8 The broken $\mathbf{SU(2)_L \otimes U(1)_Y}$ symmetry

In this case, the generators of the symmetry group will be $T^J = (\tau_1, \tau_2, \tau_3, Y)$, *i.e.* the generators of the weak isospin group and of the $U(1)_Y$ hypercharge gauge group. We introduce a complex scalar field $\Phi(x)$, which is a doublet of $SU(2)$ ($I_\Phi = \frac{1}{2}$),

$$\Phi = \frac{1}{\sqrt{2}} \begin{pmatrix} \varphi_1 & - & i\,\varphi_2 \\ \varphi_3 & - & i\,\varphi_4 \end{pmatrix} \tag{8.1}$$

and the standard scalar lagrangian

$$\mathcal{L}_S = \partial_\mu \Phi^\dagger \partial^\mu \Phi - V(\Phi), \quad V(\Phi) = -\mu^2 \Phi^\dagger \Phi \; + \; h\; (\Phi^\dagger \Phi)^2. \tag{8.2}$$

which is invariant under the rigid transformation

$$\Phi \to \Phi' = \mathrm{e}^{i\boldsymbol{\tau}\cdot\boldsymbol{\alpha}/2}\; \mathrm{e}^{iy_\Phi \beta/2}\; \Phi. \tag{8.3}$$

The minimum of the potential is obtained for (see eq. (6.4))

$$\Phi^\dagger \Phi = |\Phi|^2 = \frac{\mu^2}{2h} = \frac{v^2}{2}. \tag{8.4}$$

There is an infinite number of vacua states : all states with the norm $v/\sqrt{2}$ obtained by a gauge transformation. We choose the physical vacuum to be

$$\boxed{\Phi_0 = \begin{pmatrix} 0 \\ \frac{v}{\sqrt{2}} \end{pmatrix} \quad \text{with} \;\; v = \frac{\mu}{\sqrt{h}}.} \tag{8.5}$$

Since we require the electric charge to be conserved after symmetry breaking, following the reasoning in sec. 6.2, we have to enforce that the charge generator acting on the vacuum state should vanish. Using the Gell-Mann/Nishijima relation eq. (4.28) the charge operator acting on Φ_0 is

$$Q\;\Phi_0 = (I_3 + \frac{Y}{2})\;\Phi_0 = \frac{1}{2}\;(\tau_3 + Y)\;\Phi_0 = \begin{pmatrix} \frac{1}{2} + \frac{y_\Phi}{2} & 0 \\ 0 & -\frac{1}{2} + \frac{y_\Phi}{2} \end{pmatrix} \begin{pmatrix} 0 \\ \frac{v}{\sqrt{2}} \end{pmatrix} \;=\; 0 \tag{8.6}$$

implying that the hypercharge of the scalar field must be $y_\Phi = 1$ to ensure charge conservation in the broken theory: the charge of the classical vacuum is 0. As in the abelian case, we can study the system around the classical minimum and expand the scalar field around its vacuum expectation value

$$\Phi = \begin{pmatrix} \frac{1}{\sqrt{2}}(\omega_1(x) - i\omega_2(x)) \\ \frac{1}{\sqrt{2}}(v + \omega_0(x) - i\omega_3(x)) \end{pmatrix} \;=\; \begin{pmatrix} \omega^*(x) \\ \frac{1}{\sqrt{2}}(v + \omega_0(x) - i\omega_3(x)) \end{pmatrix}. \tag{8.7}$$

The complex field ω^* has a positive electric charge while ω_0 and ω_3 are neutral. In terms of the new variables the scalar potential $V(\Phi)$ becomes

$$V(\Phi) = hv^2\omega_0^2 \ + \ hv \ \omega_0(\omega_0^2 + \boldsymbol{\omega}^2) \ + \ \frac{h}{4} \ (\omega_0^2 + \boldsymbol{\omega}^2)^2 \tag{8.8}$$

showing that the triplet $\boldsymbol{\omega}$ of ω_i fields is massless while the neutral ω_0 field acquires a mass

$$\boxed{M_{\omega_0} = \sqrt{2hv^2}.} \tag{8.9}$$

All these fields are coupled together with a strength which can be read off the equation above.

Thus, in our model, in agreement with Noether theorem, three degrees of freedom are broken leading to three massless Goldstone bosons, and the vacuum is still left invariant under the combination $Q = I_3 + Y/2$. There is still an abelian symmetry left, namely the $U(1)_{\rm emg}$ group.

8.1 Local symmetry breaking and the Brout-Englert-Higgs mechanism

Armed with this lengthy preliminaries we now turn to spontaneous breaking of the local gauge symmetry $SU(2)_L \otimes U(1)_Y$ down to $U(1)_{\rm emg}$ in the framework of the Standard Model. Let us state the results before diving into an ocean of technicalities. The case of a global symmetry has just been analysed and led to the appearance of three massless (Goldstone) bosons and a massive one. When the symmetry is made local these massless bosons turn out to be unphysical (two charged ones, ω and ω^*, and a neutral one ω_3), in the sense that they can be gotten rid off by a gauge transformation, but instead, three gauge bosons (a neutral one and the two charged ones) become massive and therefore acquire longitudinal polarisation states which are the Goldstone modes in disguise.

To implement the breaking of the local $SU(2)_L \otimes U(1)_Y$ symmetry we first have to extend the electroweak lagrangian eq. (5.22) to include the scalar field contribution $\mathcal{L}_S$ eq. (8.2) in its locally gauge invariant form (see eq. (8.12) below) as well as the interaction of the scalar field with the fermions $\mathcal{L}_Y$ (where Y stands for Yukawa; see eq. (8.26) below) so that the complete electroweak lagrangian density is

$$\mathcal{L} \ = \ \mathcal{L}_F \ + \ \mathcal{L}_G \ + \ \mathcal{L}_S \ + \ \mathcal{L}_Y. \tag{8.10}$$

In the following we work in the unitary gauge.

8.2 The Higgs and gauge bosons sector : masses and couplings

We concentrate for the moment on $\mathcal{L}_S$ which drives the spontaneous breaking of the local electroweak symmetry. Only neutral scalar fields can acquire a vacuum expectation value : other fields, such as fermions or gauge bosons, cannot do so otherwise the physical vacuum would have some angular momentum or other non-vanishing quantum numbers. We impose now the invariance of $\mathcal{L}_S$ under a change of the local phases

$$\Phi(x) \to \Phi'(x) = \mathrm{e}^{ig\boldsymbol{\alpha}(x)\cdot\boldsymbol{\tau}/2}\ \mathrm{e}^{ig' y_\Phi \beta(x)/2}\ \Phi(x). \tag{8.11}$$

To keep gauge invariance requires substituting the covariant derivative to the partial derivative in $\mathcal{L}_S$ which then takes the form

$$\boxed{\mathcal{L}_S = D_\mu\Phi^\dagger D_\mu\Phi - \mu^2\Phi^\dagger\Phi + h(\Phi^\dagger\Phi)^2} \tag{8.12}$$

with the definition, eq. (5.20),

$$D_\mu = \partial^\mu - i\ g\ \frac{\boldsymbol{\tau}}{2}\cdot\mathbf{W}^\mu - i\ \frac{1}{2}\ g'\ B^\mu \tag{8.13}$$

and the choice, eq. (8.6), $y_\Phi = 1$ for the hypercharge. This can be easily checked using the same line of reasoning as used in sec. 5.

To study the system around the classical vacuum we parameterise the scalar field as in eq. (8.7). However we note that by an appropriate gauge transformation we can find $\boldsymbol{\alpha}(x)$, $\beta(x)$ such that :

$$\mathrm{e}^{ig' y_\Phi \beta(x)/2}\ \mathrm{e}^{ig\boldsymbol{\tau}\cdot\boldsymbol{\alpha}(x)/2}\ \Phi(x) = \begin{pmatrix} 0 \\ \frac{v+H(x)}{\sqrt{2}} \end{pmatrix}, \tag{8.14}$$

showing that the fields $\omega_i(x)$ can be removed from the lagrangian altogether and therefore are not physical. Of course, explicit gauge invariance of the vacuum state will be lost since a particular gauge has been chosen. To analyse the effects of symmetry breaking we work with the "physical" A_μ and Z_μ fields of eq. (5.31) rather than with $W_{3\mu}$ and B_μ. For this purpose we use the expression eq. (5.43) for the covariant derivative which, applied to the form eq. (8.14) of Φ (with $e_1 = 1, e_2 = 0$), yields

$$\begin{aligned} D_\mu \begin{pmatrix} 0 \\ \frac{v+H(x)}{\sqrt{2}} \end{pmatrix} &= \Bigg[\partial_\mu - i\frac{e}{\sqrt{2}\sin\theta_W}\begin{pmatrix} 0 & W_\mu^* \\ W_\mu & 0 \end{pmatrix} - ie\begin{pmatrix} A_\mu & 0 \\ 0 & 0 \end{pmatrix} \\ &\quad - i\frac{e}{\sin\theta_W\cos\theta_W}\begin{pmatrix} \frac{1}{2}-\sin^2\theta_W Z_\mu & 0 \\ 0 & -\frac{1}{2}Z_\mu \end{pmatrix}\Bigg]\begin{pmatrix} 0 \\ \frac{v+H(x)}{\sqrt{2}} \end{pmatrix} \\ &= \begin{pmatrix} -i\frac{e}{2\sin\theta_W}W_\mu^*(v+H(x)) \\ \partial_\mu\frac{H(x)}{\sqrt{2}} + i\frac{e}{2\sqrt{2}\sin\theta_W\cos\theta_W}Z_\mu(v+H(x)) \end{pmatrix} \end{aligned} \tag{8.15}$$

It is then trivial to get $D_\mu\Phi^\dagger D_\mu\Phi$ and write the scalar lagrangian density $\mathcal{L}_S$

$$\boxed{\begin{aligned} \mathcal{L}_S =& \frac{1}{2}(\partial_\mu H(x))^2 + \frac{e^2}{4\sin^2\theta_W}(v+H(x))^2 W_\mu^* W^\mu\ +\ \frac{e^2}{8\sin^2\theta_W\cos^2\theta_W}(v+H(x))^2 Z_\mu Z^\mu \\ & - hv^2H^2\ -\ hv\ H^3\ -\ \frac{h}{4}\ H^4 \end{aligned}} \tag{8.16}$$

In the last line we have used eq. (8.8) for the scalar potential dropping of course the spurious $\boldsymbol{\omega}(x)$ fields which have been gauged away. The above equation contains a lot of information since it gives masses to the gauge and the Higgs fields as well as defines the couplings between them.

• **Masses**

Combining the terms proportional to v^2 in the equation above with the stress-energy terms of $\mathcal{L}_{0_G}$, eq. (5.48), we have the pieces in the lagrangien density which lead to the free propagators of the H and gauge bosons,

$$\begin{aligned}\mathcal{L}_{OS} + \mathcal{L}_{OG} &= \frac{1}{2}(\partial_\mu H(x))^2 - hv^2 H^2 - \frac{1}{4}\mathcal{K}_{A\mu\nu}\mathcal{K}_A^{\mu\nu} \\ &-\frac{1}{2}\mathcal{K}^*_{\mu\nu}\mathcal{K}^{\mu\nu} + \frac{e^2v^2}{4\sin^2\theta_W}W^*_\mu W^\mu \\ &-\frac{1}{4}\mathcal{K}_{Z\mu\nu}\mathcal{K}_Z^{\mu\nu} + \frac{e^2v^2}{8\sin^2\theta_W\cos^2\theta_W}Z_\mu Z^\mu \end{aligned} \tag{8.17}$$

Using the same method as in sec. 7.1 we can derive the propagators of the H scalar and the gauge bosons

$$\begin{aligned} G(k) &= \frac{i}{k^2 - M_H^2 + i\epsilon} \\ G_A^{\mu\nu}(k) &= \frac{-i}{k^2 + i\epsilon}\, g^{\mu\nu} \\ G_W^{\mu\nu}(k) &= \frac{-i}{k^2 - M_W^2 + i\epsilon}(g^{\mu\nu} - k^\mu k^\nu/M_W^2) \\ G_Z^{\mu\nu}(k) &= \frac{-i}{k^2 - M_Z^2 + i\epsilon}(g^{\mu\nu} - k^\mu k^\nu/M_Z^2) \end{aligned} \tag{8.18}$$

We recover a massive H field with $\boxed{M_H = \sqrt{2h}\, v}$ as in eq. (7.8), while the W and Z bosons acquire the masses

$$\boxed{M_W = \frac{e\,v}{2\sin\theta_W}, \qquad M_Z = \frac{e\,v}{2\sin\theta_W\cos\theta_W},} \tag{8.19}$$

and the photon remains massless as no quadratic term in A_μ appears in the lagrangian. The vanishing of the photon mass is a consequence of the surviving exact gauge symmetry $U(1)_{\text{emg}}$. Note the important relation

$$\boxed{M_W = M_Z\cos\theta_W} \tag{8.20}$$

We have the relation $v = \sin\theta_W M_W/\sqrt{\pi\alpha}$ between the vacuum expectation value of the scalar field and the physical parameters and, plugging in numerical values, we find $v \sim 250$ GeV, which is the basis for the claim, made in the introduction, that the non-abelian symmetry is broken at the scale of 250 GeV.

• Couplings

We consider now all the terms of $\mathcal{L}_S$, eq. (8.16), not contained in $\mathcal{L}_{0S}$ to define the interaction lagrangian of the Higgs boson

$$\begin{aligned}\mathcal{L}_{IS} \;=\; & \frac{e^2 v}{2\sin^2\theta_W} H W^*_\mu W^\mu + \frac{e^2}{4\sin^2\theta_W} H^2 W^*_\mu W^\mu \\ & + \frac{e^2 v}{4\sin^2\theta_W\cos^2\theta_W} H Z_\mu Z^\mu + \frac{e^2}{8\sin^2\theta_W\cos^2\theta_W} H^2 Z_\mu Z^\mu \\ & - hv\, H^3 - \frac{h}{4} H^4 \end{aligned} \tag{8.21}$$

One notes that the trilinear couplings of the H boson to a pair of gauge bosons have the dimension of a mass, proportional to the vacuum expectation value v, while the quadrilinear couplings are dimensionless proportional to e^2. One can show that, in terms of Feynman diagrams,

$$\begin{aligned} &- \text{ the vertex } HW^+W^- \text{ is : } -i\frac{e^2 v}{2\sin^2\theta_W} = -i\frac{e}{\sin\theta_W} M_W; \\ &- \text{ the vertex } HZZ \text{ is : } -i\frac{e^2 v}{2\sin^2\theta_W\cos^2\theta_W} = -i\frac{e}{\sin\theta_W\cos\theta_W} M_Z; \\ &- \text{ the vertex } H^2W^+W^- \text{ is : } -i\frac{e^2}{2\sin^2\theta_W} \\ &- \text{ the vertex } H^2ZZ \text{ is : } -i\frac{e^2}{2\sin^2\theta_W\cos^2\theta_W}. \end{aligned} \tag{8.22}$$

There are furthermore the H boson self-couplings proportional respectively to hv and h. These variables are easily eliminated in favour of the observables M_W, M_H and one finds,

$$\begin{aligned} &- \text{ the vertex } H^3\text{: } i\,6hv = i\,\tfrac{3}{2}\,\tfrac{e}{\sin\theta_W}\,\tfrac{M_H^2}{M_W} \\ &- \text{ the vertex } H^4\text{: } i\,6h = i\,\tfrac{3}{4}\,\tfrac{e^2}{\sin^2\theta_W}\,\tfrac{M_H^2}{M_W^2}\;. \end{aligned} \tag{8.23}$$

It is interesting to remark that the triple and the quartic H boson vertices vary as the square of the Higgs boson mass (for fixed W mass). As an indication of the strength of the Higgs boson couplings one finds 0.2 for the vertex $H^2W^+W^-$ and 0.12 for the quartic H^4 term.

To complete this section we recall the gauge boson self-couplings defined in $\mathcal{L}_{IG}$, eq. (5.57) : they are not affected by the spontaneous breaking of the symmetry, eventhough three gauge bosons have acquired a mass.

8.3 The Yukawa lagrangian $\mathcal{L}_Y$ and fermion masses and couplings

The scalar field Φ couples to fermions. The requirement for such couplings to exist is that the corresponding terms in the lagrangian density be Lorentz invariant as well as invariant under a $SU(2)_L \otimes U(1)_Y$ transformation, before the spontaneous breaking of this symmetry is implemented. Let us recall that Ψ_{e_L} and Ψ_{q_L} of eq. (4.14) and Φ of eq. (8.1) are **2** under $SU(2)$ *i.e.* they transform as

$$\delta\Phi = i\,\frac{\boldsymbol{\tau}}{2}\,\boldsymbol{\alpha}\,\Phi,\ \cdots,\quad \delta\Psi = i\,\frac{\boldsymbol{\tau}}{2}\,\boldsymbol{\alpha}\,\Psi,\ \cdots,\quad \delta\overline{\Psi} = -i\,\overline{\Psi}\,\frac{\boldsymbol{\tau}}{2}\,\boldsymbol{\alpha},\ \cdots \tag{8.24}$$

so that $\overline{\Psi}_{e_L}\Phi$, $\overline{\Psi}_{q_L}\Phi$ are invariant under a $SU(2)_L$ transformation and so are the hermitian conjugates $\gamma^0\Phi^\dagger\Psi_{e_L}$, $\gamma^0\Phi^\dagger\Psi_{q_L}$. Considering now the transformation properties under $U(1)_Y$: the combination $\overline{\Psi}_{e_L}\Phi$ has hypercharge 2 and $\overline{\Psi}_{q_L}\Phi$ hypercharge 2/3 so that $\overline{\Psi}_{e_L}\Phi\; e_R$ and $\overline{\Psi}_{q_L}\Phi\; d_R$ (see table eq. (4.27) for the hypercharge assignments) are invariant under a $SU(2)_L \otimes U(1)_Y$ gauge transformation. Their hermitian conjugates are : $\overline{e}_R\Phi^\dagger\;\Psi_{e_L}$ and $\overline{d}_R\Phi^\dagger\;\Psi_{q_L}$. Since these terms are also Lorentz invariants they satisfy all criteria to enter $\mathcal{L}_Y$.

One can construct another type of group invariant with the help of $\tilde{\Phi} = i\tau_2\Phi^*$ which is a $SU(2)$ doublet : indeed one can show

$$\delta\tilde{\Phi} \equiv \delta(i\tau_2\Phi^*) = i\tau_2\delta\Phi^* = i\tau_2(-i\frac{\boldsymbol{\tau}^*}{2}\,\boldsymbol{\alpha})\;\Phi^* = (i\frac{\boldsymbol{\tau}}{2}\,\boldsymbol{\alpha})\;\tilde{\Phi} \tag{8.25}$$

where one has used for the last equality the property $i\tau_2\boldsymbol{\tau}^* = -\boldsymbol{\tau}(i\tau_2)$. The combinations $\overline{\Psi}_{e_L}i\tau_2\Phi^*$and $\overline{\Psi}_{q_L}i\tau_2\Phi^*$ are invariant under a $SU(2)_L$ transformation and have hypercharge 0 and -4/3 ($y_{\Phi^*} = -y_\Phi = -1$), respectively. Thus $\overline{\Psi}_{q_L}i\tau_2\Phi^*\;u_R = \overline{\Psi}_{q_L}\tilde{\Phi}\;u_R$ is invariant under a group transformation. Had we included a right-handed neutrino the contribution $\overline{\Psi}_{e_L}i\tau_2\Phi^*\;\nu_R = \overline{\Psi}_{e_L}\tilde{\Phi}\;\nu_R$ would satisfy the conditions but we will ignore it here (see sec. 12). Thus the Yukawa lagrangian then takes the form

$$\boxed{\mathcal{L}_Y = -\;c_d\;\overline{\Psi}_{q_L}\Phi\;d_R - c_u\;\overline{\Psi}_{q_L}\tilde{\Phi}\;u_R - c_e\;\overline{\Psi}_{e_L}\Phi\;e_R + \text{h.c.} + \text{other families},} \tag{8.26}$$

where we have explicitely written out the terms involving the first family of fermions $(\nu, e; u, d)$. Six other parameters should be similarly introduced for the couplings of the second and third families so that nine new parameters appear in the model.

Implementing spontaneous symmetry breaking, in the unitary gauge, *i.e.* substituting in $\mathcal{L}_Y$ the expression of Φ as given in the right-hand side of eq. (8.14), we derive

$$\boxed{\mathcal{L}_Y = -\;c_d\;\frac{v+H}{\sqrt{2}}\overline{d}d\;-\;c_u\;\frac{v+H}{\sqrt{2}}\overline{u}u\;-\;c_e\;\frac{v+H}{\sqrt{2}}\overline{e}e\;+\;\text{other families.}} \tag{8.27}$$

From this expression we relate the mass of a fermion f to the vacuum expectation value v *via*

$$\boxed{m_f = c_f \frac{v}{\sqrt{2}}.} \tag{8.28}$$

This is not a prediction of the theory since the parameters c_f are unknown and will be adjusted so as to obtain the "physical" mass of the corresponding fermion. Furthermore, no relation is expected between the masses of partners of a given family since one parameter is introduced for each of the fermion type in a family. One may remark that the only "prediction" is that the neutrino remains massless as a consequence of the absence a right-handed neutrino. On the other hand, the Higgs couplings to the fermions are predicted, if the fermion masses are known,

$$\boxed{g_f = \frac{c_f}{\sqrt{2}} = \frac{m_f}{v} = \frac{e}{2\sin\theta_W} \frac{m_f}{M_W},} \tag{8.29}$$

where eqs. (8.28) and (8.19) have been used: the Higgs particle couples to a fermion flavour in proportion to the fermion mass, implying that the top quark could play a major role in the production and/or decay of the Higgs particle ($m_t \sim 175$ GeV) while the electron and light fermion contributions can be safely neglected. It is a puzzle why one observes such a large spectrum of masses from $m_e = .511\ 10^{-3}$ GeV to $m_t = 173.21$ GeV! No model naturally "explains" this fact.

• Remark

In sec. 2.3 we mentioned a problem related to massive gauge bosons namely the bad asymptotic behavior of the cross section of W pair production in e^-e^+ colliders. This was illustrated on the simpler case $\nu\ \bar{\nu} \to W^-\ W^+$ showing that the longitudinal polarisation states yield a cross section violating the Froissart bound if one keeps only the neutrino exchange diagram. Coming back to $e^-\ e^+ \to W^-\ W^+$ we leave it to the reader to check that, keeping fermion mass terms and including all diagrams in the unitary gauge, as shown in Fig. 2, the corresponding cross section is asymptotically finite. At higher orders, loop diagrams involve massive gauge boson propagators: in the unitary gauge

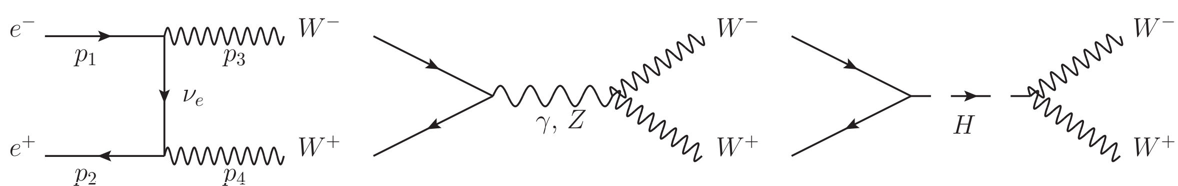

Figure 2: *The $e^-\ e^+ \to W^-\ W^+$ diagrams at lowest order in the unitary gauge.*

they do not converge to 0 when $k^2 \to \infty$ and this leads to an apparently non-renormalisable theory.

As explained in detail for the abelian case, the way out is to work in a "renormalisable" gauge (the 't Hooft gauges) where the gauge boson propagators have the form eq. (7.22) and the Golstone modes $\boldsymbol{\omega}$ are explicitely kept in the calculation.

8.4 The Higgs boson discovery

As an application we consider Higgs production in proton-proton colliders at the LHC at a center of mass energy $\sqrt{s} = 7$, 8 or 13 TeV. The Higgs boson could be produced in the annihilation of light

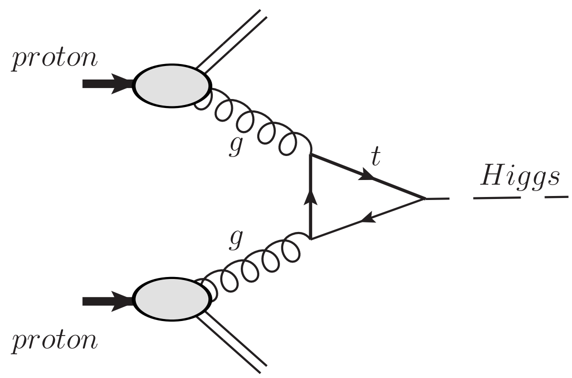

Figure 3: *Higgs production mechanism at hadron-hadron colliders : the dominant contribution arises from the subprocess where two gluons couple to the Higgs via a top quark loop. Another diagram with the fermion arrow reversed should be added.*

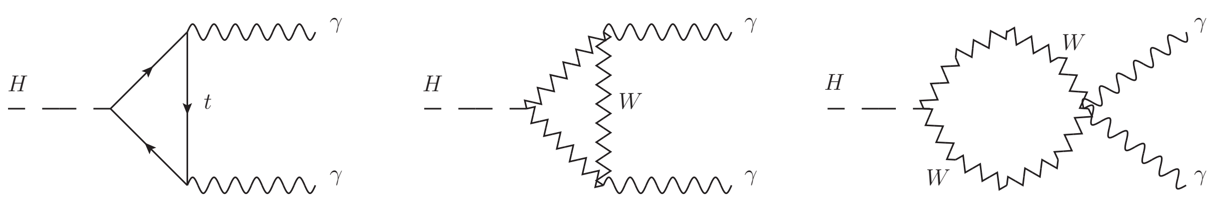

Figure 4: *Higgs decay mechanism into two photons: The dominant contributions arises from top quark loops and W boson loops. The final photons should be symmetrised in the first two diagrams.*

quarks and antiquarks of the initial hadrons, $q+\bar{q} \to H$, but such a coupling, eq. (8.29), is suppressed by a factor $m_f/v \simeq m_f/250$ with m_f, the mass of the quark, measured in GeV. The direct process $t\bar{t} \to H$ is possible but it is, of course, suppressed because of the negligibly small density of top quarks in the proton. For a Higgs mass below about 500 GeV it turns out that the dominant process is gluon-gluon fusion where the effective Higgs coupling to the gluon-gluon system is *via* a top quark loop as indicated in Fig. 3. The discovery channels of the Higgs boson have been $H \to Z\ Z^* \to 4$

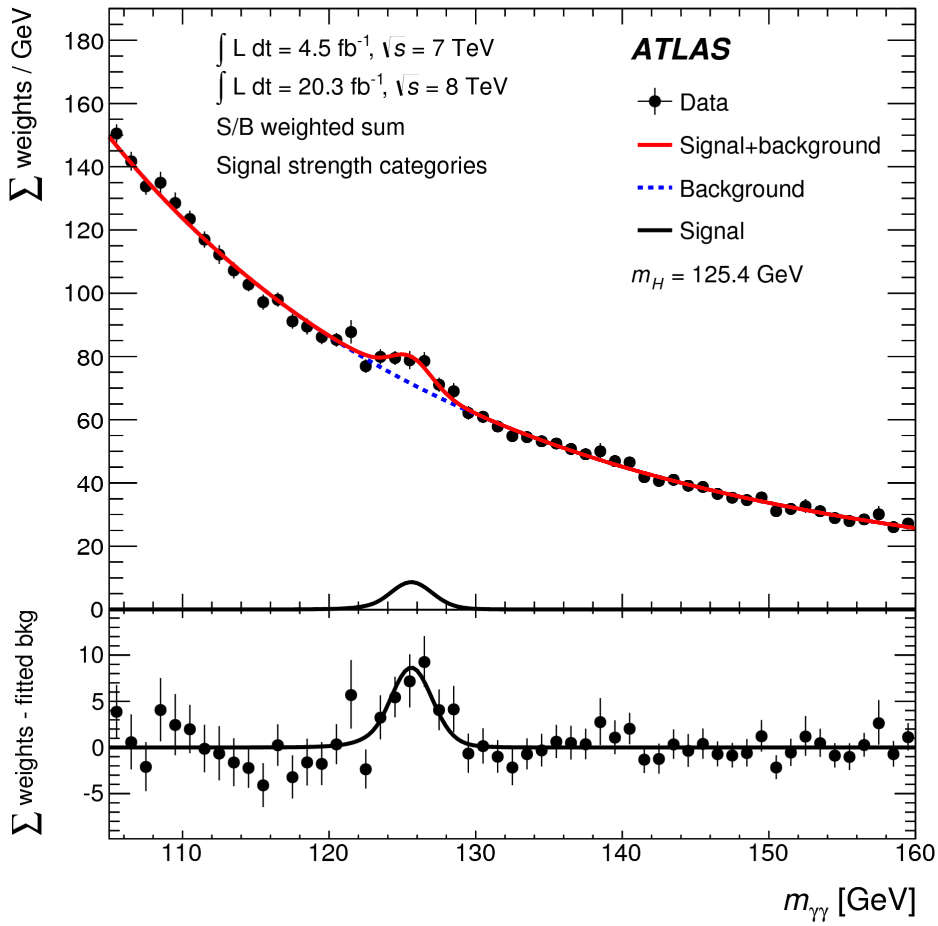

Figure 5: *ATLAS (Phys. Rev.* **D90** *(2014) 112015) results on Higgs observation through its decay into 2 photons.*

charged leptons (direct HZZ coupling $\propto e\,M_z$) and $H \to \gamma\,\gamma$. As shown in fig. 4 the two photon channel involves again a virtual top loop as well as $W^{\pm}$ loops. The results of the ATLAS and CMS collaborations are shown in Figs. 5 and 6: they illustrate the difficulty to extract the small $H \to \gamma\,\gamma$ signal from a huge background, essentially $q\,\bar{q} \to \gamma\,\gamma$ and its large associated QCD corrections. The H boson, of mass $M_H = 125.09$ GeV, cannot decay in a top pair of mass $2 * 173.21$ Gev but can decay into a bottom-antibottom pair. However the background in this channel is too large to be able to extract the Higgs signal, but the decay $H \to b + \bar{b}$, with H produced in association with a vector boson, has been studied by ATLAS[12] and CMS[13]. Other decay channels which have been considered and will studied at the High Luminosity LHC and the High Energy (27 GeV) LHC are $H \to WW^* \to l\nu l'\nu'$,

[12]ATLAS Collaboration, M. Aaboud et al., Phys. Lett. **B786** (2018) 59, arXiv:1808.08238 [hep-ex].
[13]CMS Collaboration,Phys. Rev. Lett. **121** (2018) 121801, arXiv:1808.08242 [hep-ex].

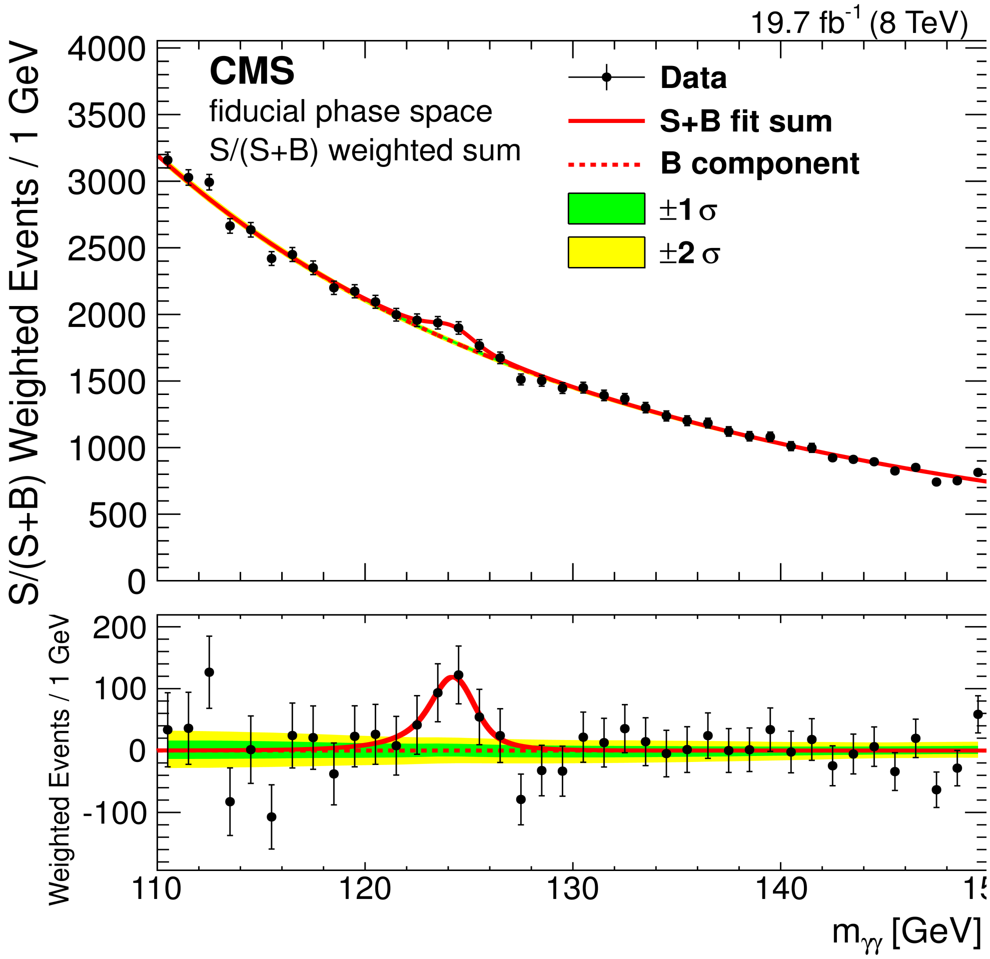

Figure 6: *CMS (Eur. Phys. J.* **C76** *(2016) 13) results on Higgs observation through its decay into 2 photons.*

$H \to \tau^- \tau^+$ with the τ's decaying leptonically or hadronically and $H \to \mu^- \mu^+$.[14]

The next two sections are devoted to a discussion of important production and decay mechanisms of the Higgs boson. Other exercises on the Standard Model can be found at `https://lectures.lapth.cnrs.fr/standard_model/cours/exo_en.pdf`.

8.5 Conclusions

At this point one has, in a first approximation, a complete model for the electroweak interactions. It contains a massive scalar particle, a massless and three massive gauge bosons, with propagators as in eqs. (8.18). All couplings between bosons and bosons to fermions are given assuming no mixing between the three generations of matter fields. The generation mixing is dealt with in sec. 11.

[14] Higgs Physics at the HL-LHC and HE-LHC, arXiv:1902.00134, [hep-ph].