Introduction to Higgs Sector: Theory and Phenomenology

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1 Introduction

This review is based on lectures delivered at International Summer School on High Energy Physics, in Akyaka, Muğla, Turkey, September 2006. The goal of the review is to introduce junior graduate students, having a some knowledge of field theory, to the problem of mass generation and its physical consequences. The review starts with the statements and a toy model of mass generation in Sect. 2. In Sect. 3 a model of weak interactions is constructed and its predictions are compared with experimental results obtained so far. Given in Sect. 4 is search strategies for Higgs boson, the only remnant of the mechanism that gives mass to matter, by referring to model predictions and experimental facilities. Finally, in Sect. 5, given are a summary of results plus reasons for and expectations from models which rehabilitate the minimal model constructed in Sect. 3[1].

2 Mass Generation

The inertial mass of a particle is a measure of its resistance to external forces acting upon, and of its energy content when at rest. In field theory language, mass of a particle is given by the pole of its propagator – the probability amplitude for particle to travel with a given energy and momentum. What is nontrivial about mass is that it is correlated with the symmetry properties of the field. In other words, symmetries and mass of a field are two intimately related structures. For a clearer view of this connection, it may be instructive to examine the electromagnetic field. This field is strictly massless. In fact,

[1] For the ease of grasping the fundamental concepts, and learning about alternative viewpoints present in the literature the references are chosen mostly from books and review papers and/or e-prints.

it has to be so since a finite mass explicitly breaks gauge invariance of electromagnetism, and hence, violates conservation of electric charge. Therefore, one arrives at a definite conclusion about the mass of photon on the basis of electric charge conservation alone [1,2].

The connection between symmetry and mass of a field is not special to electromagnetism. A massless fermion field $\psi(x)$ exhibits chiral invariance, that is, its left– and right–handed components

$$\psi_L \equiv \frac{1-\gamma_5}{2}\psi\ ,\ \ \psi_R \equiv \frac{1+\gamma_5}{2}\psi \tag{1}$$

do not mix, and its Lagrangian

$$\mathcal{L}_{fermion} = \overline{\psi} i \not{D} \psi \tag{2}$$

is invariant under

$$\psi(x) \rightarrow e^{i\alpha_c(x)\gamma_5}\,\psi(x) \tag{3}$$

with the conserved current $J_\mu = \overline{\psi_R}\gamma_\mu\psi_R - \overline{\psi_L}\gamma_\mu\psi_L$. A finite mass for fermion explicitly breaks chiral invariance and spoils conservation of current: $\partial^\mu J_\mu = 2im_\psi\overline{\psi}\gamma_5\psi \neq 0$ [1,2].

The weak interactions, which govern several phenomena ranging from radioactivity to burning of sun, proceed with the exchange of massive vector bosons. The nature of this interaction can be revealed by examining the decay

$$n \rightarrow p + \left(W^- \rightarrow\right) e^- + \nu_e \tag{4}$$

which rests on the existence of a charged intermediate vector boson $W^\pm_\mu$. This vector particle must be massive for decay amplitude to reduce Fermi's empirical contact interaction

$$\frac{G_F}{\sqrt{2}}\,\overline{p_L}\,\gamma^\mu n_L\,\overline{\nu_e}\,\gamma_\mu e_L \tag{5}$$

at low energies. The correspondence requires $G_F/\sqrt{2} = g_W^2/(8M_W^2)$ where g_W is strength of $W^\pm$ coupling to charged currents, and M_W is its mass [1,2].

This brief discussion of weak interactions should convince oneself that massive vector bosons do exist in Nature and they must be incorporated in any theory of fundamental interactions. This is not straightforward, however, since, as mentioned in the case of electromagnetism, a hard mass for a vector boson explicitly breaks the gauge invariance. For a clearer view of this point it may be instructive to examine a generic Abelian vector field $A_\mu(x)$ (gauge field of some U(1) invariance [1,2]) which transforms as

$$A_\mu \rightarrow A_\mu - \frac{1}{g}\partial_\mu\alpha \tag{6}$$

where g is gauge coupling and $\alpha(x)$ is an arbitrary scalar field. If $A_\mu(x)$ is given a mass M_A then its Lagrangian

$$\mathcal{L} = -\frac{1}{4}F^{\mu\nu}F_{\mu\nu} + \frac{1}{2}M_A^2 A^\mu A_\mu\,, \tag{7}$$

with $F_{\mu\nu} \equiv \partial_\mu A_\nu - \partial_\nu A_\mu$ being the field strength tensor, transforms as

$$\mathcal{L} \to \mathcal{L} - \frac{1}{2g}M_A^2\left(2A^\mu\partial_\mu\alpha - \frac{1}{g}\partial^\mu\alpha\partial_\mu\alpha\right) \tag{8}$$

under (6). The excess terms, those proportional to M_A^2, in this expression show that vector boson mass explicitly breaks invariance under $U(1)$ symmetry.

This rather brief exercise shows how nontrivial it is to induce a mass for a vector boson. To have a clue of correct description it proves useful to focus on vector boson self energy

$$\Pi_{\mu\nu}\left(q^2\right) = \left(g_{\mu\nu} - \frac{q_\mu q_\nu}{q^2}\right)\Pi\left(q^2\right) \tag{9}$$

whose proportionality to the transverse projector $q^2 g_{\mu\nu} - q_\mu q_\nu$ follows from Ward identity $q^\mu \Pi_{\mu\nu}\left(q^2\right) = 0$ meaning current conservation. The scalar self energy $\Pi\left(q^2\right)$ is generated by interactions of A_μ with itself and other fields. One obvious contribution is provided by the vector boson mass term in (7):

$$\Pi_{\mu\nu}^{mass}\left(q^2\right) = iM_A^2 g_{\mu\nu} \tag{10}$$

It is clear that to generate part of the projector proportional to $q_\mu q_\nu/q^2$ we obviously need a strictly massless scalar field (to identify $1/q^2$ factor with its propagator) which couples to A_μ via its derivative (to generate $q_\mu q_\nu$ in the numerator). This observation is guided by interaction between $\alpha(x)$ and $A_\mu(x)$ given in (8). Naming this massless scalar as ϕ, letting it couple to A_μ via $M_A^2 A^\mu \partial_\mu \phi$, and interpreting M_A^2 at the vertex to be a background field one obtains

$$\Pi_{\mu\nu}^{\phi}\left(q^2\right) = -iM_A^2\frac{q_\mu q_\nu}{q^2} \tag{11}$$

which adds up to (10) to generate the requisite projector correctly. From these observations one concludes that a vector boson can acquire a mass if it couples to a strictly massless scalar field (See the book by Peskin and Schroeder listed in [1]).

The place where one can find requisite strictly massless scalar fields is given by Goldstone's theorem. This theorem states that there appear strictly massless scalar fields, the Goldstone bosons, for each broken symmetry generator when a continuous symmetry is spontaneously broken, that is, when the minimum energy configuration (ground state) does not respect the symmetry of the Lagrangian. Armed with this theorem, it is now clear that vector boson masses can be naturally generated in gauge theories with spontaneously broken symmetries. A toy model based on a complex scalar $\phi(x)$ and fermion field $\psi(x)$ that realizes aforementioned properties may be taken as

$$\mathcal{L} = -\frac{1}{4}F^{\mu\nu}F_{\mu\nu} + \overline{\psi_L}\left(\not{D} - m_\psi\right)\psi_L + \overline{\psi_R}\left(\not{D} - m_\psi\right)\psi_R + \left[h_\psi \overline{\psi_L}\phi\psi_R + h.c.\right] \\ + \left(D^\mu \phi\right)^\dagger D_\mu \phi - V\left(\phi^\dagger \phi\right) \tag{12}$$

where h_ψ is the Yukawa coupling and $V\left(\phi^\dagger\phi\right)$ is potential energy density of the scalar field. This Lagrangian represents an Abelian gauge theory invariant under

$$\phi(x) \rightarrow e^{iQ_\phi \alpha(x)}\phi(x),\ \psi_L(x) \rightarrow e^{iQ_{\psi_L}\alpha(x)}\psi_L(x),\ \psi_R(x) \rightarrow e^{iQ_{\psi_R}\alpha(x)}\psi_R(x) \tag{13}$$

accompanied by transformation of A_μ in (6). For the purpose of illustration the local U(1) invariance is chosen to treat left– and right–chirality fermions differently. The covariant derivatives $D_\mu\phi = \left(\partial_\mu + igQ_\phi A_\mu\right)\phi$, $D_\mu\psi_L = \left(\partial_\mu + igQ_{\psi_L}A_\mu\right)\psi_L$ and $D_\mu\psi_R = \left(\partial_\mu + igQ_{\psi_R}A_\mu\right)\psi_R$ are needed to make derivatives of fields transform as fields themselves under (13). In these expressions Q_ϕ, Q_{ψ_L} and Q_{ψ_R} are, respectively, charges of ϕ, ψ_L and ψ_R under local U(1) symmetry such that $Q_{\psi_R} + Q_\phi - Q_{\psi_L} = 0$ so as to enable the Yukawa interaction ([1,2]).

Focussing on a simple polynomial structure, potential of the scalar field may be taken as

$$V\left(\phi^\dagger\phi\right) = m^2\phi^\dagger\phi + \frac{\lambda}{2}\left(\phi^\dagger\phi\right)^2 \tag{14}$$

where $\lambda > 0$ and m^2 are model parameters. For $m^2 > 0$ this potential has a minimum at $\phi \equiv \phi^S_{min} = 0$, and an expansion in $\phi - \phi^S_{min}$ reveals a massive scalar interacting with massless A_μ and ψ.

When $m^2 < 0$, that is, when ϕ is a complex tachyon then potential (14) develops a minimum at $\phi \equiv \phi^B_{min}$ where

$$\phi^B_{min} = \left(-\frac{m^2}{\lambda}\right)^{1/2} \tag{15}$$

with an arbitrary phase reachable by a U(1) transformation. Expanding ϕ around this vacuum expectation value (VEV)

$$\phi(x) = \phi^B_{min} + \frac{1}{\sqrt{2}}\left(h(x) + ig(x)\right) \tag{16}$$

and using in (12) one finds that particle spectrum now consists of a massive vector boson A_μ, a massive real scalar $h(x)$ and a massive fermion $\psi(x)$ with the masses

$$M_A^2 = 2g^2 \left(\phi_{min}^B\right)^2,\ m_h^2 = 2\lambda \left(\phi_{min}^B\right)^2,\ m_\psi = h_\psi \phi_{min}^B \tag{17}$$

where real scalar $g(x)$ in (16) remains strictly massless, and gets swallowed by $A_\mu(x)$ to become massive. In other words, $g(x)$ can be absorbed in $A_\mu(x)$ by a U(1) gauge transformation

$$A_\mu(x) \to A_\mu(x) - \frac{1}{\sqrt{2}} \frac{\partial_\mu g(x)}{g\phi_{min}^B} \tag{18}$$

so that $g(x)$ disappears to generate the longitudinal component of the now-massive vector boson $A_\mu(x)$. This transformation is a reexpression of the fact that a massive vector field, with respect to a massless one such as electromagnetic field, is endowed with a third polarization component parallel to its propagation direction (momentum). (See, for instance, the review volume by Abers and Lee listed in [2]).

In the formulae above, ϕ_{min} stands for the field value for which potential is a minimum, and superscripts S and B designate, respectively, symmetric and broken phases of the U(1) gauge theory at hand. The minimum of the potential at ϕ_{min}^B corresponds to a specific choice for phase of ϕ and thus it breaks or better hides the U(1) invariance. It is in this minimum that gauge boson and fermion acquire their masses.

The toy U(1) gauge theory above illustrates some salient features of mass generation mechanism, that is, the Higgs mechanism. It is not a realistic model of weak interactions; however, it possesses all the features one needs to construct a realistic model. For future reference, one is reminded of the fact that the real scalar field $g(x)$ above is Goldstone boson (of spontaneously broken U(1) invariance), and $h(x)$ is a real field, the so-called Higgs field, that remains in the spectrum as a fingerprint of the Higgs mechanism [1, 2].

3 Electroweak Theory

Armed with the knowledge of Higgs mechanism gained in last section, one can now attempt at constructing a realistic model of weak interactions. A realistic model must, first of all, take into account the fact that gauge bosons to be given mass are electrically charged. Therefore, Abelian models like (12) do not work at all. In fact, a moment of thinking reveals that the simplest gauge group whose gauge bosons consist of a charged sector is SU(2) (the so-called isospin group). This point is best understood by giving the explicit expressions of SU(2) generators (the Pauli matrices)

$$\sigma_3 = \begin{pmatrix} 1 & 0 \\ 0 & -1 \end{pmatrix},\ \sigma_+ = \frac{1}{\sqrt{2}}(\sigma_1 + i\sigma_2) = \begin{pmatrix} 0 & \sqrt{2} \\ 0 & 0 \end{pmatrix},\ \sigma_- = \frac{1}{\sqrt{2}}(\sigma_1 - i\sigma_2) = \begin{pmatrix} 0 & 0 \\ \sqrt{2} & 0 \end{pmatrix} \tag{19}$$

with respective gauge bosons A_μ^3, W_μ^+ and W_μ^-. However, it is easy to see that an SU(2)–invariant theory alone is not sufficient for describing Nature.

The reason is that, upon spontaneous symmetry breaking, all these gauge bosons acquire masses, and thus, A^3_μ, the candidate for electromagnetic field, becomes massive. This is an unacceptable result. The way out is to introduce an additional gauge group the simplest of which being a U(1) factor so that gauge boson of this group combines with A^3_μ to form a massless state. The nature of this new U(1) invariance is determined by the Gell-Mann–Nishijima rule

$$Q = I_3 + \frac{Y}{2} \tag{20}$$

which relates electric charge Q to third component of isospin I_3 (eigenvalue of σ_3 above) and hypercharge Y of a particle. Therefore, correct gauge group which describes weak interactions and electromagnetism involves SU(2) and $\mathrm{U}(1)_Y$ symmetries.

At this stage problem is how to arrange matter into SU(2) and $\mathrm{U}(1)_Y$ gauge symmetries. The Yukawa interaction in (12) makes it clear that if both left– and right–handed matter fermions are charged under SU(2) then Higgs field ϕ must be a singlet under both SU(2) and $\mathrm{U}(1)_Y$. This, however, prevents $W^\pm$ bosons to acquire masses – *la raison de entre* for Higgs mechanism and SU(2) invariance. Hence, only matter fields of one handedness, say left-handed ones, must be charged under SU(2). The ones with opposite handedness must be singlets under SU(2) for Higgs field to be an SU(2) doublet to give masses to $W^\pm$ bosons. These observations lead one to the group structure G_{SM} =$\mathrm{SU}(2)_L \otimes \mathrm{U}(1)_Y$ meaning that left-handed fermions are SU(2) doublets, and right-handed ones are SU(2) singlets. In this context, gauge quantum numbers of leptons are

$$\begin{aligned} &L_1 = \begin{pmatrix} \nu_{e_L} \\ e_L \end{pmatrix}_{-1}, \; L_2 = \begin{pmatrix} \nu_{\mu_L} \\ \mu_L \end{pmatrix}_{-1}, \; L_3 = \begin{pmatrix} \nu_{\tau_L} \\ \tau_L \end{pmatrix}_{-1} \\ &e_{R_1} = (e_R)_{-2} \; , \; e_{R_2} = (\mu_R)_{-2} \; , \; e_{R_3} = (\tau_R)_{-2} \end{aligned} \tag{21}$$

and those of quarks are where subscript of each field refers to its hypercharge under the definition (20). On the other hand, charge assignments for quarks are given by

$$\begin{aligned} &Q_1 = \begin{pmatrix} u_L \\ d_L \end{pmatrix}_{1/3}, \; Q_2 = \begin{pmatrix} c_L \\ s_L \end{pmatrix}_{1/3}, \; Q_3 = \begin{pmatrix} t_L \\ b_L \end{pmatrix}_{1/3} \\ &u_{R_1} = (u_R)_{4/3} \; , \; u_{R_2} = (c_R)_{4/3} \; , \; u_{R_3} = (t_R)_{4/3} \\ &d_{R_1} = (d_R)_{-2/3} \; , \; d_{R_2} = (s_R)_{-2/3} \; , \; d_{R_3} = (b_R)_{-2/3} \end{aligned} \tag{22}$$

where subscript of each field refers to its hypercharge under the definition (20). The Higgs field H (which corresponds to the scalar field ϕ in (12) of the toy U(1) model discussed in last section) is given by

$$H = \begin{pmatrix} \varphi^+ \\ \varphi^0 \end{pmatrix}_1 \tag{23}$$

where superscripts on its component fields refer to their electric charges [1,2].

The $\mathrm{SU}(2)_L$ and $\mathrm{U}(1)_Y$ symmetries refer to electroweak sector which is the main topic of this review volume. Apart from these, as a theory of the constituents of hadrons, that is, quarks and gauge bosons gluing them together, one has to include an additional gauge invariance, $\mathrm{SU}(3)_c$, associated with color degrees of freedom. Quarks transform in fundamental of $\mathrm{SU}(3)_c$ whose gauge fields are gluons.

Then complete $\mathrm{SU}(3)_c \otimes \mathrm{SU}(2)_L \otimes \mathrm{U}(1)_Y$ invariant SM Lagrangian, explicating all gauge and flavor quantum numbers, is given by [3]

$$\begin{aligned}
\mathcal{L}_{\mathrm{SM}} = & -\frac{1}{4} G^{r_a}_{\mu\nu} G^{\mu\nu}_{r_a} - \frac{1}{4} W^{s_a}_{\mu\nu} W^{\mu\nu}_{s_a} - \frac{1}{4} B_{\mu\nu} B^{\mu\nu} \\
& + \overline{L_i}^{s_f} i \not{D} L_i^{s_f} + \overline{e}_{R_i} i \not{D} e_{R_i} + \overline{Q_i}^{r_f, s_f} i \not{D} Q_i^{r_f, s_f} + \overline{u}^{r_f}_{R_i} i \not{D} u^{r_f}_{R_i} + \overline{d}^{r_f}_{R_i} i \not{D} d^{r_f}_{R_i} \\
& - \Big[(h_u)_{ij} \overline{Q_i}^{r_f, s_f} H^{c, s_f} u^{r_f}_{R_j} + (h_d)_{ij} \overline{Q_i}^{r_f, s_f} H^{s_f} d^{r_f}_{R_j} \\
& + (h_e)_{ij} \overline{L_i}^{s_f} H^{s_f} e_{R_j} + h.c. \Big] \\
& + (D_\mu H)^\dagger D^\mu H - m^2 H^\dagger H - \frac{\lambda}{2} \left(H^\dagger H \right)^2
\end{aligned} \tag{24}$$

where $H^c \equiv i\sigma_2 H^\star$ is the charge-conjugate of H, and ranges of various indices are: $r_a = 1, \ldots, 8$ for adjoint of $\mathrm{SU}(3)_c$, $r_f = 1, \ldots, 3$ for fundamental of $\mathrm{SU}(3)_c$, $s_a = 1, \ldots, 3$ for adjoint of $\mathrm{SU}(2)_L$, $s_f = 1, 2$ for fundamental of $\mathrm{SU}(2)_L$, $i, j = 1, \ldots, 3$ for three flavors of quarks and leptons. The Yukawa matrices are 3×3 complex matrices in the space of fermion flavors. For a generic matter field, including Higgs field H, the gauge-covariant derivative is defined by

$$D_\mu \psi = \left(\partial_\mu - \frac{i}{2} g_s \lambda_{r_a} g^{r_a}_\mu - \frac{i}{2} g_w \sigma_{s_a} W^{s_a}_\mu - \frac{i}{2} g_Y Y_q B_\mu \right) \psi \tag{25}$$

where g^a_μ, W^s_μ and B_μ are gauge fields, $\lambda^{r_a}/2$ (3×3 Gell-Mann matrices), $\sigma^{s_a}/2$ (2×2 Pauli matrices), and Y_ψ (the hypercharge) are generators, and g_s, g_w and g_Y are gauge couplings of $\mathrm{SU}(3)_c$, $\mathrm{SU}(2)_L$ and $\mathrm{U}(1)_Y$, respectively. The field-strength tensors of the gauge fields are given by

$$\begin{aligned}
G^{r_a}_{\mu\nu} &= \partial_\mu g^{r_a}_\nu - \partial_\nu g^{r_a}_\mu + g_s\, f^{r_a bc} g^b_\mu g^c_\nu \\
W^{s_a}_{\mu\nu} &= \partial_\mu W^{s_a}_\nu - \partial_\nu W^{s_a}_\mu + g_w\, \epsilon^{s_a bc} W^b_\mu W^c_\nu \\
B_{\mu\nu} &= \partial_\mu B_\nu - \partial_\nu B_\mu
\end{aligned} \tag{26}$$

where f^{abc} and ϵ^{abc} are, respectively, the structure constants of $\mathrm{SU}(3)_c$ and $\mathrm{SU}(2)_L$.

The color group $\mathrm{SU}(3)_c$ is never broken; quantum chromodynamic (QCD) interactions are (non-)perturbative at (low)high energies – the infamous asymptotic freedom of color group. It is the non-perturbativity of $\mathrm{SU}(3)_c$ at the infrared that is responsible for confinement of quarks to form hadrons [4].

As was illustrated for toy U(1) gauge theory in the last section, upon spontaneous breakdown of $SU(2)_L \otimes U(1)_Y$ by the Higgs sector, all fermions acquire masses with a hierarchy determined by their Yukawa matrices:

$$m_u = h_u \frac{v}{\sqrt{2}}\,,\; m_d = h_d \frac{v}{\sqrt{2}}\,,\; m_e = h_e \frac{v}{\sqrt{2}}. \tag{27}$$

where

$$v = \left(-\frac{m^2}{\lambda}\right)^{1/2} \tag{28}$$

is the Higgs VEV. The physical fermion masses arise after diagonalization of these mass matrices. This procedure has no remnant in lepton sector (the flavor structure of h_e is unphysical as long as neutrinos are taken to be massless). However, quark mass matrices, upon diagonalization, leave a physical finger print by dressing all $W^{\pm}\overline{u}^{phys}_{L_i} d^{phys}_{L_j}$ vertices by $(V_{CKM})_{ij}$, the (i,j) entry of the Cabibbo–Kobayashi–Maskawa matrix

$$V_{CKM} = (V_L^u)^{\dagger}\, V_L^d \tag{29}$$

where $V_L^u h_u h_u^{\dagger} (V_L^u)^{\dagger} = \mathbf{1}$ and $V_L^d h_d h_d^{\dagger} \left(V_L^d\right)^{\dagger} = \mathbf{1}$ with unitary transformations of the quark fields $u_L^{phys} = V_L^u u$ and $d_L^{phys} = V_L^d d_L$. The unitary matrices $V_R^{u,d}$ that define physical right-handed quarks do not appear in interactions of mass eigenstate fields.

Like fermions, gauge bosons acquire their masses upon $SU(2)_L \otimes U(1)_Y$ breaking. The physical gauge bosons

$$W^{\pm} = \frac{1}{\sqrt{2}}(A_{\mu}^1 \mp i A_{\mu}^2)\,,\; Z_{\mu} = \frac{g_w A_{\mu}^3 - g_Y B_{\mu}}{\sqrt{g_w^2 + g_Y^2}}\,,\; A_{\mu} = \frac{g_w W_{\mu}^3 + g_Y B_{\mu}}{\sqrt{g_w^2 + g_Y^2}} \tag{30}$$

obtain masses $M_W = \frac{1}{2} g_w v$, $M_Z = \frac{1}{2}\sqrt{g_w^2 + g_Y^2}\, v$ and $M_A = 0$. The model correctly predicts massive charged bosons $W^{\pm}$ plus a massless gauge boson, A_{μ}, corresponding to electromagnetic field. However, it also predicts a neutral massive boson Z_{μ} experimental confirmation of which forms an important test of electroweak theory. At this point it may be instructive to determine cost of giving mass to gauge bosons. Before electroweak breaking (corresponding to ϕ^S_{min} in last section), all gauge bosons W_{μ}^i and B_{μ} are massless, and the Higgs doublet (23) consists of four real fields $\varphi^{\pm}$, $\Re[\varphi^0 - v]$ and $\Im[\varphi^0 - v]$. In course of spontaneous breakdown of $SU(2)_L \otimes U(1)_Y$ (corresponding to ϕ^B_{min} of last section), $\varphi^{\pm}$ are spent for giving mass to $W^{\pm}$ bosons, and $\Im[\varphi^0 - v]$ is spent for giving mass to Z boson. Hence, three Goldstone bosons are spent for giving mass to gauge bosons. The real part of neutral component of Higgs field, $\Re[\varphi^0 - v]$, is nothing but the Higgs boson – an important finger print of spontaneous breakdown of $SU(2)_L \otimes U(1)_Y$. Its mass-squared

$$m_h^2 = 2\lambda v^2 = -2m^2 \tag{31}$$

is a direct probe of the tachyonic mass of the Higgs field H [2].

Since their first observation at CERN, the $W^{\pm}$ and Z boson masses have been measured with increasing precision:

$$M_W = 80.403\,\text{GeV}\,,\ M_Z = 91.1876\,\text{GeV} \tag{32}$$

as listed by PDG [5]. Continuing with PDG tables, the fermion masses read as

$$\begin{aligned} &m_e = 0.511 \text{ MeV}\,,\ m_\mu = 105.658 \text{ MeV}\,,\ m_\tau = 1.777 \text{ GeV} \\ &m_u = 1.5 - 3 \text{ MeV}\,,\ m_c = 1.25 \text{ GeV}\,,\ m_t = 174.2 \text{ GeV} \\ &m_d = 3 - 7 \text{ MeV}\,,\ m_s = 95 \text{ MeV}\,,\ m_b = 4.2 \text{ GeV} \end{aligned} \tag{33}$$

where u, d, s quark masses are $\overline{\text{MS}}$ masses evaluated at 2 GeV scale. The quark mixings, parameterized by V_{CKM}, have been measured and will continue to be measured with increasing precision via experiments on various B meson decays (the b quark is heavy enough to apply perturbative QCD).

Consequently, in the SM the only particle which has not yet been observed is the Higgs boson h, and its experimental discovery and theoretical understanding is one of the hottest topics of high energy physics today [1, 2, 6]. From theoretical perspective, the model itself does not provide a means of computing m_h in terms of known particle masses. There are certain physical requirements, however, which determine likely ranges of m_h. An important constraint in the class is unitarity which requires cross section for a scattering process to behave as $1/\text{energy}^2$ at high energies. The importance of unitarity constraint is best seen by examining scattering of massive vector bosons. Indeed, scattering of longitudinal $W^{\pm}/Z$ (massive vector bosons, with respect to photon, develop a third polarization component parallel to their momentum) grow with their momenta. More explicitly, for $s \gg M_W^2$,

$$\mathcal{A}\left(W_a^+ W_b^- \to W_c^+ W_d^-\right) = \frac{1}{v^2}\left(s + t - \frac{s^2}{s - m_h^2} - \frac{t^2}{t - m_h^2}\right) \tag{34}$$

where Mandelstam variables are defined by $s = (p_a + p_b)^2$ and $t = (p_a - p_c)^2$. Clearly, this scattering amplitude must approach to that of charged Goldstone bosons contained in (23) at high energies, and hence, it can cause violation of unitarity. This yields an upper bound of $m_h \lesssim 900$ GeV. One arrives at similar bounds via triviality (related to renormalization group flow of the quartic coupling in Higgs potential) or vacuum stability arguments. Hence, one expects Higgs boson to weigh below TeV scale for Higgs mechanism to account for observed mass spectra of matter and force carriers. In fact, global fits to various electroweak precision observables, including gauge boson masses and several asymmetries, agree with the SM predictions. These fits require Higgs boson to weigh below 260 GeV at 95% C. L. [2].

Higgs searches at various experiments have given negative results. The LEP II experiment at CERN, e^+e^- collisions at $\sqrt{s} = 209$ GeV, has ended by reporting the lower bound [7]

$$m_h \geq 114.4\,\text{GeV} \tag{35}$$

at 90% C. L. Search for Higgs boson will continue at next generation colliders: the Large Hadron Collider (LHC) at CERN, the International Linear Collider (ILC), CLIC, and others. In each experiment, collisions of two beams of hadrons/leptons generates Higgs boson in certain channels. The question of what channel is more viable for Higgs discovery depends exclusively on the Higgs mass since couplings of the Higgs boson to itself and other matter species are determined entirely by its mass alone [6]:

$$g_{hff} = \frac{m_f}{v}\ ,\ g_{hVV} = 2\frac{M_V^2}{v}\ ,\ g_{hhVV} = 2\frac{M_V^2}{v^2}\ ,\ g_{hhh} = 3\frac{m_h^2}{v}\ ,\ g_{hhhh} = 3\frac{m_h^2}{v^2} \tag{36}$$

where $V = (W^{\pm}, Z)$ and f is a generic fermion. The relations among certain couplings reveal hidden $\mathrm{S}(2)_L \otimes \mathrm{U}(1)_Y$ symmetry. The SM Lagrangian (24), being a renormalizable quantum field theory, has all the ingredients for inclusion of quantum theoretic effects in Higgs (and any other field in the spectrum) boson production and decays. These corrections, which are crucial for high-energy high-intensity collider processes, will not be detailed in this review volume [2].

4 Collider Searches for Higgs Boson

In this section we shall discuss, somewhat briefly, Higgs boson search strategies at high-energy colliders, in particular, LHC. (See the excellent book [6] for a detailed analysis of Higgs boson search.) The experiments at other colliders, especially at ILC as a precision machine, are certainly important; however, as the earliest experiment to start, we focus on the LHC signals here.

The LHC experiment (see the URL: `http://greybook.cern.ch/` as well as [5]), which is expected to start running this year, is a pp collider where each proton beam has 7 TeV energy. As for any collision process, a crucial parameter is luminosity L which relates rate of occurrence, R, of a given event to its interaction cross section σ_{int} via $R = L\sigma_{int}$. The luminosity is a measure of how frequent the beams of particles are showering on the interaction point:

$$L = \frac{\nu}{4\pi}\frac{n_1 n_2}{\sigma_{hor}\sigma_{ver}} \tag{37}$$

where n_1 and n_2 are numbers of particles in bunches (forming the beam), ν frequency with which the bunches collide, σ_{hor} and σ_{ver} characterize Gaussian beam spreads in transverse plane (in horizontal and vertical directions). The accelerator parameters at the LHC are such that planned luminosity is $10^{34}\ \mathrm{cm}^{-2}\mathrm{s}^{-1}$.

The LHC complex consists of four detectors: ATLAS and CMS (general purpose experiments for new physics search), ALICE (a special experiment for quark-gluon plasma search), and LHCb (a special experiment for measurements of flavor and CP violation effects in B meson system). There are other experiments as well: the LHCf experiment which will examine high-rapidity region for laboratory generation of ultra-high energy cosmic rays, and TOTEM experiment which will measure cross sections and diffractive processes in conjunction with the CMS detector.

In pp collisions, production and detection of the Higgs signal proceed via different channels at different Higgs mass values and luminosities. The fundamentally relevant processes are

$$\begin{aligned}
&\text{Gluon fusion: } gg \rightarrow h \rightarrow \gamma\gamma\,, VV^{\star} \\
&\text{Weak boson fusion: } qq \rightarrow qqV^{\star}V^{\star} \rightarrow qqh \\
&\text{Drell-Yan: } q\overline{q} \rightarrow hV^{\star} \\
&\text{Radiation off top: } gg, q\overline{q} \rightarrow ht\overline{t}
\end{aligned} \tag{38}$$

each of which can be estimated by using (36) and contrasted with experimental data available then. Clearly, gluon fusion $g\,g \rightarrow h$ and diphoton decay $h \rightarrow \gamma\gamma$ arise at the loop level. In all these channels, except for gluon fusion, detection of the Higgs boson requires reconstruction of its mass and couplings from decay products.

Figure 1 depicts integrated luminosity needed to discover a Higgs boson of given mass at the LHC. The figure treats production channels one by one, and it provides requisite information for determining what channel dominates at what values of Higgs boson masses for an anticipated integrated luminosity in units of fb^{-1}.

From the figure it is clear that an integrated luminosity of 100 fb^{-1} suffices to discover a Higgs boson as heavy as 600 GeV (which is well inside the range expected from unitarity and other arguments). In case LHC achieves this size of luminosity it will certainly discover a Higgs boson within 5 σ accuracy. The largest contribution among all channels comes from $h \rightarrow Z^{\star}Z^{\star} \rightarrow 4\ell$ for $m_h \sim 2M_Z$. For small m_h, close to LEP II lower bound, a 30 fb^{-1} luminosity is sufficient to detect Higgs boson via gluon fusion or diphoton decay or $WW^{\star}$ production. For high-mass Higgs, $m_h \gtrsim 300$ GeV, $h \rightarrow ZZ$ dominates. Irrespective of what channel is actually dominating, LHC will certainly discover a Higgs boson within 5 σ accuracy if electroweak breaking is the correct mechanism for generating mass of matter.

5 Summary and Outlook

In this brief review we have provided a pedagogical introduction into 'the problem of mass generation' by putting emphasis on connection between mass and symmetry properties of a field. Then we have proceeded to explain basic mechanism – the Higgs mechanism – by a toy U(1) gauge theory. Afterwards, we have mentioned all the reasonings and constraints for constructing a realistic theory of weak interactions. This led us to a theory of electroweak interactions. Finally, we have given a brief discussion of LHC as a collider and Higgs

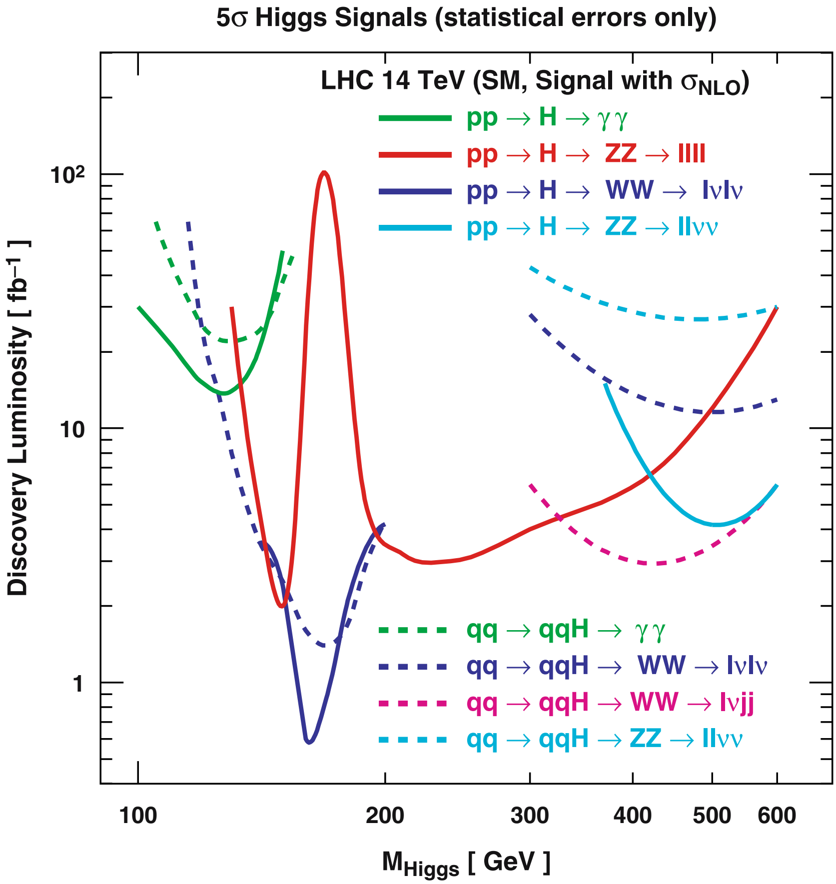

Fig. 1. 5σ discovery luminosity vs. Higgs boson mass at the LHC. The figure is taken from [8].

boson searches therein. The LHC, if reaches a luminosity $\mathcal{O}(100 \text{ fb}^{-1})$ will be capable of discovering a Higgs boson within the mass range indicated by unitarity arguments.

The SM, however, hardly stands as a complete theory of Nature above the electroweak scale. The reason is that the Higgs sector develops a strong sensitivity to ultraviolet (UV) cut-off once effects of quantum fluctuations are included. More explicitly,

$$m_h^{2\,corr} = m_h^2 + \frac{3\Lambda^2}{8\pi v^2}\left[m_h^2 + 2M_W^2 + M_Z^2 - 4m_t^2\right] \tag{39}$$

where Λ is the UV cut-off, m_h is tree-level (classical) Higgs mass defined in (31), and $m_h^{2\,corr}$ is the one-loop (quantal) Higgs boson mass. This divergence is quadratic, not logarithmic, and hence it poses a serious problem since larger the Λ (the highest energy scale up to which SM works well) larger the quantal correction. For instance, if SM is a theory of fundamental particles and interactions among them up to the gravitational scale $\Lambda = (8\pi G_N)^{-1/2}$ then quantum correction turns out to be some 16 orders of magnitude larger than the tree-level Higgs mass m_h (which is expected to lie at the electroweak scale). One way of evading this is to fine-tune quantum correction to zero which requires $m_h \sim 320$ GeV. This requirement involves a huge fine-tuning, however. A more natural way out of this impasse should soften dependence on Λ or should give a physical meaning to Λ beyond what one might know of in the SM.

The way out is to embed SM into another theoretical framework which exhibits good UV behavior. One such playground is provided by 'supersymmetry' which controls Higgs mass just like chiral symmetry controls the fermion masses. It doubles the particle spectrum of SM by introducing a fermion for each boson and vice versa so that loops of a particle and its supersymmetry partner cancel out in accord with Spin-Statistics Theorem (for a pedagogical introduction to supersymmetry, see [9]).

Other than supersymmetry, a higher dimensional spacetime whose gravitational interactions possess a Newton constant $\mathcal{O}$(TeV) is known to bring an alternative solution to UV sickness of the SM Higgs sector. This framework, large extra dimensions, promotes electroweak scale TeV to be a fundamental constant of Nature, and treats Newton constant in four dimensions as a derived quantity (for a pedagogical introduction to extra dimensions, see [10]).

A beginner having sufficient 'theory minimum' about quantum field theory is expected to follow observations and arguments given in this review volume.

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References

1. T.P. Cheng and L.F. Li, *Gauge theory of Elementary Particle Physics*, Clarendon Press, 1984; C. Itzykson and J.B. Zuber, *Quantum Field Theory*, Mac Graw Hill, 1985; D. Bailin and A. Love, *Introduction to Gauge Field Theory*, Adam Hilger, 1986; M.E. Peskin and D.V. Schroeder, *An introduction to Quantum Field Theory*, Addison-Wesley, 1995.

2. E. Abers and B. Lee, Phys. Rep. 9 (1973) 1; M.A.B. Beg and A. Sirlin, Phys. Rept. 88 (1982) 1; M.E. Peskin, hep-ph/9705479; S. Dawson, hep-ph/9901280; M. Spira and P.M. Zerwas, hep-ph/9803257; C. Quigg, hep-ph/9905369; M. Carena and H.E. Haber, Prog. Part. Nucl. Phys. **50**, 63 (2003) [arXiv:hep-ph/0208209]; A. Djouadi, arXiv:hep-ph/0503172.
3. S. Glashow, Nucl. Phys. 22 (1961) 579; S. Weinberg, Phys. Rev. Lett. 19 (1967) 1264; A. Salam, in "Elementary Particle Theory", ed. N. Svartholm, Almqvist and Wiksells, Stockholm (1969) p. 367.
4. W.J. Marciano and H. Pagels, Phys. Rept. **36**, 137 (1978); V.A. Novikov, L.B. Okun, M.A. Shifman, A.I. Vainshtein, M.B. Voloshin and V.I. Zakharov, Phys. Rept. **41**, 1 (1978); E. Reya, Phys. Rept. **69**, 195 (1981).
5. Particle Data Group, K. Hagiwara *et al.*, Phys. Rev. D66 (2002) 010001; S. Eidelman *et al.*, Phys. Let. B592 (2004) 1.
6. J.F. Gunion, H.E. Haber, G.L. Kane and S. Dawson, *The Higgs Hunter's Guide*, Westview Press; New Ed. 2001.
7. The LEP Collaborations and LEP Higgs Working Group, ALEPH-2000-074 CONF 200-051.
8. M. Dittmar and S. Nicollerat, CMS–Note 2001-036 (See the URL: `http://cmsdoc.cern.ch/documents/01/note01-036.pdf`).
9. H.P. Nilles, Phys. Rept. **110**, 1 (1984); M. E. Peskin, *Supersymmetry: Theory*, World Scientific, 1996; S. P. Martin, arXiv:hep-ph/9709356.
10. N. Arkani-Hamed, Talk at SLAC Summer Institute, 1998 (See the URL: `http://www.slac.stanford.edu/gen/meeting/ssi/1998/arkani.html`); C. Csaki, J. Hubisz and P. Meade, arXiv:hep-ph/0510275.