The standard model [1] is a gauge theory [2] of the microscopic interactions. The strong interaction part (QCD [3]) is described by the Lagrangian

\[ L_{SU_3} = -\frac{1}{4} F_{\mu\nu}^i F^{i\mu\nu} + \sum_r \bar{q}_r \gamma_\mu D^\alpha_\mu q^\beta_r, \]  

(1)

where \( g_s \) is the QCD gauge coupling constant,

\[ F_{\mu\nu}^i = \partial_\mu G_\nu^i - \partial_\nu G_\mu^i - g_s f_{ijk} G_\mu^j G^K_\nu \]  

(2)

is the field strength tensor for the gluon fields \( G_\mu^i, \ i = 1, \ldots, 8 \), and the structure constants \( f_{ijk} \) (\( i, j, k = 1, \ldots, 8 \)) are defined by

\[ [\lambda^i, \lambda^j] = 2i f_{ijk} \lambda^k, \]  

(3)

where the \( SU_3 \) \( \lambda \) matrices are defined in Table 1. The \( F^2 \) term leads to three and four-point gluon self-interactions. The second term in \( L_{SU_3} \) is the gauge covariant

\[ \text{Reprinted from } \textit{Precision Tests of the Standard Electroweak Model}, \text{ ed. P. Langacker} (World, Singapore, 1995). \]
Table 1: The $SU_3$ matrices.

\[
\begin{align*}
\lambda^i &= \left( \begin{array}{ccc}
\tau^i & 0 & 0 \\
0 & 0 & 0 \\
0 & 0 & 0 \\
\end{array} \right), & i &= 1, 2, 3 \\
\lambda^4 &= \left( \begin{array}{ccc}
0 & 1 & 0 \\
0 & 0 & 0 \\
0 & 0 & 0 \\
\end{array} \right) \\
\lambda^5 &= \left( \begin{array}{ccc}
0 & 0 & -i \\
i & 0 & 0 \\
0 & 0 & 0 \\
\end{array} \right) \\
\lambda^6 &= \left( \begin{array}{ccc}
1 & 0 & 0 \\
0 & 0 & 1 \\
0 & 1 & 0 \\
\end{array} \right) \\
\lambda^7 &= \left( \begin{array}{ccc}
0 & 0 & -i \\
i & 0 & 0 \\
0 & i & 0 \\
\end{array} \right) \\
\lambda^8 &= \frac{1}{\sqrt{3}} \left( \begin{array}{ccc}
1 & 0 & 0 \\
0 & 1 & 0 \\
0 & 0 & -2 \\
\end{array} \right)
\end{align*}
\]

The electroweak theory is based on the $SU_2 \times U_1$ Lagrangian (5).

\[\mathcal{L}_{SU_2 \times U_1} = \mathcal{L}_{\text{gauge}} + \mathcal{L}_\varphi + \mathcal{L}_f + \mathcal{L}_{\text{Yukawa}},\]

where the gauge part is

\[\mathcal{L}_{\text{gauge}} = -\frac{1}{4} F_{\mu\nu}^i F^{\mu\nu}_i - \frac{1}{4} B_{\mu\nu} B^{\mu\nu},\]

and

\[B_{\mu\nu} = \partial_\mu B_\nu - \partial_\nu B_\mu, \quad F_{\mu\nu} = \partial_\mu W_\nu^i - \partial_\nu W_\mu^i - g \epsilon_{ijk} W_\mu^j W_\nu^k,\]

where $g(g')$ is the $SU_2$ ($U_1$) gauge coupling and $\epsilon_{ijk}$ is the totally antisymmetric symbol. The $SU_2$ fields have three and four-point self-interactions. $B$ is a $U_1$ field associated with the weak hypercharge $Y = Q - T_3$, where $Q$ and $T_3$ are respectively

\[D_\mu^\alpha = (D_\mu)_{\alpha\beta} = \partial_\mu \delta_{\alpha\beta} + i g_s G^i_\mu L^i_{\alpha\beta},\]
the electric charge operator and the third component of weak $SU_2$. It has no self-interactions. The $B$ and $W_3$ fields will eventually mix to form the photon and $Z$ boson.

The scalar part of the Lagrangian is

$$L_\phi = (D^\mu \phi)^\dagger D_\mu \phi - V(\phi),$$

where $\phi = \begin{pmatrix} \varphi^+ \\ \varphi^0 \end{pmatrix}$ is a complex Higgs scalar, which is a doublet under $SU_2$ with $U_1$ charge $Y_\phi = +\frac{1}{2}$. The gauge covariant derivative is

$$D_\mu \phi = \left( \partial_\mu + ig \frac{\tau^i}{2} W^i_\mu + ig' B_\mu \right) \phi,$$

where the $\tau^i$ are the Pauli matrices. The square of the covariant derivative leads to three and four-point interactions between the gauge and scalar fields $[1]$. $V(\phi)$ is the Higgs potential. The combination of $SU_2 \times U_1$ invariance and renormalizability restricts $V$ to the form

$$V(\phi) = +\mu^2 \phi^\dagger \phi + \lambda (\phi^\dagger \phi)^2.$$  

For $\mu^2 < 0$ there will be spontaneous symmetry breaking. The $\lambda$ term describes a quartic self-interaction between the scalar fields. Vacuum stability requires $\lambda > 0$.

The fermion term is

$$L_F = \sum_{m=1}^F \left( \overline{q}^0_m L \! D q^0_m L + \overline{\ell}^0_m L \! D \ell^0_m L + \overline{u}^0_m R \! D u^0_m R + \overline{d}^0_m R \! D d^0_m R + \overline{e}^0_m R \! D e^0_m R \right).$$

In (11) $m$ is the family index, $F \geq 3$ is the number of families, and $L(R)$ refer to the left (right) chiral projections $\psi_{L(R)} \equiv (1 \mp \gamma_5)\psi/2$. The left-handed quarks and leptons

$$q^0_m L = \begin{pmatrix} u^0_m \\ d^0_m \end{pmatrix}, \quad \ell^0_m L = \begin{pmatrix} \nu^0_m \\ e^{-0}_m \end{pmatrix}$$

transform as $SU_2$ doublets, while the right-handed fields $u^0_{mR}$, $d^0_{mR}$, and $e^{-0}_m$ are singlets. Their $U_1$ charges are $Y_{\psi_L} = \frac{1}{6}$, $Y_{\psi_R} = -\frac{1}{2}$. The superscript 0 refers to the weak eigenstates, i.e., fields transforming according to definite $SU_2$ representations. They may be mixtures of mass eigenstates (flavors). The quark color indices $\alpha = r, g, b$ have been suppressed. The gauge covariant derivatives are

$$D_\mu q^0_{mL} = \left( \partial_\mu + ig \frac{\tau^i}{2} W^i_\mu + ig' B_\mu \right) q^0_{mL} \quad D_\mu u^0_{mR} = \left( \partial_\mu + ig' B_\mu \right) u^0_{mR}$$

$$D_\mu l^0_{mL} = \left( \partial_\mu + ig \frac{\tau^i}{2} W^i_\mu - ig' B_\mu \right) l^0_{mL} \quad D_\mu d^0_{mR} = \left( \partial_\mu - ig' B_\mu \right) d^0_{mR}$$

$$D_\mu e^0_{mR} = \left( \partial_\mu - ig B_\mu \right) e^0_{mR}.$$
from which one can read off the gauge interactions between the $W$ and $B$ and the fermion fields. The different transformations of the $L$ and $R$ fields (i.e., the symmetry is chiral) is the origin of parity violation in the electroweak sector. The chiral symmetry also forbids any bare mass terms for the fermions.

The last term in (5) is

$$ - L_{\text{Yukawa}} = \sum_{m,n=1}^{F} \left[ \Gamma_{mn}^{u} q_{mL} \tilde{\varphi}_{mR}^{0} + \Gamma_{mn}^{d} q_{mL} \varphi_{nR}^{0} + \Gamma_{mn}^{e} \bar{e}_{mL} \varphi_{nR}^{0} \right] + \text{H.C.}, $$

where the matrices $\Gamma_{mn}$ describe the Yukawa couplings between the single Higgs doublet, $\varphi$, and the various flavors $m$ and $n$ of quarks and leptons. One needs representations of Higgs fields with $Y = \frac{1}{2}$ and $-\frac{1}{2}$ to give masses to the down quarks, the electrons, and the up quarks. The representation $\varphi \dagger$ has $Y = -\frac{1}{2}$, but transforms as the $2^*$ rather than the 2. However, in $SU_2$ the $2^*$ representation is related to the 2 by a similarity transformation, and $\tilde{\varphi} \equiv i\tau^2 \varphi \dagger = \left( \begin{array}{c} \varphi^0 \\ -\varphi^- \end{array} \right)$ transforms as a 2 with $Y_{\tilde{\varphi}} = -\frac{1}{2}$. All of the masses can therefore be generated with a single Higgs doublet if one makes use of both $\varphi$ and $\tilde{\varphi}$. The fact that the fundamental and its conjugate are equivalent does not generalize to higher unitary groups. Furthermore, in supersymmetric extensions of the standard model the supersymmetry forbids the use of a single Higgs doublet in both ways in the Lagrangian, and one must add a second Higgs doublet. Similar statements apply to most theories with an additional $U_1$ gauge factor, i.e., a heavy $Z'$ boson.

2 Spontaneous Symmetry Breaking

Gauge invariance (and therefore renormalizability) does not allow mass terms in the Lagrangian for the gauge bosons or for chiral fermions. Massless gauge bosons are not acceptable for the weak interactions, which are known to be short-ranged. Hence, the gauge invariance must be broken spontaneously [5], which preserves the renormalizability [6]. The idea is simply that the lowest energy (vacuum) state does not respect the gauge symmetry and induces effective masses for particles propagating through it.

Let us introduce the complex vector

$$ v = \langle 0 | \varphi | 0 \rangle = \text{constant}, $$

which has components that are the vacuum expectation values of the various complex scalar fields. $v$ is determined by rewriting the Higgs potential as a function of $v$, $V(\varphi) \rightarrow V(v)$, and choosing $v$ such that $V$ is minimized. That is, we interpret $v$ as the lowest energy solution of the classical equation of motion. The quantum theory is obtained by considering fluctuations around this classical minimum, $\varphi = v + \varphi'$.

2It suffices to consider constant $v$ because any space or time dependence $\partial_{\mu} v$ would increase the
The single complex Higgs doublet in the standard model can be rewritten in a Hermitian basis as

\[ \varphi = \begin{pmatrix} \varphi^+ \\ \varphi^0 \end{pmatrix} = \begin{pmatrix} \frac{1}{\sqrt{2}} (\varphi_1 - i\varphi_2) \\ \frac{1}{\sqrt{2}} (\varphi_3 - i\varphi_4) \end{pmatrix}, \]

where \( \varphi_i = \varphi_i^\dagger \) represent four Hermitian fields. In this new basis the Higgs potential becomes

\[ V(\varphi) = \frac{1}{2} \mu^2 \left( \sum_{i=1}^{4} \varphi_i^2 \right) + \frac{1}{4} \lambda \left( \sum_{i=1}^{4} \varphi_i^2 \right)^2, \]

which is clearly \( O_4 \) invariant. Without loss of generality we can choose the axis in this four-dimensional space so that \( \langle 0|\varphi_i|0 \rangle = 0, \ i = 1, 2, 4 \) and \( \langle 0|\varphi_3|0 \rangle = \nu \). Thus, \( V(\varphi) \rightarrow V(v) = \frac{1}{2} \mu^2 \nu^2 + \frac{1}{4} \lambda \nu^4 \),

which must be minimized with respect to \( \nu \). Two important cases are illustrated in Figure 1. For \( \mu^2 > 0 \) the minimum occurs at \( \nu = 0 \). That is, the vacuum is empty space and \( SU_2 \times U_1 \) is unbroken at the minimum. On the other hand, for \( \mu^2 < 0 \) the \( \nu = 0 \) symmetric point is unstable, and the minimum occurs at some nonzero value of \( \nu \) which breaks the \( SU_2 \times U_1 \) symmetry. The point is found by requiring \( V'(\nu) = \nu (\mu^2 + \lambda \nu^2) = 0 \),

which has the solution \( \nu = (-\mu^2/\lambda)^{1/2} \) at the minimum. (The solution for \( -\nu \) can also be transformed into this standard form by an appropriate \( O_4 \) transformation.) The dividing point \( \mu^2 = 0 \) cannot be treated classically. It is necessary to consider the one loop corrections to the potential, in which case it is found that the symmetry is again spontaneously broken [7].

We are interested in the case \( \mu^2 < 0 \), for which the Higgs doublet is replaced, in first approximation, by its classical value \( \varphi \rightarrow \frac{1}{\sqrt{2}} \begin{pmatrix} 0 \\ \nu \end{pmatrix} \equiv v \). The generators \( L^1, L^2, \) and \( L^3 - Y \) are spontaneously broken (e.g., \( L^1 v \neq 0 \)). On the other hand, the vacuum carries no electric charge (\( Qv = (L^3 + Y)v = 0 \)), so the \( U_{1Q} \) of electromagnetism is not broken. Thus, the electroweak \( SU_2 \times U_1 \) group is spontaneously broken down, \( SU_2 \times U_{1Y} \rightarrow U_{1Q} \).

To quantize around the classical vacuum, write \( \varphi = v + \varphi' \), where \( \varphi' \) are quantum fields with zero vacuum expectation value. To display the physical particle content it is useful to rewrite the four Hermitian components of \( \varphi' \) in terms of a new set of variables using the Kibble transformation [8]:

\[ \varphi = \frac{1}{\sqrt{2}} e^{i \sum \xi^i L^i} \begin{pmatrix} 0 \\ \nu + H \end{pmatrix}. \]

energy of the solution. Also, one can take \( \langle 0|\psi|0 \rangle = \langle 0|A_\mu|0 \rangle = 0 \), because any non-zero vacuum values would violate Lorentz invariance. These extensions are involved in (higher energy) topological defects, such as monopoles, strings, domain walls, and textures.
Figure 1: The Higgs potential $V(\nu)$ for $\mu^2 > 0$ (dashed line) and $\mu^2 < 0$ (solid line).

$H$ is a Hermitian field which will turn out to be the physical Higgs scalar. If we had been dealing with a spontaneously broken global symmetry the three Hermitian fields $\xi^i$ would be the massless pseudoscalar Goldstone bosons \footnote{Goldstone bosons} that are necessarily associated with broken symmetry generators. However, in a gauge theory they disappear from the physical spectrum. To see this it is useful to go to the unitary gauge

$$ \varphi \rightarrow \varphi' = e^{-i \sum \xi^i t^i} \varphi = \frac{1}{\sqrt{2}} \begin{pmatrix} 0 \\ \nu + H \end{pmatrix}, \quad (21) $$

in which the Goldstone bosons disappear. In this gauge, the scalar covariant kinetic energy term takes the simple form

$$ (D_\mu \varphi)^\dagger D^\mu \varphi = \frac{1}{2} (0 \nu) \left[ \frac{g}{2} t^i W^i_\mu + \frac{g'}{2} B_\mu \right]^2 \begin{pmatrix} 0 \\ \nu \end{pmatrix} + H \text{ terms} $$

$$ \rightarrow \frac{M_W^2 W^+_{\mu} W^-_{\mu}}{2} + \frac{M_Z^2}{2} Z^\mu Z_{\mu} + H \text{ terms}, \quad (22) $$

where the kinetic energy and gauge interaction terms of the physical $H$ particle have been omitted. Thus, spontaneous symmetry breaking generates mass terms for the $W$ and $Z$ gauge bosons

$$ W^\pm = \frac{1}{\sqrt{2}} (W^1 \mp iW^2) $$

$$ Z = - \sin \theta_W B + \cos \theta_W W^3. \quad (23) $$

The photon field

$$ A = \cos \theta_W B + \sin \theta_W W^3 $$

$$ (24) $$
remains massless. The masses are

\[ M_W = \frac{g\nu}{2} \]  

(25)

and

\[ M_Z = \sqrt{g^2 + g'^2\nu^2} = \frac{M_W}{\cos \theta_W} \]  

(26)

where the weak angle is defined by \( \tan \theta_W \equiv g'/g \). One can think of the generation of masses as due to the fact that the \( W \) and \( Z \) interact constantly with the condensate of scalar fields and therefore acquire masses, in analogy with a photon propagating through a plasma. The Goldstone boson has disappeared from the theory but has reemerged as the longitudinal degree of freedom of a massive vector particle.

It will be seen below that \( \frac{G_F}{\sqrt{2}} \sim \frac{g^2}{8M_W^2} \), where \( G_F = 1.16639(2) \times 10^{-5} \text{ GeV}^{-2} \) is the Fermi constant determined by the muon lifetime. The weak scale \( \nu \) is therefore

\[ \nu = 2M_W/g \approx (\sqrt{2}G_F)^{-1/2} \approx 246 \text{ GeV}. \]  

(27)

Similarly, \( g = e/\sin \theta_W \), where \( e \) is the electric charge of the positron. Hence, to lowest order

\[ M_W = M_Z \cos \theta_W \sim \left(\frac{\pi \alpha/\sqrt{2}G_F}{\sin \theta_W}\right)^{1/2}, \]  

(28)

where \( \alpha \approx 1/137.036 \) is the fine structure constant. Using \( \sin^2 \theta_W \sim 0.23 \) from neutral current scattering, one expects \( M_W \sim 78 \text{ GeV} \), and \( M_Z \sim 89 \text{ GeV} \). (These predictions are increased by \( \sim (2-3) \text{ GeV} \) by loop corrections.) The \( W \) and \( Z \) were discovered at CERN by the UA1 [10] and UA2 [11] groups in 1983. Subsequent measurements of their masses and other properties have been in perfect agreement with the standard model expectations (including the higher-order corrections), as is described in the articles of by Schäile and Einsweiler.

After symmetry breaking the Higgs potential becomes

\[ V(\varphi) = -\frac{\mu^4}{4\lambda} - \mu^2H^2 + \lambda\nu H^3 + \frac{\lambda}{4}H^4. \]  

(29)

The third and fourth terms represent the cubic and quartic interactions of the Higgs scalar. The second term represents a (tree-level) mass

\[ M_H = \sqrt{-2\mu^2} = \sqrt{2\lambda\nu}. \]  

(30)

The weak scale is given in (27), but the quartic Higgs coupling \( \lambda \) is unknown, so \( M_H \) is not predicted. A priori, \( \lambda \) could be anywhere in the range \( 0 \leq \lambda < \infty \). There is now an experimental lower limit \( M_H \gtrsim 60 \text{ GeV} \) from LEP [12]. Otherwise, the decay \( Z \rightarrow Z^*H \) would have been observed. (There are also theoretical lower limits on \( M_H \) in the \( 0-10 \text{ GeV} \) range, depending on \( m_t \), when higher-order corrections are included [13].)
There are also plausible theoretical upper limits. If $\lambda > O(1)$ the theory becomes strongly coupled. ($M_H > O(1\ TeV)$). There is not really anything wrong with strong coupling a priori. However, there are fairly convincing triviality limits, which basically say that the running quartic coupling would become infinite within the domain of validity of the theory if $\lambda$ and therefore $M_H$ is too large. If one requires the theory to make sense to infinite energy, one may run into problems with the increasing quartic coupling for any $\lambda$. However, one only needs for the theory to hold up to the next mass scale $\Lambda$, at which point the standard model breaks down. In that case \[13\],

$$M_H < \begin{cases} O(200) \text{ GeV}, & \Lambda \sim M_P \\ O(600) \text{ GeV}, & \Lambda \sim 2M_H \end{cases} \quad (31)$$

The more stringent limit of $O(200) \text{ GeV}$ obtains for $\Lambda$ of order of the Planck scale $M_P = G_N^{-1/2} \sim 10^{19} \text{ GeV}$. If one makes the less restrictive assumption that the scale $\Lambda$ of new physics can be small, one obtains a weaker limit. Nevertheless, for the concept of an elementary Higgs field to make sense one should require that the theory be valid up to something of order of $2M_H$, which implies that $M_H < O(600) \text{ GeV}$. These limits may be relaxed if there are other heavy particles in the theory.

The first term in (29) is the vacuum expectation value

$$\langle 0 | V | 0 \rangle = -\mu^4 / 4\lambda \quad (32)$$

of the Higgs potential when evaluated at the minimum. This is a $c$-number which has no significance for the microscopic interactions. However, it assumes great importance when the theory is coupled to gravity, because a constant energy density plays the role of a cosmological constant \[14\]. The cosmological constant becomes

$$\Lambda_{\text{cosm}} = \Lambda_{\text{bare}} + \Lambda_{\text{SSB}}, \quad (33)$$

where $\Lambda_{\text{bare}}$ is the primordial cosmological constant, which can be thought of as the value of the energy of the vacuum in the absence of spontaneous symmetry breaking. (Eqn. (10) implicitly assumed $\Lambda_{\text{bare}} = 0$.) $\Lambda_{\text{SSB}}$ is the part generated by the Higgs mechanism:

$$|\Lambda_{\text{SSB}}| = 8\pi G_N |\langle 0 | V | 0 \rangle| \sim 10^{50} |\Lambda_{\text{obs}}|. \quad (34)$$

It is some $10^{50}$ times larger than the observational upper limit $\Lambda_{\text{obs}}$. This is clearly unacceptable. Technically, one can solve the problem by adding a constant $+\mu^4 / 4\lambda$ to $V$, so that $V$ is equal to zero at the minimum (i.e., $\Lambda_{\text{bare}} = 2\pi G_N \mu^4 / \lambda$). However, with our current understanding there is no reason for $\Lambda_{\text{bare}}$ and $\Lambda_{\text{SSB}}$ to be related;

\[3\]This is true for a pure $\lambda H^4$ theory. The presence of other interactions may eliminate the problems for small $\lambda$.\]
to have to invoke such an incredibly fine-tuned cancellation to 50 decimal places is a major unsatisfactory feature of the standard model.

The Yukawa interaction in the unitary gauge becomes

$$- L_{\text{Yukawa}} \rightarrow \sum_{m,n=1}^{F} \bar{u}_{mL}^{0} \Gamma^{u}_{mn} \left( \frac{\nu + H}{\sqrt{2}} \right) u_{mR}^{0} + (d,e) \text{ terms} + \text{ H.C.}$$

where in the second form $u_{L}^{0} = (u_{1L}^{0}, u_{2L}^{0}, \ldots, u_{FL}^{0})^{T}$ is an $F$-component column vector, with a similar definition for $u_{R}^{0}$. $M^{u}_{mn} = \Gamma^{u}_{mn} \nu/\sqrt{2}$ is induced by spontaneous symmetry breaking, and $h^{u} = M^{u}/\nu = g M^{u}/2 M_{W}$ is the Yukawa coupling matrix.

In general $M$ is not diagonal, Hermitian, or symmetric. To identify the physical particle content it is necessary to diagonalize $M$ by separate unitary transformations $A_{L}$ and $A_{R}$ on the left- and right-handed fermion fields. (In the special case that $M^{u}$ is Hermitian one can take $A_{L} = A_{R}$.) Then,

$$A_{L}^{u \dagger} M^{u} A_{R}^{v} = M^{D}_{D} = \begin{pmatrix} m_{u} & 0 & 0 \\ 0 & m_{c} & 0 \\ 0 & 0 & m_{t} \end{pmatrix}$$

is a diagonal matrix with eigenvalues equal to the physical masses of the charge $2/3$ quarks. Similarly, one diagonalizes the down quark and charged lepton mass matrices by

$$A_{L}^{d \dagger} M^{d} A_{R}^{d} = M_{D}^{d}$$
$$A_{L}^{e \dagger} M^{e} A_{R}^{e} = M_{D}^{e}.$$  \hspace{1cm} (37)

In terms of these unitary matrices we can define mass eigenstate fields $u_{L} = A_{L}^{u \dagger} u_{L}^{0} = (u_{L}, c_{L}, t_{L})^{T}$, with analogous definitions for $u_{R} = A_{R}^{u \dagger} u_{R}^{0}$, $d_{L,R} = A_{L,R}^{d \dagger} d_{L,R}^{0}$, and $e_{L,R} = A_{L,R}^{e \dagger} e_{L,R}^{0}$. Assuming the neutrinos are massless, their mass eigenstates are arbitrary. It is convenient to define them in terms of the charged lepton unitary transformation, $\nu_{L} = A_{L}^{e \dagger} \nu_{L}^{0}$. That is, we define $\nu_{e}, \nu_{\mu}, \nu_{\tau}$ as the weak interaction partners of the $e, \mu,$ and $\tau$. Typical estimates of the quark masses are $m_{u} = 5.6 \pm 1.1 \text{ MeV}$, $m_{d} = 9.9 \pm 1.1 \text{ MeV}$, $m_{s} = 199 \pm 33 \text{ MeV}$, $m_{c} = 1.35 \pm 0.05 \text{ GeV}$, $m_{b} \sim 4.7 \text{ GeV}$, and $m_{t} > 131 \text{ GeV}$. These are the current masses: for QCD their effects are identical to bare masses in the QCD Lagrangian. They should not be confused with the constituent masses of order 300 MeV generated by the spontaneous breaking of chiral symmetry in the strong interactions. Including QCD renormalizations, the $u, d, s$ and $c$ masses are running masses evaluated at 1 GeV$^{2}$, while $m_{b}$ and $m_{t}$ are pole masses.

Thus,

$$L_{\text{Yukawa}} = \sum_{i} \bar{\psi}_{i} \left( -m_{i} - \frac{g m_{i}}{2 M_{W}} H \right) \psi_{i}. \hspace{1cm} (38)$$
The coupling of the physical Higgs boson to the $i^{\text{th}}$ fermion is $g m_i / 2 M_W$, which is very small except for the top quark. The coupling is flavor-diagonal in the minimal model: there is just one Yukawa matrix for each type of fermion, so the mass and Yukawa matrices are diagonalized by the same transformations. In generalizations in which more than one Higgs doublet couples to each type of fermion there will in general be flavor-changing Yukawa interactions involving the physical neutral Higgs fields [18]. There are stringent limits on such couplings [19]; for example, the $K_L - K_S$ mass difference implies $h / M_H < 10^{-6} \text{GeV}^{-1}$, where $h$ is the $\bar{d}s$ Yukawa coupling.

3 The Gauge Interactions

The major quantitative tests of the electroweak standard model involve the gauge interactions of fermions and the properties of the gauge bosons. The charged current weak interactions of the Fermi theory and its extension to the intermediate vector boson theory are incorporated into the standard model, as is quantum electrodynamics. The theory successfully predicted the existence and properties of the weak neutral current. Here I will summarize the structure of the interactions. Later chapters will discuss the phenomenology and tests in more detail.

3.1 The Charged Current

The interaction of the $W$ bosons to fermions is given by

$$L = -\frac{g}{2 \sqrt{2}} \left( J_W^\mu W^-_\mu + J_W^{\mu\dagger} W^+_\mu \right) ,$$

where the weak charge-raising current is

$$J_W^{\mu\dagger} = \sum_{m=1}^F \left[ \bar{\nu}_m \gamma^\mu (1 - \gamma^5) e_m + \bar{u}_m \gamma^\mu (1 - \gamma^5) d_m \right] (40)$$

$J_W^{\mu\dagger}$ has a $V - A$ form, i.e., it violates parity and charge conjugation maximally. The mismatch between the unitary transformations relating the weak and mass eigenstates for the up and down-type quarks leads to the presence of the $F \times F$ unitary matrix $V = A_L^\dagger A_L$ in the current. This is the Cabibbo-Kobayashi-Maskawa (CKM) matrix [20], which is ultimately due to the mismatch between the weak and Yukawa interactions. For $F = 2$ families $V$ takes the familiar form

$$V = \begin{pmatrix} \cos \theta_c & \sin \theta_c \\ -\sin \theta_c & \cos \theta_c \end{pmatrix} ,$$

(41)
Figure 2: A weak interaction mediated by the exchange of a $W$ and the effective four-fermi interaction that it generates if the four-momentum transfer $Q$ is sufficiently small.

where $\sin \theta_c \simeq 0.22$ is the Cabibbo angle. This form gives a good zeroth-order approximation to the weak interactions of the $u, d, s$ and $c$ quarks; their coupling to the third family, though non-zero, is very small. Including these couplings, the 3-family CKM matrix is

$$V = \begin{pmatrix} V_{ud} & V_{us} & V_{ub} \\ V_{cd} & V_{cs} & V_{cb} \\ V_{td} & V_{ts} & V_{tb} \end{pmatrix},$$

(42)

where the $V_{ij}$ may involve a CP-violating phase.

There is nothing to distinguish massless neutrinos except their weak interactions, so one simply defines the $\nu_e$ as the weak partner of the electron, and similarly for $\nu_\mu$ and $\nu_\tau$. If there were non-zero neutrino mass then one would have to introduce a leptonic mixing matrix in the current, but its effects would not be important in any process that is not actually sensitive to the masses.

The interaction between fermions mediated by the exchange of a $W$ is illustrated in Figure 2. In the limit $|Q^2| \ll M_W^2$ the momentum term in the $W$ propagator can be neglected, leading to an effective zero-range (four-fermi) interaction

$$-L_{\text{eff}}^{cc} = \frac{G_F}{\sqrt{2}} J_W^\mu J_W^{\mu*},$$

(43)

where the Fermi constant is identified as

$$\frac{G_F}{\sqrt{2}} \simeq \frac{g^2}{8M_W^2} = \frac{1}{2\nu^2}.$$  

(44)

Thus, the Fermi theory is an approximation to the standard model valid in the limit of small momentum transfer. From the muon lifetime, $G_F = 1.16639(2) \times 10^{-5} \text{ GeV}^{-2}$, which implies that the weak interaction scale defined by the VEV of the Higgs field is $\nu = \sqrt{2}(0)|\phi^0|0) \simeq 246 \text{ GeV}$. 
The charged current weak interaction as described by (43) has been successfully tested in a large variety of weak decays [21], including $\beta$, $K$, hyperon, heavy quark, $\mu$, and $\tau$ decays. In particular, high precision measurements of $\beta$, $\mu$, and $\tau$ decays are a sensitive probe of extended gauge groups involving right-handed currents and other types of new physics, as is described in the chapters by Deutsch and Quin; Fetscher and Gerber; and Herczeg. Tests of the unitarity of the CKM matrix are important in searching for the presence of fourth family or exotic fermions and for new interactions, as described by Sirlin and by London. The standard theory has also been successfully probed in neutrino scattering processes such as $\nu_\mu e \rightarrow \mu^- \nu_e$, $\nu_\mu n \rightarrow \mu^- p$, $\nu_\mu N \rightarrow \mu^- X$. It works so well that the neutrino-hadron interactions are used more as a probe of the structure of the hadrons and QCD than as a test of the weak interactions.

Weak charged current effects have also been observed in higher orders, such as in the mass difference $M_{K_S} - M_{K_L}$, $CP$ violation in the kaon system [22], and in $B \leftrightarrow \bar{B}$ mixing [23]. For these higher order processes the full theory must be used because large momenta occur within the loop integrals.

3.2 QED

The standard model incorporates all of the (spectacular) successes of quantum electrodynamics (QED) [24], which is based on the $U_{1Q}$ subgroup that remains unbroken after spontaneous symmetry breaking. The relevant part of the Lagrangian is

$$L = -\frac{gg'}{\sqrt{g^2 + g'^2}} J^\mu_Q(\cos \theta_W B_\mu + \sin \theta_W W^3_\mu),$$

(45)

where the linear combination of neutral gauge fields is just the photon field $A_\mu$. This reproduces the QED interaction provided one identifies the combination of couplings

$$e = g \sin \theta_W$$

(46)

as the electric charge of the positron, where $\tan \theta_W \equiv g'/g$. The electromagnetic current is given by

$$J^\mu_Q = \sum_{m=1}^{F} \left[ \frac{2}{3} \bar{u}_m \gamma^\mu u_m - \frac{1}{3} \bar{d}_m \gamma^\mu d_m - \bar{e}_m \gamma^\mu e_m \right]$$

(47)

It takes the same form when written in terms of either weak or mass eigenstates because all fermions which mix with each other have the same electric charge. Thus, the electromagnetic current is automatically flavor-diagonal.
Figure 3: Typical neutral current interaction mediated by the exchange of the Z, which reduces to an effective four-fermi interaction in the limit that the momentum transfer $Q$ can be neglected.

3.3 The Neutral Current

The third class of gauge interactions is the weak neutral current \[25\], which was predicted by the $SU_2 \times U_1$ model. The relevant interaction is

$$L = -\frac{\sqrt{g^2 + g'^2}}{2} J^\mu_Z \left( -\sin \theta_W B\mu + \cos \theta_W W^3_{\mu} \right),$$

where the combination of neutral fields is the massive Z boson field. The strength is conveniently rewritten as $g/(2 \cos \theta_W)$, which follows from $\cos \theta_W = g/\sqrt{g^2 + g'^2}$.

The weak neutral current is given by

$$J^\mu_Z = \sum_m \left[ \bar{u}_{mL} \gamma^\mu u_{mL} - \bar{d}_{mL} \gamma^\mu d_{mL} + \bar{\nu}_{mL} \gamma^\mu \nu_{mL} - \bar{e}_{mL} \gamma^\mu e_{mL} \right] - 2 \sin^2 \theta_W J^\mu_Q$$

Like the electromagnetic current $J^\mu_Z$ is flavor-diagonal in the standard model; all fermions which have the same electric charge and chirality and therefore can mix with each other have the same $SU_2 \times U_1$ assignments, so the form is not affected by the unitary transformations that relate the mass and weak bases. It was for this reason that the GIM mechanism \[26\] was introduced into the model, along with its prediction of the charm quark. Without it the $d$ and $s$ quarks would not have had the same $SU_2 \times U_1$ assignments, and flavor-changing neutral currents would have resulted. The absence of such effects is a major restriction on many extensions of the standard model involving exotic fermions, as described in the article by London.

The neutral current has two contributions. The first only involves the left-chiral fields and is purely $V - A$. The second is proportional to the electromagnetic current with coefficient $\sin^2 \theta_W$ and is purely vector. Parity is therefore violated in the neutral current interaction, though not maximally.
In an interaction between fermions in the limit that the momentum transfer is small compared to $M_Z$ one can neglect the $Q^2$ term in the propagator, and the interaction reduces to an effective four-fermi interaction

$$-L_{NC}^{\text{eff}} = \frac{G_F}{\sqrt{2}} J^\mu Z J_{Z\mu}. \quad (50)$$

The coefficient is the same as in the charged case because

$$\frac{G_F}{\sqrt{2}} = \frac{g^2}{8M_W^2} = \frac{g^2 + g'^2}{8M_Z^2}. \quad (51)$$

That is, the difference in $Z$ couplings compensates the difference in masses in the propagator. The weak neutral current was discovered at CERN in 1973 by the Gargamelle [27] collaboration and by HPW at Fermilab [28] shortly thereafter, and since that time it has been extensively studied in many interactions, including $\nu e \to \nu e$, $\nu N \to \nu N$, $\nu N \to \nu X$; $e^\uparrow \downarrow D \to eX$; atomic parity violation; $e^+e^-$ and $Z$-pole reactions. These have been the primary quantitative test of the unification part of the standard electroweak model, and all aspects will be discussed extensively in later chapters.

### 3.4 Gauge Self-interactions

The self-interactions of the gauge bosons in the standard model are displayed in Figure 4. Their form is predicted by the underlying gauge invariance, but they have not yet been tested. A sensitive test will have to wait for a study of $e^+e^- \to W^+W^-$ in the second phase, LEP II, of the $e^+e^-$ collider at CERN, as described in the article by Treille [29], and future possible $e^+e^-$ and hadron colliders at higher energy. To lowest order there are three diagrams, as shown in Figure 5. Two of them involve the three-point interaction between a photon or $Z$ boson and $W^+W^-$. The cross section from any one of these diagrams rises with center of mass energy, but gauge invariance relates these three-point vertices to the couplings of the fermions in such a way that at high energies there is a cancellation. It is another manifestation of the same cancellation which brings higher-order loop integrals under control, leading to a renormalizable theory (otherwise, vector theories would have severe divergences). At LEP II one will be able to observe the cancellation; it would be even more dramatic at possible future colliders at higher energies. Detailed studies of $e^+e^- \to W^+W^-$ should be sensitive to deviations from the standard model, especially those associated with such non-gauge physics as compositeness. In practice, however, many of the types of new physics which could lead to observable effects are already excluded by other observables at LEP I and elsewhere [30].

The processes $(\bar{q} q \to VV')$ would also be sensitive to gauge self-interactions. Finally, one can study the gauge-gauge three and four point vertices in the processes $e^+e^- \to e^+e^-VV'$ and $(\bar{q} q \to \bar{q} qVV')$. These tests involve the same reactions that
Figure 4: The three and four point-self-interactions of gauge bosons in the standard electroweak model.

\[ F_{\nu\lambda\mu} = g_{\nu\lambda}(k_1 - k_2)_\mu + g_{\lambda\mu}(k_2 - k)_\nu + g_{\mu\nu}(k - k_1)_\lambda \]

\[ \frac{ieF_{\nu\lambda\mu}}{\tan \theta_W} \]

\[ -ie^2 G_{\alpha\beta\mu\nu} \]

\[ G_{\alpha\beta\mu\nu} = 2g_{\alpha\beta}g_{\mu\nu} - g_{\alpha\mu}g_{\beta\nu} - g_{\alpha\nu}g_{\beta\mu} \]

\[ \frac{-ie^2}{\tan \theta_W} G_{\alpha\beta\mu\nu} \]

\[ \frac{-ie^2}{\tan^2 \theta_W} G_{\alpha\beta\mu\nu} \]

\[ \frac{ie^2}{\sin^2 \theta_W} G_{\alpha\beta\mu\nu} \]
would be used to search for a very heavy Higgs boson at a high energy hadron collider and will be important not only for their own sake but as necessary background for the Higgs search.

4 Problems With the Standard Model

The Lagrangian for the standard model after spontaneous symmetry breaking is

\[
L = L_{\text{gauge}} + L_{\text{Higgs}} + \sum_i \bar{\psi}_i \left( i \frac{\partial}{\partial i} - m_i - \frac{m_i H}{\nu} \right) \psi_i \\
- \frac{g}{2\sqrt{2}} \left( J_{W\mu}^\mu W^- + J_{W\mu}^{\dagger\mu} W^+_\mu \right) - e J_{Q\mu} A_\mu - \frac{g}{2\cos\theta_W} J_{Z\mu} Z_\mu.
\] (52)

The standard electroweak model is a mathematically-consistent renormalizable field theory which predicts or is consistent with all experimental facts. It successfully predicted the existence and form of the weak neutral current, the existence and masses of the W and Z bosons, and the charm quark, as necessitated by the GIM mechanism. The charged current weak interactions, as described by the generalized Fermi theory, were successfully incorporated, as was quantum electrodynamics. When combined with quantum chromodynamics for the strong interactions and general relativity for classical gravity, the standard model is almost certainly the approximately correct description of nature down to at least $10^{-16}$ cm, with the possible exception of the Higgs sector. However, the theory has far too much arbitrariness to be the final story. For example, the minimal version of the model has 21 free parameters, assuming massless neutrinos and not counting electric charge (Y) assignments. Most physicists believe that this is just too much for the fundamental theory. The complications of the standard model can also be described in terms of a number of problems.

1. Gauge Problem

The standard model is a complicated direct product of three sub-groups, $SU_3 \times SU_2 \times U_1$, with separate gauge couplings. There is no explanation for why only the
electroweak part is chiral (parity-violating). Similarly, the standard model incorporates but does not explain another fundamental fact of nature: charge quantization, i.e., why all particles have charges which are multiples of $e/3$. This is important because it allows the electrical neutrality of atoms ($|q_p| = |q_e|$). Possible explanations include: grand unified theories, the existence of magnetic monopoles, and constraints from the absence or cancellation of anomalies.

2. Fermion Problem

All matter under ordinary terrestrial conditions can be constructed out of the fermions ($\nu_e, e^-, u, d$) of the first family. Yet we know from laboratory studies that there are $\geq 3$ families: ($\nu_\mu, \mu^-, c, s$) and ($\nu_\tau, \tau^-, t, b$) are heavier copies of the first family with no obvious role in nature. (The $t$ and $\nu_\tau$ have not yet been directly observed, although there are candidate $t$ events from CDF.) They are assumed to exist because the weak interactions of the $b$ and $\tau$ have been well measured and are in agreement with the assumptions that they have $SU_2$-doublet partners. The standard model gives no explanation for the existence of these heavier families and no prediction for their numbers. Furthermore, there is no explanation or prediction of the fermion masses, which vary over at least five orders of magnitude, or of the CKM mixings. There are many possible suggestions of new physics that might shed light on this, including composite fermions; family symmetries; radiative hierarchies, in which the fermion masses are generated at the loop-level, with the lighter families requiring more loops; and the topology of extra space-time dimensions, such as in superstring models. Despite all of these ideas there is no compelling model and none of these yields detailed predictions. The problem is just too complicated. Simple grand unified theories don’t help very much with this, except for the prediction of $m_\nu$ in terms of $m_\tau$ in the simplest versions.

3. Higgs/hierarchy Problem

In the standard model one introduces an elementary Higgs field into the theory to generate masses for the $W$, $Z$, and fermions. For the model to be consistent the Higgs mass should not be too different from the $W$ mass, i.e., $M_H^2 = O(M_W^2)$. If $M_H$ were to be larger than $M_W$ by many orders of magnitude there would be a hierarchy problem, and the Higgs self-interactions would be excessively strong. Combining theoretical arguments with laboratory limits one obtains $M_H \lesssim 1$ TeV. (See (31)).

However, there is a complication. The tree-level (bare) Higgs mass receives quadratically-divergent corrections from the loop diagrams in Figure 6. One finds

$$M_H^2 = (M_H^2)_{\text{bare}} + O(\lambda, g^2, h^2)\Lambda^2,$$

where $\Lambda$ is the next higher scale in the theory. If there were no higher scale one would simply interpret $\Lambda$ as an ultraviolet cutoff and take the view that $M_H$ is a measured parameter and that $(M_H)_{\text{bare}}$ is not an observable. However, the theory is

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4The absence of anomalies is not sufficient to determine all of the $Y$ assignments without additional assumptions, such as family universality.
presumably embedded in some larger theory that cuts off the integral at the finite scale of the new physics\textsuperscript{5}. For example, if the next scale is gravity $\Lambda$ is the Planck scale $M_P = G_N^{-1/2} \sim 10^{19}$ GeV. If there is a simple grand unified theory \cite{31}, one would expect $\Lambda$ to be of order the unification scale $M_X \sim 10^{14}$ GeV. Hence, the natural scale for $M_H$ is $O(\Lambda)$, which is much larger than the expected value. There must be a fine-tuned and apparently highly contrived cancellation between the bare value and the correction, to more than 30 decimal places in the case of gravity. If the cutoff is provided by a grand unified theory there is a separate hierarchy problem at the tree-level. The tree-level couplings between the Higgs field and the superheavy fields lead to the expectation that $M_H$ is equal to the unification scale unless unnatural fine-tunings are done.

One solution to this Higgs/Hierarchy problem is the possibility that the $W$ and $Z$ bosons are composite. However, in this case one would apparently be throwing out the successes of the $SU_2 \times U_1$ gauge theory. Another approach is to eliminate elementary Higgs fields in favor of a dynamical mechanism in which they are replaced by bound states of fermions. Technicolor and composite Higgs models are in this category \cite{38}. The third possibility is supersymmetry \cite{39}, which prevents large renormalizations by enforcing cancellations between the various diagrams in Figure 6. However, most grand unified versions do not explain why $(M_W/M_X)^2$ is so small in the first place.

4. Strong $CP$ Problem

Another fine-tuning problem is the strong $CP$ problem \cite{40}. One can add an additional term $\frac{\theta}{32\pi^2} g_s^2 F\tilde{F}$ to the QCD Lagrangian which breaks $P$, $T$ and $CP$ symmetry.

\textsuperscript{5}There is no analogous fine-tuning associated with logarithmic divergences, such as those encountered in QED, because $\alpha \ln(\Lambda/m_e) < O(1)$ even for $\Lambda = M_P$. 

Figure 6: Radiative corrections to the Higgs mass, including self-interactions, interactions with gauge bosons, and interactions with fermions.
$\tilde{F}_{\mu\nu} = \epsilon_{\mu\nu\alpha\beta}F^{\alpha\beta}/2$ is the dual field. This term, if present, would induce an electric dipole moment $d_N$ for the neutron. The rather stringent limits on the dipole moment [41] lead to the upper bound $\theta < 10^{-10}$. The question is, therefore, why is $\theta$ so small? It is not sufficient to just say that it is zero because $CP$ violation in the weak interactions leads to a radiative correction or renormalization of $\theta$ by $O(10^{-3})$. Therefore, an apparently contrived fine-tuning is needed to cancel this correction against the bare value. Solutions include the possibility that $CP$ violation is not induced directly by phases in the Yukawa couplings, as is usually assumed in the standard model, but is somehow violated spontaneously [40]. $\theta$ then would be a calculable parameter induced at loop level, and it is possible to make $\theta$ sufficiently small. However, such models lead to difficult phenomenological and cosmological problems\(^6\). Alternately, $\theta$ becomes unobservable if there is a massless $u$ quark [43]. However, most phenomenological estimates are not consistent with $m_u = 0$ [15, 44]. Another possibility is the Pececi-Quinn mechanism [45], in which an extra global $U_1$ symmetry is imposed on the theory in such a way that $\theta$ becomes a dynamical variable which is zero at the minimum of the potential. Such models imply the existence of very light pseudoscalar particles called axions. Laboratory, astrophysical, and cosmological constraints allow only the range $10^8 - 10^{12}$ GeV for the scale at which the $U_1$ symmetry is broken.

5. **Graviton Problem**

Gravity is not fundamentally unified with the other interactions in the standard model, although it is possible to graft on classical general relativity by hand. However, this is not a quantum theory, and there is no obvious way to generate one within the standard model context. In addition to the fact that gravity is not unified and not quantized there is another difficulty, namely the cosmological constant. The cosmological constant can be thought of as energy of the vacuum. The energy density induced by spontaneous symmetry breaking is some 50 orders of magnitude larger than the observational upper limit (see Eqns. (33) and (34)). This implies the necessity of severe fine-tuning between the generated and bare pieces, which do not have any a priori reason to be related. Possible solutions include Kaluza-Klein [46] and supergravity theories [39]. These unify gravity but do not solve the problem of quantum gravity or yield renormalizable theories of quantum gravity, nor do they provide any obvious solution to the cosmological constant problem. Superstring theories [36] unify gravity and may yield finite theories of quantum gravity and all the other interactions. It is not clear whether or not they solve the cosmological constant problem.

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