going into detail here, we now catalogue a few of the experimental implications of the preceding section.

3.1. Search for $W^{\pm}$ and Z^{o}

The next major step in the evolution of electroweak theory will clearly be the testing of the Weinberg-Salam predictions for the masses and properties of the intermediate bosons $W^{\pm}$ and Z^{o}. Overt production evidently requires high energies; energies accessible by the forthcoming generations of colliding beam machines. The possibilities include

(i) $p\bar{p} \to W^{\pm}$ or Z^{o} + anything: Assuming the Weinberg-Salam model parameters together with the reasonably successful quark-parton Drell-Yan model for its production, the CERN $p\bar{p}$ collider should have a good chance of finding the Z^{o} via its decay into $\mu^{+}\mu^{-}$ and $e^{+}e^{-}$. This process has been extensively studied [52] and will not be reviewed here. The $W^{\pm}$ decay into lepton $+\nu$ is more difficult than $Z^{o} \to \ell^{+}\ell^{-}$ and may require a higher luminosity collider such as the pp machine ISABELLE, now under construction at Brookhaven.

(ii) $e^{+}e^{-} \to Z^{o}$, resonantly: The LEP machine, again at CERN, should produce at least 10 Z^{o} per minute, allowing not only precise determination of its properties but also a rich source of all its decay products. The physics is so clean, well-defined, and predictable that there is already little to do but wait [53], or else to contemplate alternative theories of neutral currents.

(iii) $e^+e^- \to W^+W^-$: Ultimately this process can be observed at LEP and is important in explicitly testing the gauge-structure of the theory. It may not be too easy. For example, if the orthodox σ-model Higgs picture is correct, but the Higgs boson is too heavy to detect ($m_H > 100$ GeV), we might hope that this would manifest itself in this process. However, the calculated effects are small. I am not sure, but suspect this may be due in part to the fact that W_L is essentially the same as the Higgs, and production of spin zero pairs by e^+e^- is small. Furthermore, Higgs effects can only occur via W - W rescattering. Near threshold, these phase shifts are all small, in large part due to Adler-like low-energy theorems [54]. However, if W's are composite and/or interact strongly with each other for reasons outside gauge theory ideology, the $e^+e^- \to W^+W^-$ process should be especially revealing.

(iv) $e^-p \to \begin{matrix} e^- + \ldots \\ \nu_e + \ldots \end{matrix}$. While these are only virtual W and Z exchange processes, propagator effects, revealing the finite range of the weak force, are prominent for proposed colliders (10 - 30 GeV electrons against 400-1000 GeV protons). This would be very important were $W^\pm$ and Z^0 not found by the other methods. In particular, the general arguments of the previous section using essentially only the structure of the Weinberg-Salam low-energy effective Lagrangian require these propagator effects be observable in the charged-current reactions (at least a factor two effect in $d\sigma/dQ^2$ for $Q^2 > 10^4$ GeV^2).

3.2. Search for the Standard Higgs

The standard ("σ-model") Higgs boson of the Weinberg-Salam model is probably best searched for in e^+e^- annihilation:

If $m_h \lesssim 10$ GeV, the decay $\Upsilon, \Upsilon', \Upsilon'' \to h + \gamma$ may be observable [55]. This mass is what is obtained from a calculation of lowest order $W^\pm$, Z, radiative corrections. Perhaps this contribution dominates--although it is hard to justify why. If $m_h \lesssim 40$ GeV, the decay $Z^o \to \begin{cases} h + \mu^+\mu^- \\ h + e^+e^- \end{cases}$ is observable [56].

If $m_h \lesssim 100$ GeV, the Higgs boson may be able to be observed at the ultimate LEP (E = 140 GeV + 140 GeV) via $e^+e^- \to Z^o + h$.

If $m_h > 100$ GeV, the situation simply looks very difficult. Probably even higher energy e^+e^- colliders are the best hope.

3.3. Search for other states

i) Gauge-bosons: The electroweak group may be bigger than SU(2) × U(1), even at an accessible energy scale. In any enlarged group, there will be extra Z^o's. As we already discussed, there probably would be an indirect effect on the ordinary Z^o via mixing, driving it to lower mass. Another natural alternative would be an extra $SU(2)_R$ (or larger group), as especially suggested by the asymmetry in the σ-model description of the Higgs mechanism. If there does exist left-right symmetry which is dynamically broken, then we could expect right-handed neutrinos coupling to these right-handed currents and all the concomitant phenomenology of neutrino mass and mixing [3]. Present lower limits on masses of W_R or $Z_R{}^o$ are typically $> 3\, m_W$. Proposed $p\bar{p}$ or pp colliders in the multi-TeV range may produce such bosons at a measurable rate if luminosities of $\gtrsim 10^{31} cm^{-2} sec^{-1}$ can be attained.

ii) Extra quarks and charged leptons: The methods are well known and need no elaboration; e^+e^- machines are very good here. However as the mass gets into the 50-100 GeV region, quarks may also be observable in $p\bar{p}$ colliders, even with modest luminosity ($\gtrsim 10^{30} cm^{-2} sec^{-1}$). For example, even for the t-quark the process

$$\begin{array}{l} p\bar{p} \to t\bar{t} + \ldots \\ \qquad\quad \bar{t} \to \bar{b}\ \text{jet} + \ell + \nu \\ \qquad\quad t \to 3\ \text{jets} \end{array} \tag{3.1}$$

may be reconstructable if $m_t > 30 - 50$ GeV. As the mass increases, the method should become easier.

(iii) Extra Higgs bosons: Almost any extension of the standard model leads to extra Higgs bosons, some of which may not be too massive [57]. Such extensions must take care not to upset limits on flavor-changing processes; conversely this means one should also search for Higgs-exchange effects in rare decays.

If extra Higgs bosons exist, some may be charged. These can be produced in e^+e^- collisions. However, with $\Delta R = \frac{1}{4}$ and a $(1 - 4m^2/s^2)^{3/2}$ phase-space factor near threshold, it will not be a prominent signal. Technicolor models, discussed separately in Section 4, predict existence of pseudoscalar mesons (Pseudogoldstone bosons) with properties quite similar (but not identical) to such Higgs bosons. Some of these are expected to be charged and light enough to be produced in the present PEP/PETRA energy regime.

The decay products of these objects are uncertain; but the best guess is that these bosons decay mainly into the most massive fermion pair kinematically accessible consistent with charge-conservation.

iv) Neutral leptons: We should not forget that $e^+e^- \to L^o\bar{L}^o$ via Z^o exchange is a nontrivial channel even at PEP/PETRA energies; a neutral member of an electroweak doublet is produced at a rate of a few percent of the $\mu^+\mu^-$ rate. Furthermore, this ratio grows as the fourth power of the center-of-mass energy. Multilepton final states with high sphericity provide a reasonable signature.

3.4. Search for rare processes

Almost any idea that complicates the theory beyond the standard model produces new bosons (Higgs or gauge) which mediate new four-fermi interactions. Therefore, as mentioned above, the study of decay amplitudes at the sub-weak level deserves to be pursued in as many directions as practical. This frontier may yet be as exciting a one as the push to higher energies.

3.5. Search for quark or lepton compositeness

All colliding-beam techniques can evidently be used here. If at the ultimate LEP, as many $\mu\mu$, $\tau\tau$, $q\bar{q}$ events were to be accumulated as at PETRA, with agreement with theory, then the limits on gross internal structure of leptons and quarks could be pushed to $\Lambda \gtrsim 1$ TeV. This might be difficult to improve upon with the presently proposed ep colliders. If hard qq scattering at the $p\bar{p}$ or pp colliders were well understood in terms of QCD and parton-model concepts, one also might be tempted to use that process to bound any quark structure. A 50% cross-section measurement at $p_\perp \gtrsim 300$ GeV gives a limit on Λ of $\gtrsim 1$ TeV. For reasons of rate it may be difficult to improve this limit--even at much higher energy, without observing directly the breakup of the quark into its constituents [58].

4. Technicolor Phenomenology

It is not our purpose to develop in this section a comprehensive theoretical overview of technicolor (hypercolor) theories, (including the complications of extended technicolor) but only some phenomenological implications. We shall concentrate on those which are least dependent on the details of extended technicolor, and follow the work of Dimopoulos, Kane and Raby [59] . We recall from Section 2 that the essential ingredients are

i) A new set of technifermions which have no bare mass, which transform nontrivially under the electroweak group $SU(2) \times U(1)$ as well as transforming nontrivially under another nonabelian technicolor group.

ii) A set of gauge-technigluons which confine techniquarks into technicolor-singlet bound states in a manner similar to QCD--with the exception that the confinement-scale is now ~ 1 TeV instead of ~ 1 GeV.

iii) Given N such technifermions, then the technicolor Lagrangian, in the absence of other interactions (electroweak, QCD, extended technicolor, ...), will have an $SU(N)_L \times SU(N)_R$ chiral symmetry which is assumed to spontaneously break. This leads to a large number of O^- Goldstone bosons. When the other interactions are included, three (and only three) should survive as massless particles; these will be eaten by $W^{\pm}$ and Z^0.

However, the remainder, if any, should survive as almost massless (pseudo-Goldstone) bosons. For moderately low energy applications ($E_{CMS} \lesssim 1 - 10$ TeV), these seem to be the most important for phenomenology.

In the minimal example given in Section 2, there were no such extra pseudo-Goldstone states. However, there seemed to be little hope in accounting for fermion mass at that state, and "realistic" attempts using more degrees of freedom invariably lead to some such particles.

Hereafter we shall, for definiteness, assume:

i) One extra flavor-family of technifermions U, D, E, N.

ii) An SU(N) technicolor group (N = 3 or 4).

iii) The technifermions in the N-dimensional representation

and only sketch the phenomenological consequences for the 63 Goldstone or pseudo-Goldstone states. The remaining particles (technivector mesons, technibaryons,...) are expected to have masses $\gtrsim 1$ TeV, although there is clearly a lot of uncertainty in that estimate.

The 64 candidate O^- states are

1) Techni- η': This is an SU(8) singlet which couples to an anomaly-ridden, nonconserved axial current. It presumably has a very big mass (500 - 1000 GeV).

2) 32 color-octet technipions and techni- η's": These are

$$\bar{U}\gamma_5\lambda_i D \quad \bar{D}\gamma_5\lambda_i U$$

$$\bar{U}\gamma_5\lambda_i U \quad \bar{D}\gamma_5\lambda_i D$$

They get a (reasonably calculable) mass $\backsim$250-300 GeV via gluon emission and reabsorption. The calculation parallels the PCAC calculation of the pion electromagnetic mass.

3) <u>24 color-triplet technileptoquarks:</u> These are

$$\bar{E}\gamma_5 U \quad \bar{E}\gamma_5 D$$

$$\bar{N}\gamma_5 U \quad \bar{N}\gamma_5 D$$

and their antiparticles. Their mass is 2/3 that of the color octet bosons.

4) <u>3 true Goldstone bosons eaten by $W^{\pm}$ and Z^o:</u> In the limit $\sin^2\theta_W = 0$, when the usual U(1)-generator complications can be thrown away, these are

$$\bar{U}\gamma_5 D + \bar{N}\gamma_5 E$$

$$\bar{D}\gamma_5 U + \bar{E}\gamma_5 N$$

$$(\bar{U}\gamma_5 U - \bar{D}\gamma_5 D) + (\bar{N}\gamma_5 N - \bar{E}\gamma_5 E)$$

5) <u>4 color-singlet pseudo-Goldstone bosons, including a charged pair:</u> These are (for $\sin^2\theta_W = 0$)

$$\bar{U}\gamma_5 D - \bar{N}\gamma_5 E$$

$$\bar{D}\gamma_5 U - \bar{E}\gamma_5 N$$

$$(\bar{U}\gamma_5 U - \bar{D}\gamma_5 D) - (\bar{N}\gamma_5 N - \bar{E}\gamma_5 E)$$

$$(\bar{U}\gamma_5 U + \bar{D}\gamma_5 D) - 3(\bar{N}\gamma_5 N + \bar{E}\gamma_5 E) \quad .$$

The charged boson gets a mass of ~ 8 GeV via photon, Z, and $W^{\pm}$ emission and reabsorption. The remaining two do not. Extended technicolor must be invoked to give these (dangerous?) axion-like objects masses. Dimopoulos, Raby, and Kane, using Pati-Salam leptoquarks, get a small contribution ($\sim 2 - 3$ GeV) from this source. However, this part of that estimate is very uncertain. But more open-minded guesses [60] still leave these masses in the 10 - 50 GeV range. Of course, the same mystery mechanism which provides mass to the neutral bosons will in all likelihood contribute to the charged pseudo-Goldstone bosons as well. In any case, a careful search for such particles should be carried out at the e^+e^- colliding-beam machines.

The low-mass color-singlet particles appear to be the best candidates for experimental search. It may not be an inevitability, for example, that colored technifermions are even necessary. However, there do exist opportunities for observation of colored technibosons. The best candidate appears to be single production of the techni-eta in pp or $p\bar{p}$ colliding beams [59,61]. The techni-eta couples to a pair of QCD gluons, much as the ordinary η to two photons, via the celebrated triangle diagram, with the technifermions running around the loop. There is not much uncertainty in the calculation of the width into two gluons (it turns out to be $\sim$50 MeV) and hence not much uncertainty in its production in hadron-hadron (i.e. g - g resonant) collisions. At Fermilab Tevatron energies the cross-section is a quite respectable $2 \times 10^{-34}\text{cm}^2$. The problem is in its observation. The dominant channel is two-body: either gg or $t\bar{t}$. Dimopoulos, Raby, and Kane believe it to be $t\bar{t}$, but this depends upon the vagaries of extended technicolor. In either case there is expected to be a rather big background from QCD $q\bar{q} \to gg$ and $q\bar{q} \to t\bar{t}$ subprocesses.

Pair production of the colored technihadrons--whether color triplet or color octet--can proceed via $q\bar{q}$ annihilation via single-gluon exchange. Kinematically, there is enough energy at the Tevatron to produce such a pair (mass $\gtrsim 350 + 350$ GeV), but the cross sections are small (10^{-36} - 10^{-37}cm^2) for a variety of reasons (unfavorable color factor, spin zero, β^3 threshold factor, falling parton distributions). However if the nominal mass-scale for vector technimesons is a bit lower than 1 TeV (and this mass scale is in fact quite uncertain) the cross section could be quite a bit bigger. This occurs because there exists a 1^- color-octet techni-ω which has quantum numbers identical to the gluon. Thus there can be resonant production of this techni-ω in q - $\bar{q}$ collisions analogous to resonant production of ω in e^+e^- collisions. The cross section depends very sensitively on the mass of this techni-ω: were it 500 GeV, it would be produced at an observable level at the Fermilab 1 + 1 TeV $p\bar{p}$ collider. But with a 1 TeV mass, the cross section is probably too small.

The decay products of the techni-ω are spectacular, and include pairs of colored technipions (charged or neutral) or colored techni-etas. Technileptoquarks will also be present in the decay products. The charged technipions decay into $W^{\pm}$ + g (or possibly $t\bar{b}$) and the neutral technipion and techni-η into two gluons (or possibly $t\bar{t}$). A technileptoquark decays into a lepton and a quark, although here the details depend very much upon how extended technicolor is implemented. It is clear that were the techni-ω to exist and were it able to be produced in $p\bar{p}$ colliding beams it would be an extremely rich source of information on the Higgs-sector in its natural energy domain. Nevertheless, we must again stress that even within the technicolor picture, the existence of colored technifermions involves an additional assumption. The most reliable search method for now is probably to look for charged scalar bosons in e^+e^- colliding beams.

5. Technicolor Ideology

5.1. "Tumbling"

Why is the technicolor idea an attractive one? Thus far, the difficulties with the phenomenology of extended technicolor make it seem a rather extravagant and apparently awkward way of circumventing problems associated with a Higgs sector built from elementary scalar fields. However, the basic conceptual structure of the idea is in principle economical and beautiful, and worth briefly describing here [62].

One envisages at some very high energy--perhaps 10^{15} - 10^{19} GeV, perhaps lower--beginning with some large non-abelian gauge group G with a set of left-handed fermions which transform irreducibly under G. If there are N such fermions, there will be an enormous ungauged SU(N) chiral symmetry in the theory. At this energy the running coupling-constant (or constants) associated with G are assumed small. However at energies lower than this initial energy scale the coupling constant (or constants) will become larger, as a consequence of the asymptotic freedom property of nonabelian gauge theories. At some critical energy Λ_1, it can be expected--and is <u>assumed</u>--that some spontaneous chiral symmetry breakdown will occur. Thus some of the fermions will pair off and behave at lower energies as massive Dirac fermions (but perhaps confined like quarks) and others will remain massless, associated with a residual chiral-symmetry. Also, some of the gauge bosons will acquire mass a la Higgs, and therefore the original gauge-group G will be broken down to $G_1 \subset G$ at energies small compared to the critical energy Λ_1.

As the energy is decreased below Λ_1, the residual gauge coupling constants will again grow, triggering another such process. Dimopoulos, Raby and Susskind dub this phenomenon "tumbling." After some number n of tumbles, the original group G is then envisioned to have been broken down to $G_n = G_{TC} \times SU(3)_{QCD} \times (SU(2) \times U(1))_{WS}$ and the final tumble of G_{TC}--as we described previously--is to give mass to $W^{\pm}$ and Z^o, and leave only $SU(3)_{QCD} \times U(1)_{QED}$ unbroken. Meanwhile, the original gigantic chiral symmetry breaks down along with the breakdown of G, with most of the original fermion degrees of freedom pairing up into massive degrees of freedom, leaving behind only the observed spectrum of quarks and leptons.

An attractive feature of this scenario is that it can naturally account for disparate mass scales. Because the running coupling constants vary only logarithmically with energy, it is probable that the critical mass scales where tumbling occurs are widely spaced: $\Lambda_1 \gg \Lambda_2 \gg \ldots \Lambda_n$; hence one may also be able to account for the hierarchy of fermion masses.

However, the most attractive feature is its in principle absence of parameters: the input parameter (or parameters?) is only the initial coupling constant (or constants?) for the original gauge group G. The only other parameters are discrete--namely, which group G and which fermion representation of G is chosen. In return for this conceptual simplicity there is the difficult, non-perturbative dynamical problem of the "tumbling" mechanism to attack.

Dimopoulos, Raby, and Susskind have suggested a very plausible starting point--namely that examination of single-gluon emission and absorption suffices to determine the group structure of the spontaneous breakdown mechanism. The mechanism is illustrated in Fig. 15. In a given stage of tumbling, some set of fermion pairs of spin zero

$$(\sigma_2)_{\alpha\beta} f_{L_\alpha}^{\ i} f_{L_\beta}^{\ j} = \phi^{ij} \tag{5.1}$$

forms a "condensate"; i.e. it is assumed that $\langle 0 | \phi^{ij} | 0 \rangle \neq 0$. The ϕ^{ij} can be decomposed into channels A which transform irreducibly under the unbroken (but about to be spontaneously broken) group. Dimopoulos, Raby and Susskind suggest that one simply look in each of these channels A at the strength of the fermion-fermion force coming from single-gluon exchange. They then assume, quite plausibly, that the channel A with the strongest attractive force (MAC, or "most attractive channel") is the one which undergoes the spontaneous breakdown. This assumption is quite powerful, and allows the analysis of large clases of models, using group theory alone. We refer the interested reader to the literature for examples. A reasonably realistic example is given by Dimopoulos, Raby, and Sikivie [63], along with a useful critique and list of the problems that must be overcome in a completely realistic model. These include:

1) No triangle anomalies in the gauged current associated with the group G.

2) A satisfactory "tumbling" sequence.

3) SU(2) symmetry for the low-energy intrinsic weak interaction (nontrivial, because of the SU(2) breaking effects which occur at the extended-technicolor stage).

4) Baryon conservation. Leptoquark gauge bosons are needed to ensure breaking of a chiral symmetry associated with the overall relative phase of quark and lepton degrees of freedom. (If that symmetry is not broken, an unacceptable axion appears in the spectrum.) Instanton effects (which are potentially strong) also must be watched with care.

5) Strong CP violation (induced by instantons) must be avoided.

6) Axions can occur, and any axions which do must be compatible with experimental bounds.

7) Fermion masses and Cabibbo angles must be correctly generated. It seems to be especially difficult (but in principle not impossible) to account for the Cabibbo mixings.

8) The final unbroken gauge group should be $SU(3)_{QCD} \times U(1)_{QED}$ with no additional unbroken U(1) generators (extra photons).

As the authors remark, the trick is not only in showing each such condition can be met, but also to show that they can all be met simultaneously.

5.2. "Complementarity"

Another aspect of the technicolor ideology has recently emerged: it is the possibility that the residual fermions of low mass that we see can be regarded as composites of other basic fields. If this possibility is realized, then it most likely should be associated with some kind of chiral symmetry. There is plenty of chiral symmetry in the previous considerations; what needs to be shown is how residual massless fermions can emerge from a "tumbling" scheme, and in what sense they are composite. Again no realistic model is available, but Dimopoulos, Raby and Susskind [64] have provided a nice example of the ideas involved.

The starting gauge group is SU(5) and the starting fermion representation is a $5 + \overline{10}$ just as in the grand unified models. We represent

$$5 = \psi_i$$

$$\overline{10} = \chi_{ijk} \quad \text{i, j, k antisymmetric} \quad . \tag{5.2}$$

However, the starting coupling is taken as somewhat larger than in the grand unified models, so that spontaneous breakdown occurs. According to the previous rules, the MAC is a 5; the $\overline{10}$ attracts itself (via SU(5) gluon exchange). The nonvanishing vacuum expectation value of the 5 can be rotated into the 5th direction; hence the residual gauge group is SU(4), with

$$\overline{10} = \text{massive 6 formed from } \chi_{5AB} + \text{massless } \overline{4} = \chi_{ABC} \quad (A,B,C, = 1,2,3,4) \quad .$$

At somewhat lower energies, when the SU(4) coupling evolves to a large enough magnitude, a secondary spontaneous breakdown occurs, with the 4 (from the $5 = \psi_i$) and the $\overline{4}$ (from the $\overline{10} = \chi_{ijk}$) forming the MAC. This breaks down the gauge symmetry completely, all SU(4) bosons acquiring a mass at this stage. The only massless object remaining is the ψ_5 degree of freedom.

Thus far we have only exhibited an example of how the MAC game is played. However, now a new idea enters. We have considered the original SU(5) theory as a spontaneously broken theory--*a la* Higgs. But 't Hooft has emphasized [65] that in a real sense the SU(5) theory can be regarded as *unbroken*--that owing to the fact that the ψ_i and χ_{ijk} are coupled to gauge fields, the vacuum state remains gauge-invariant and unique. We may equally well regard the theory--and analyze it--as an unbroken, confining theory like QCD [66]. If that is the case, the physical particles must be SU(5)-singlet composites of the original fields. This is nicely described by 't Hooft for the Weinberg-Salam model itself. There the SU(2) invariant composites for the fermions are objects like $\psi_i\phi_j\varepsilon^{ij}$ or $\psi_i(\phi^+)^i$. (Notice that with $\phi_i \to <\phi_i>$, this projects out the individual components of the fermion fields). For the gauge-boson fields, the composites are $\phi_i^+\overleftrightarrow{\partial}_\mu\phi_i$, or $\phi_i\partial_\mu\phi_j\varepsilon^{ij}$.

If our SU(5) example is regarded from this point of view, we may ask where the massless fermion occurs. The basic SU(5)-singlet composite-fermion fields which can be constructed are

$$\Psi = \psi_i \chi^{\dagger\, jk\ell} \chi^{\dagger\, mni} \varepsilon_{jk\ell mn}$$

$$\Phi_1 = \psi_i\psi_j\chi_{\ell mn}\varepsilon^{ij\ell mn}$$

$$\Phi_2 = \chi_{ipv}\chi_{jqw}\chi_{krx}\chi_{msy}\chi_{ntz}\varepsilon^{ijkmn}\varepsilon^{pqrst}\varepsilon^{vwxyz}$$

$$\Phi_3 = \psi_i\psi_j\psi_k\psi_m\psi_n\varepsilon^{ijkmn} \quad . \tag{5.3}$$

Is there a candidate among these fermions for the massless particle?? Since $\chi^\dagger\chi^\dagger$ possessed a vacuum expectation value in the Higgs picture, the most natural candidate is the Ψ-fermion, inasmuch as $\psi_5 \propto \psi<\chi^\dagger\chi^\dagger>$. Is there a way of checking? The answer is yes; one may try using a consistency criterion proposed by 't Hooft [65]. There exists in the unbroken SU(5) theory one global chiral symmetry

$$\psi_i \to e^{3i\theta}\psi_i$$

$$\chi_{jk\ell} \to e^{-i\theta}\chi_{jk\ell} \tag{5.4}$$

with a conserved charge

$$Q = 3N_\psi - N_\chi \tag{5.5}$$

which is not an SU(5) generator. According to the 't Hooft argument, if massless composite fermions occur in the physical spectrum, then the triangle anomaly for the composite fermion which is built from the aforementioned current should precisely equal the triangle anomaly computed from the original elementary fields. We may do the calculation both ways. For the single presumed composite massless field Ψ, a chiral rotation takes Ψ to $e^{5i\theta}\Psi$; hence the QQQ anomaly is, up to an overall factor

$$(\text{anomaly})_\Psi = (5)^3 = 125 \quad . \tag{5.6}$$

On the other hand, this anomaly when computed in terms of the "elementary" ψ's and χ's, becomes

$$\begin{aligned}(\text{anomaly})_{\psi,\chi} &= 5(3)^3 + 10(-1)^3 \\ &= 135 - 10 \\ &= 125!! \quad . \end{aligned} \tag{5.7}$$

The 't Hooft relation indeed works, and appears in fact to be quite nontrivial: why should 125 = 135 - 10?? Dimopoulos, Raby and Susskind have constructed further examples of tumbling schemes, which when seen from the Higgs point of view leave massless fermions, and which, when seen from the unbroken, composite point of view lead to identifiable candidates for massless composite fermions which satisfy the 't Hooft consistency criterion. The way this works appears quite magical, and cries out for a general interpretation.

This dual way of viewing the structure of gauge theories has been called complementarity by Dimopoulos, Raby and Susskind. Combined with the 't Hooft criterion (with which there at least exist these examples of compatibility), this concept may prove useful in developing further the technicolor models of gauge-boson and fermion mass generation. However, it is not clear whether, say, the electron should be viewed as an approximately massless composite particle with its mass arising from some dirt effect not included in the formalism. The detailed role of these ideas is not clear--and of course, may not have applicability at all. Nevertheless, they seem to go quite deep, and should at the very least be of utility in understanding better the structure of nonabelian gauge theories in general.

We may conclude from all of this that while there is as yet no good model of extended technicolor, the conceptual structure is extremely attractive, quite rich, and runs quite deep. The game seems to be all-or-nothing, and therefore it is well worth considerable additional effort to search for a scheme which really works.

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	$e\nu_e$	$\mu\nu_\mu$	$\tau\nu_\tau$	ud	us	ub	cd	cs	cb
$e\nu_e$	*	***	**	***	***	✓	✓	*	✓
$\mu\nu_\mu$			**	***	***	✓	*	**	✓
$\tau\nu_\tau$				**	✓	✓	✓	✓	✓
ud				*	**	✓	*	**	✓
us									
ub									
cs									✓
cs									
cb									

Table I: Charged currents

Key:

* Checked to some extent.

✓ Can be checked in future.

	b	c	s	d	u	τ	μ	e	ν_τ	ν_μ	ν_e
b								?		✓	
c								?		✓	
s										✓	
d				?	?		✓	**	?	***	*
u					?		✓	**	?	***	*
τ								✓			
μ								✓		✓	
e								✓		*	*
ν_τ											
ν_μ											
ν_e											

Table II: Neutral currents

Key:

* Checked to some extent.

✓ Can be checked in future.

? Very difficult, but maybe not hopeless.

FIGURE CAPTIONS

Fig. 1: Photon-exchange contribution to weak neutral currents.

Fig. 2: Vacuum-polarization insertions to the photon propagator.

Fig. 3: "Penguin" diagrams which may contribute to $\Delta I = \frac{1}{2}$ enhancement of nonleptonic weak decays.

Fig. 4: Photon-exchange contribution to neutral currents in intermediate-boson models.

Fig. 5: Constraints on Z-boson masses in a two-Z model. Data from PETRA; the analysis is from Ref. 33.

Fig. 6: Lower bound on production of weak quanta in e^+e^- annihilation as function of the range m_W^{-1} of the charged-current weak force.

Fig. 7: Pion contribution to vacuum polarization of intermediate bosons.

Fig. 8: Spectrum of fermion masses versus generation.

Fig. 9: Fermion contribution to gauge-boson vacuum polarization.

Fig. 10: Feynman graphs for $e^+e^- \rightarrow W^+W^-$.

Fig. 11: Higgs-induced contributions to gauge boson vacuum polarization (compare Fig. 9): (a) Transverse part (b) longitudinal part.

Fig. 12: σ-exchange contribution to $e^+e^- \rightarrow W^+W^-$.

Fig. 13: Mechanism for mass generation in technicolor (hypercolor) models.

Fig. 14: Schematic mechanism for dynamical mass-generation for technifermions.

Fig. 15: Dynamical mass-generation via single-gluon emission and absorption.

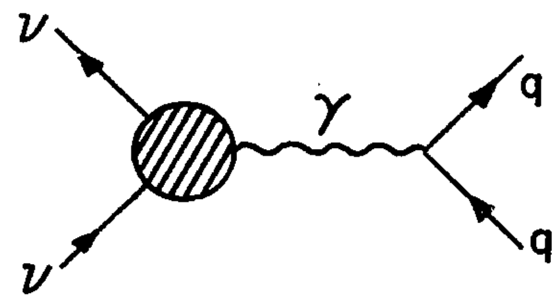

Fig. 1

Fig. 2

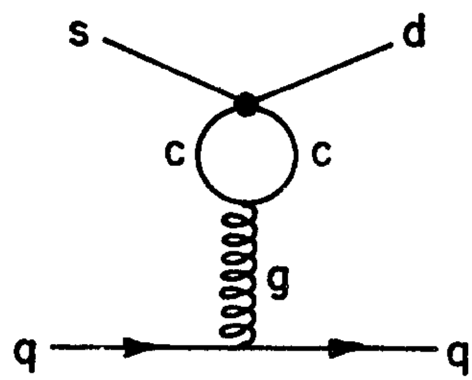

Fig. 3

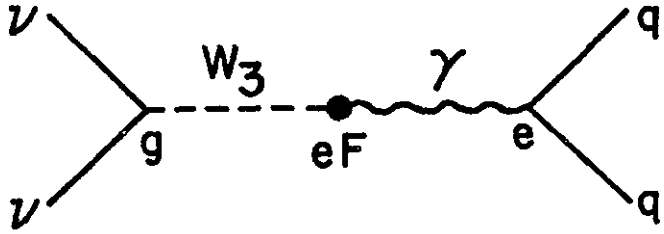

Fig. 4

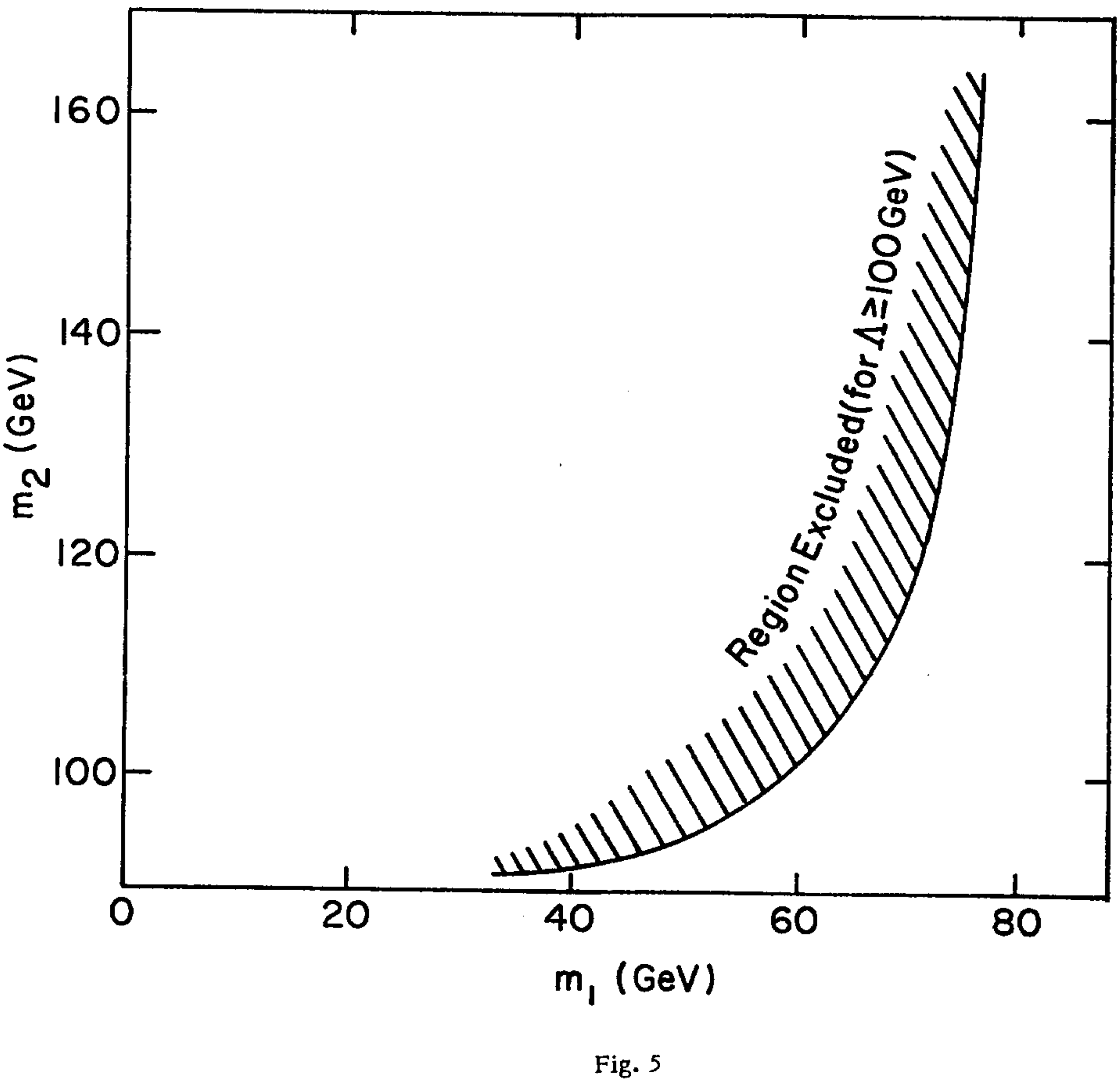
160
140
120
100
m2 (GeV)
Region Excluded (for Λ≥100 GeV)
0
20
40
60
80
m1 (GeV)

Fig. 5

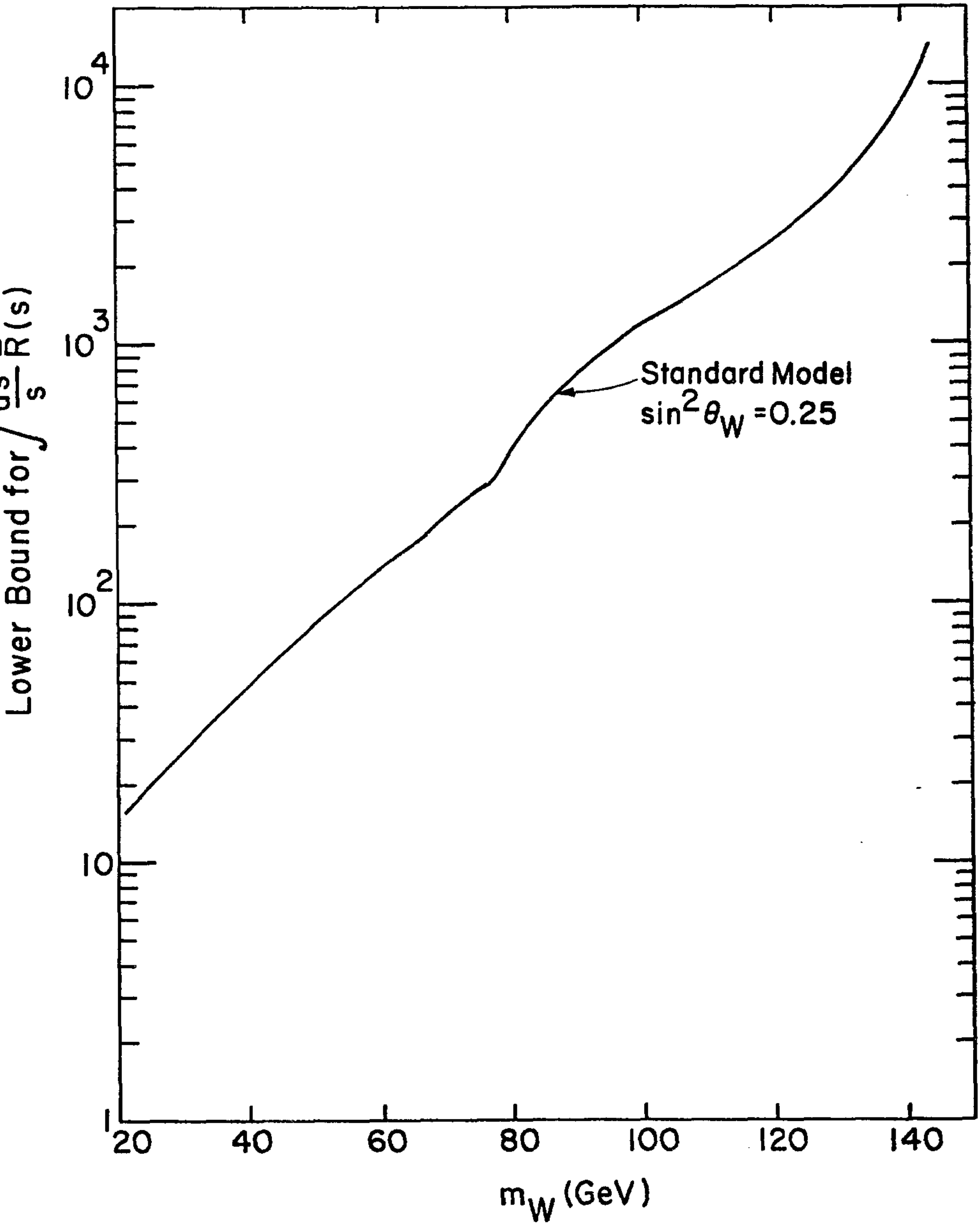
Lower Bound for $\int \frac{ds}{s} \bar{R}(s)$
10^4
10^3
10^2
10
1
Standard Model
$\sin^2\theta_W = 0.25$
20
40
60
80
100
120
140
m_W (GeV)

Fig. 6

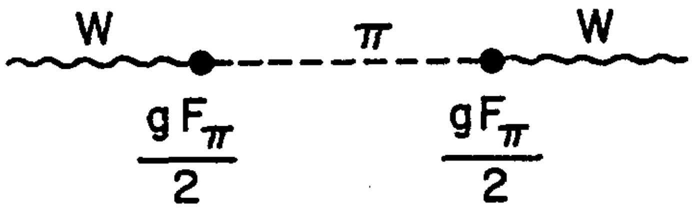

Fig. 7

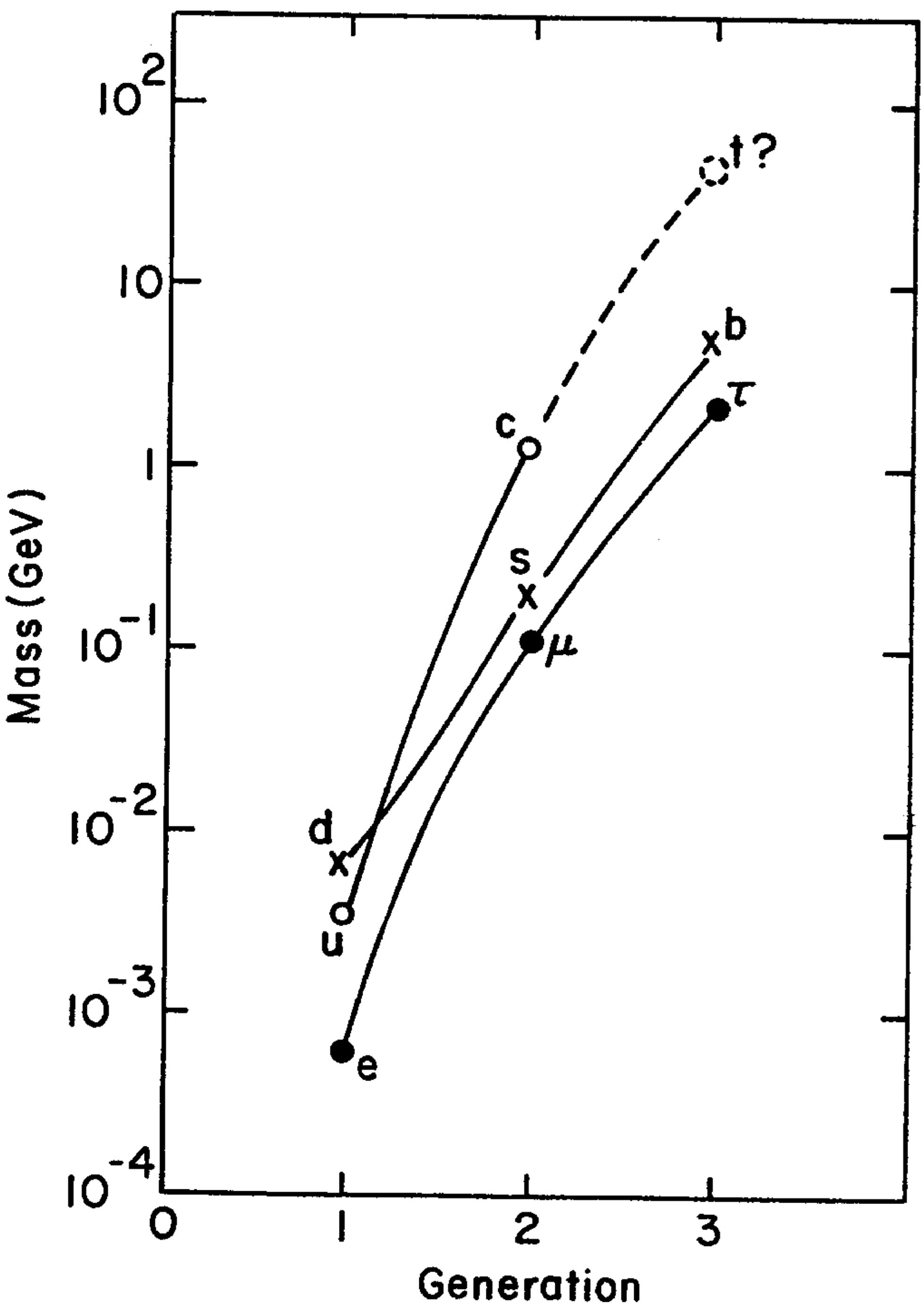

Fig. 8

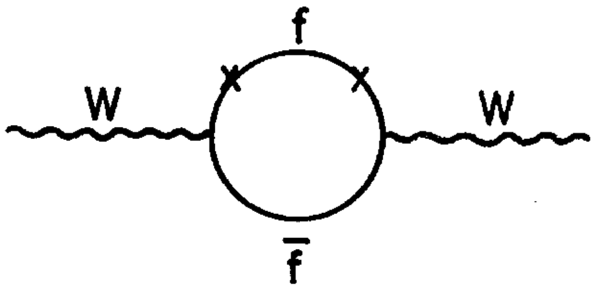

Fig. 9

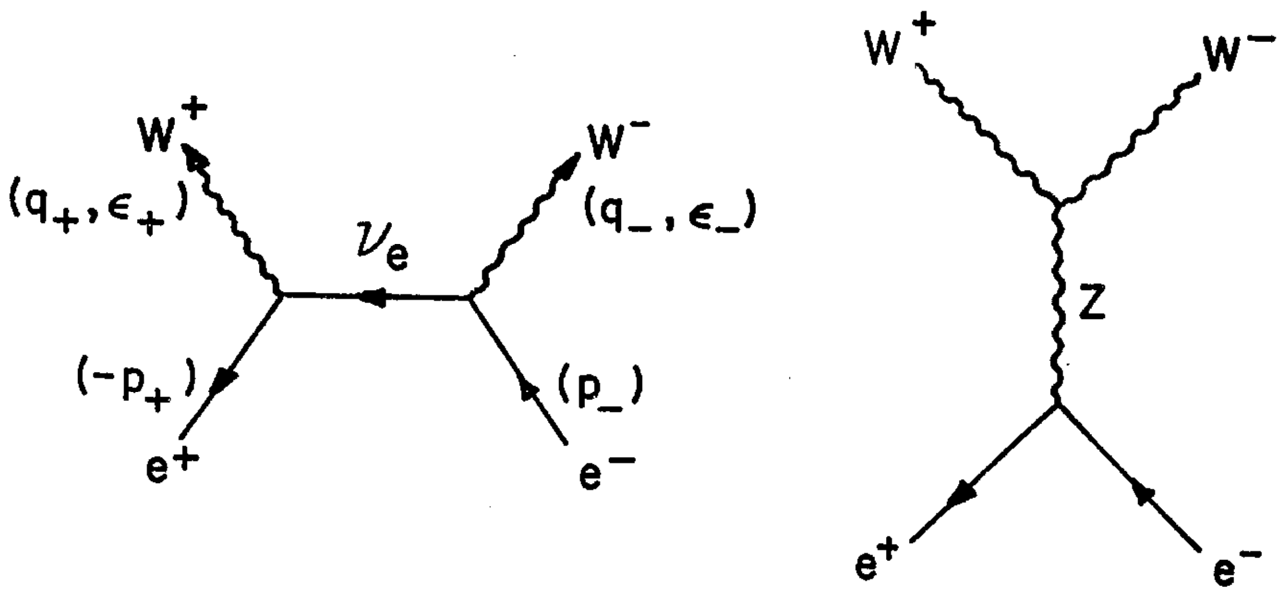

Fig. 10

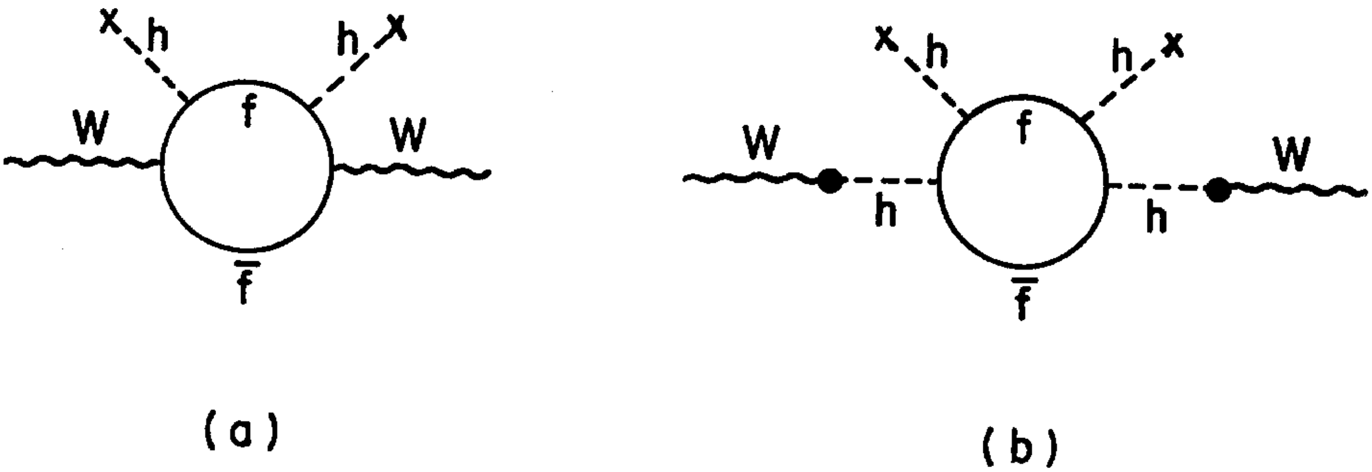

Fig. 11

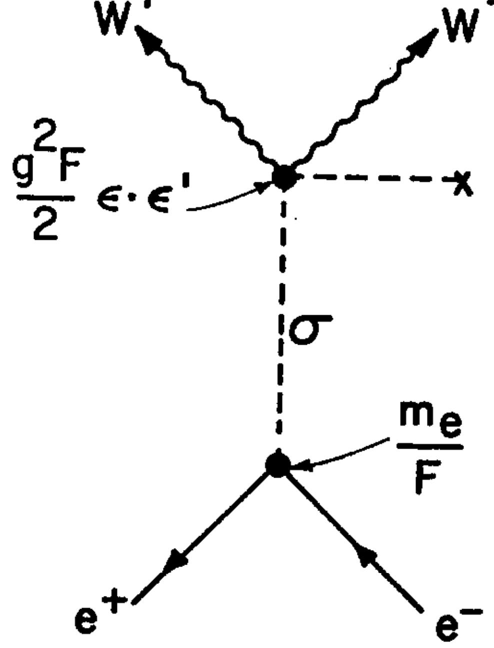

Fig. 12

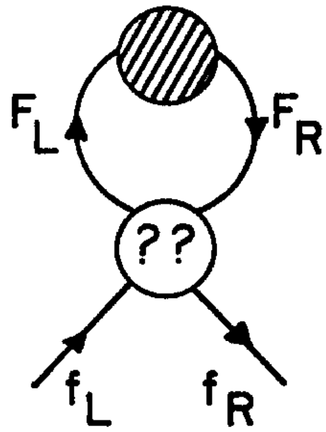

Fig. 13

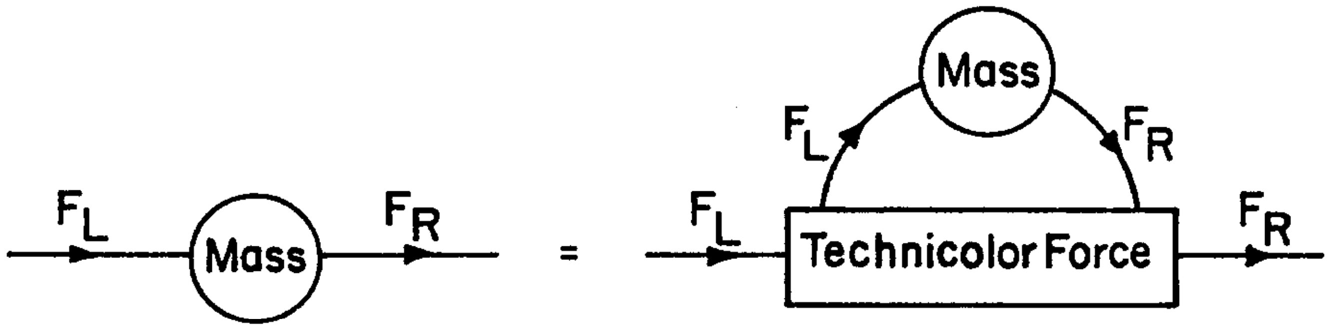

Fig. 14

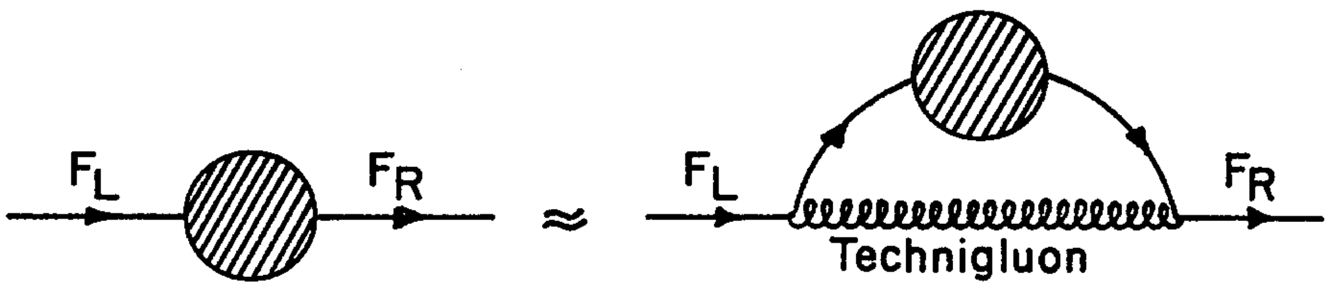

Fig. 15