Electroweak symmetry breaking, Higgs and Technicolor

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February 13, 2009

Abstract

This short review centers on electroweak symmetry and symmetry breaking in Standard Model of particle physics. The Higgs mechanism as well as Technicolor models of dynamical symmetry breaking are described and their advantages and shortcomings are discussed. Detection channels at the Large Hadron Collider are briefly described.

Contents

1 Intro: Electroweak theory

1.1 Motivation

The choice of topic deserves a few words of explanation, which is what I will attempt at in this (very) brief intoductory section. One could argue that electroweak theory (EW) is an 'old' part of the Standard Model (SM) of particle physics, verified to a great precision in many an experiment in the 1980s and 90s – in particular the massive vector bosons Z and W have been discovered as predicted. The electroweak symmetry is broken and while the features of this broken gauge symmetry have been indeed observed, the precise mechanism by which this occurs is still only conjectured. The best-known candidate theory for the (spontaneous) electroweak symmetry breaking is the famous Higgs mechanism, which is also responsible for fermion mass in the Standard Model. Thus we see that the EW theory is intimately linked to the problem of mass generation. In fact, the main scientific goals of the Large Hadron Collider (LHC) investigations, physics at $1\ TeV$ is directly related to EW symmetry breaking and so we come to realise that this subject is of prime importance.

We shall first revise the basics of electroweak theory and of the Standard Model Higgs mechanism following closely [27, 26], but we shall then proceed to describing various shortcomings of this proposition. This discussion will lead us to one of the alternative, beyond the Standard Model theories: the Technicolor. We shall briefly describe the main idea thereof and some important features, advertising the elegance of the dynamical solution, but also emphasizing the problem of flavour changing neutral currents in the theory. After a short review of various technicolor extensions we shall conclude with a section on experimental searches for Higgs boson at the LHC.

1.2 $SU(2)_L \otimes U(1)_Y$ theory

We shall not attempt describing the history of electroweak theory, the interested reader will find an account thereof in one of the references [27]. Let us just state, that experiments up to the 1950s prompted Fermi to write an effective Lagrangian in a *vector-axial current* form:

$$\frac{-G_F}{\sqrt{q}}\bar{\nu}\gamma_\mu(1+\gamma_5)e\bar{e}\gamma^\mu(1-\gamma_5)\nu \;+\; h.c. \tag{1}$$

Subsequent investigations revealed family structure of leptons, thus more terms of this type were included, but interestingly the strength of those *current-current* interaction proved to be the same, regardless of the family, hence the weak coupling constant was proven to be *universal*. Later on, quark doublets were added to the picture. We shall take those facts as starting point assumptions for constructing EW model:

1. experiment dictates the existence of quark and lepton left-handed weak isospin dublets: $\begin{pmatrix} \nu_e \\ e \end{pmatrix}_L$, $\begin{pmatrix} \nu_\mu \\ \mu \end{pmatrix}_L$, $\begin{pmatrix} \nu_\tau \\ \tau \end{pmatrix}_L$ and analogously for quarks: $\begin{pmatrix} u \\ d \end{pmatrix}_L$, $\begin{pmatrix} c \\ s \end{pmatrix}_L$, $\begin{pmatrix} t \\ b \end{pmatrix}_L$
2. experiment tells us also about the universality of the weak coupling constant
3. we take neutrinos to be massless

For notational convenience we shall only use one leptonic doublet in subsquent formluae, the generalisation of the formalism is easy. Hence we start with a theory of leptons only, transforming under the gauge group $SU(2)$. To incorporate electomagnetic interactions we need to add a $U(1)_Y$ weak hypercharge symmetry to the weak isospin $SU(2)$. Since no right-handed neutrinos were seen in an experiment, right-handed fermions must transform in singlet of $SU(2)$ - we now have a left-handed weak isospin doublet $\begin{pmatrix} \nu_e \\ e \end{pmatrix}_L$ and a right-handed singlet e_R with an assignment of hypercharges $Y_L = -1$ and $Y_R = -2$, respectively. The rule of thumb is hypercharge is twice the average electric charge in a doublet/singlet.

A local symmetry implies the existence of associated gauge fields, in this case isovector (b^i_μ) and isoscalar (A_μ) (this is not the A^μ of electromagnetism yet!) with couplings g, g'. The field-strength tensors read:

$$F^a_{\mu\nu} = \partial_\mu b^a_\nu - \partial_\nu b^a_\mu - g f^a_{bc} b^b_\mu b^c_\nu \quad \text{and} \quad f_{\mu\nu} = \partial_\mu A_\nu - \partial_\nu A_\mu \tag{2}$$

The lagrangian $\mathcal{L}$ contains as usual the kinetic terms for gauge fields and the terms for leptons with an appropriate covariant derivative and splits into two parts: $\mathcal{L} = \mathcal{L}_{gauge} + \mathcal{L}_{leptonic}$, where

$$\mathcal{L}_{gauge} = -\frac{1}{4}F^a_{\mu\nu}F^{a\mu\nu} - \frac{1}{4}f_{\mu\nu}f^{\mu\nu}, \tag{3}$$

$$\mathcal{L}_{leptonic} = \bar{R} i\gamma^\mu \left(\partial_\mu + i\frac{g'}{2} A_\mu Y_R \right) R + \bar{L} i\gamma^\mu \left(\partial_\mu + i\frac{g'}{2} A_\mu Y_L + i\frac{g}{2} b^i_\mu \tau^i \right) L, \tag{4}$$

where R is the righ-handed singlet (electron) and L is the left-handed doublet (electron and neutrino). We note that the covariant derivative for the right-handed fields does not contain a $SU(2)$ generator, as those trasform trivially under this symmetry.

There are two immediate observations about this lagrangian:

1. due to the gauge symmetry explicit mass terms are forbidden, as they would couple left and right-handed fields.

2. there are four massless gauge bosons in theory, but in reality we only have one - the photon.

The conclusion is that the gauge symmetry has to be broken, which is hardly surprising, since we know the observed symmetry of the theory is that of electromagnetism (i.e. it is the electric charge that is conserved). The necessary mechanism had already been known before, but in the context of condensed matter theory, rather than particle physics.

2 Breaking the symmetry: elementary Higgs

2.1 Ginzburg-Landau theory

There is a beautiful analogy in [27], which we shall present here. Ginzburg-Landau theory describes a superconducting phase transition. It assumes two types of carriers: resistive and the superconducting ones being responsible for charge transfer in the material. It is a phenomenological model, though it can be derived from miscorscopic BCS theory. The free energy of the superconducting phase in terms of the free energy of the normal phase is given by:

$$G_s(0) = G(0) + a|\psi|^2 + b|\psi|^4, \tag{5}$$

where $|\psi|^2$ is the density of superconductive charge carriers, or in other words ψ is the wavefunction of the superconducting state, a and b are parameters. The only obvious restriction is that $b \geq 0$ so that energy is bounded from below, a is not restricted. It is easy to see that if $a \geq 0$ then the absolute minimum of the energy corresponds to $|\psi|^2 = 0$, i.e. in the groundstate there are no superconducting carriers.

Change of sign of the a parameter below some critical temperature (*Curie* temperature T_c in fact) triggers the transition, when the the new minimum is at value of $\psi = \psi_0 \neq 0$ as shown in figure 1.

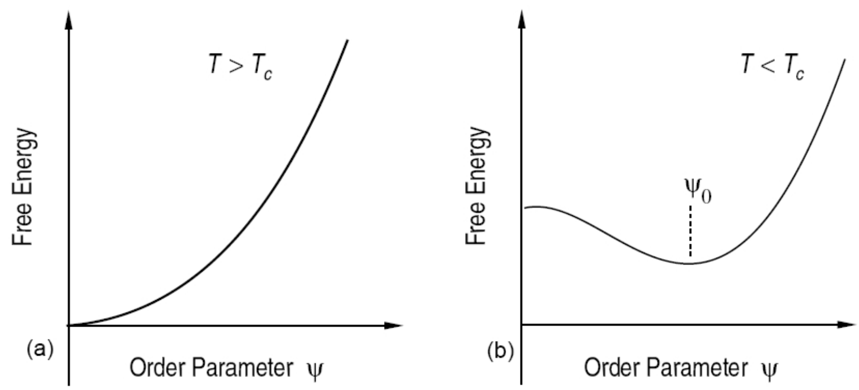

Figure 1: Ginzburg-Landau potential below and above the critical temperature. Taken from [26].

Now the minimum is at $|\psi|^2 = \psi_0^2 \neq 0$, hence the groundstate is superconducting.

We can extend this discussion to include external magnetic field $\vec{H}$. The expression for free energy in that case reads:

$$G_s(\vec{H}) = G_s(0) + \frac{\vec{H}^2}{8\pi} + \frac{1}{2m^*}| - i\hbar\nabla\psi - \frac{e^*}{c}\vec{A}\psi|^2 \tag{6}$$

From this we can obtain equations of motion for the fields. In slow-varying, weak field approximation we derive:

$$\nabla^2\vec{A} - \frac{4\pi e^*}{m^* c^2}|\psi_0|^2\vec{A} = 0, \tag{7}$$

where e^*,m^* are the effective charge and mass and c is the speed of light. Apparently the photon has become massive in the superconductor. The Higgs mechanism is a relativistic generalisation of this phenomenon.

2.2 Minimal Higgs model

This model was introduced in the early 1960s in numerous works:[1], [2], [25], [29],[16], [11], [15], but is now commonly refered to as the 'Higgs model'. We want to modify our theory in such way, that the lagrangian $\mathcal{L}$ is still invariant under the same $SU(2) \otimes U(1)$ gauge symmetry, but the groundstate of the theory is not. This is refered to as the *spontaneous symmetry breaking*, and was explored in the context of particles physics by – among others – Y. Nambu. To keep the lagrangian invariant we need correct field assignments:

- introduce a complex scalar doublet $\phi = \begin{pmatrix} \phi' \\ \phi^0 \end{pmatrix}$, with hypercharge $Y = 1$,
- add to the lagrangian a gauge invariant kinetic term and a potential: $(D^\mu\phi)^+(D_\mu\phi) - V(\phi^+\phi)$
- add possible coupling between new scalars and fermions of the theory in form of the Yukawa term

$$-\zeta_e \left(\bar{R}(\phi^+ L) + (\bar{L}\phi)R\right) \tag{8}$$

The potential has the form (analogy with Ginzburg-Landau is obvious):

$$V(\phi^+\phi) = \mu^2(\phi^+\phi) + \lambda(\phi^+\phi)^2. \tag{9}$$

When the parameter $\mu^2 \leq 0$ we develop a new minimum and the electroweak (EW for convenience) symmetry is spontaneously broken. The minimum of the energy may be chosen to correspond to $\langle\phi\rangle_0 = \begin{pmatrix} 0 \\ v/\sqrt{2} \end{pmatrix}$, where $v = \sqrt{-\mu^2/\lambda}$. This is referd to as the unitary gauge: in general the scalar doublet can be parametrised as $e^{iu^a(x)t_a} \begin{pmatrix} 0 \\ s(x)/\sqrt{2} \end{pmatrix}$, where t_a are the $SU(2)$ generators and s, u^a are four real fields - since the exponent is precisely of the form of a local $SU(2)$ gauge transformation, one can fix the gauge so that $u^a(x) = 0$. With this gauge choice we obtain minimum in the form mentioned before, i.e.

$$\phi = \begin{pmatrix} 0 \\ u(x)/\sqrt{2} \end{pmatrix} = \begin{pmatrix} 0 \\ (v + h(x))/\sqrt{2} \end{pmatrix}, \tag{10}$$

where $h(x)$ is the Higgs field.

Let us examine explicitly how symmetry is broken by the choice of the vacuum. If the vacuum were invariant the generators of the symmetry would annihilate it. What we obtain is however:

$$\begin{pmatrix} 0 & 1 \\ 1 & 0 \end{pmatrix} \begin{pmatrix} 0 \\ v/\sqrt{2} \end{pmatrix} = \begin{pmatrix} v/\sqrt{2} \\ 0 \end{pmatrix} \neq 0$$

$$\begin{pmatrix} 0 & -i \\ i & 0 \end{pmatrix} \begin{pmatrix} 0 \\ v/\sqrt{2} \end{pmatrix} = \begin{pmatrix} -iv/\sqrt{2} \\ 0 \end{pmatrix} \neq 0$$

$$\begin{pmatrix} 1 & 0 \\ 0 & -1 \end{pmatrix} \begin{pmatrix} 0 \\ v/\sqrt{2} \end{pmatrix} = \begin{pmatrix} 0 \\ -v/\sqrt{2} \end{pmatrix} \neq 0$$

$$Y\langle\phi\rangle_0 = \begin{pmatrix} 1 & 0 \\ 0 & 1 \end{pmatrix} \begin{pmatrix} 0 \\ v/\sqrt{2} \end{pmatrix} = \begin{pmatrix} 0 \\ v/\sqrt{2} \end{pmatrix} \neq 0 \tag{11}$$

So it looks like all 4 symmetries are broken, but in fact there is one vanishing combination:

$$Q\langle\phi\rangle_0 = \frac{1}{2}(\tau_3 + Y)\langle\phi\rangle_0 = \frac{1}{2} \begin{pmatrix} 1+1 & 0 \\ 0 & -1+1 \end{pmatrix} \begin{pmatrix} 0 \\ v/\sqrt{2} \end{pmatrix} = \begin{pmatrix} 0 \\ 0 \end{pmatrix} \tag{12}$$

This particular linear combination of generators correspods to charge generator Q of the electromagnetism. Hence we have achieved $SU(2)_L \otimes U(1)_Y \to U(1)_{EM}$.

The Goldstone theorem tells us that for every spontaneously broken continuous global symmetry we should have a massless boson, but since we have broken a local symmetry the bosons get 'eaten' – they become the longitudinal components of the vector gauge bosons, which therefore acquire mass (massive vector fields have 3 degrees of freedom vs. 2 for the massless case). We shall have three massive gauge bosons and one masless – the photon. After a suitable field redefinition we recover our usual Z and W bosons:

$$W^{\pm} = \frac{b_1 \mp i b_2}{\sqrt{2}} \quad \text{with mass} \quad M_W = \frac{gv}{2}$$

$$Z = b_3 \cos\theta_W - A \sin\theta_W \quad \text{with mass} \quad M_Z^2 = \frac{M_W^2}{\cos^2\theta_W} \tag{13}$$

where $g' = g \tan\theta_W$

The construction described above is known as the *Higgs mechanism*. We note that that the value of the gauge bosons' mass depends on the vacuum expectation value of the Higgs field, and that there is a particular relation (exact at tree-level only) between the masses of W and Z. This mass relation has been verified experimentally, hence any other theory will have to reproduce it. Higgs mechanism also solves the problem of fermion masses: the standard mass term for fermions would couple left-handed fermions transforming as singlets of $SU(2)$ and right-handed fermions transforming as doublets, breaking the invariance of the lagrangian. In Higgs model the masses are generated in the Yukawa terms – which first couple the scalar doublet with fermion doublet and only then take product with fermion singlet, hence the mass term is $SU(2)$ singlet as it should be – and they depend not only on the v, but also on the coupling constant ζ_e as can be seen from equation 8. Most of the free parameters of the Standard Model are actually Yukawa couplings. It is important to remark, however, that it is not required that the same mechanism generates masses for both the gauge bosons and the fermions – it is a particular feature of this model. The Higgs mass is $M_H^2 = -2\mu^2$, but this is an *a priori* arbitrary parameter of the theory, hence there is no prediction for this value! We shall see in a moment that requirements of consistency of the theory place some bounds on values of M_H.

2.3 FCNC problem

This is a somewhat historical digression, but one that will prove to be important, since we shall encounter similiar problems in our discussion of Technicolor. Originally one quark doublet $\begin{pmatrix} u \\ d \end{pmatrix}_L$ was known, but that proved to be at odds with the experimental results involving charged currents.

Cabibbo postulated the following solution: $\begin{pmatrix} u \\ d_\theta \end{pmatrix}_L$, where $d_\theta = d\sin\theta_C + s\cos\theta_C$., i.e. a doublet with strange and down quark mixing, where the angle θ_C is known as the Cabibbo angle. This allowed to handle the charged currents discrepancy, but at a cost of introducing Flavor Changing Neutral Currents (FCNC) i.e. interactions of neutral vector boson with quarks in which the flavor of the quark changes. The Z-quark-quark term in the Lagrangian of Cabbibo contained for example the following:

$$Z_\mu \left[C_1 \cdot \bar{d}\Gamma_1^\mu s + C_2 \cdot \bar{s}\Gamma_2^\mu d \right], \tag{14}$$

where d,s denote the down and strange quarks, Γ_i^μ stand for generic gamma-matrix structures and C_i are constants. It is now clearly visible that a strange quark may turn into a down quark and vice versa by emission of a neutral Z boson. Precision electroweak measurements place a very tight bound on such currents and virtualy exclude this possibility (see for example discussion of Kaon system in and oblique corrections in [27]).

A further development was needed, known as the GIM soltion: introduce a new (at that time unobserved!) c quark and make two doublets: $\begin{pmatrix} u \\ d_\theta \end{pmatrix}_L, \begin{pmatrix} c \\ s_\theta \end{pmatrix}_L$, where $s_\theta = s\cos\theta_C - d\sin\theta_C$, so that the cross-terms cancel out and there are no more FCNCs in theory.This proved to be a correct approach - charm quark was discovered and theory reconciled with experiment. We shall later discover that

FCNCs are going to be the main trouble of Technicolor models as well.

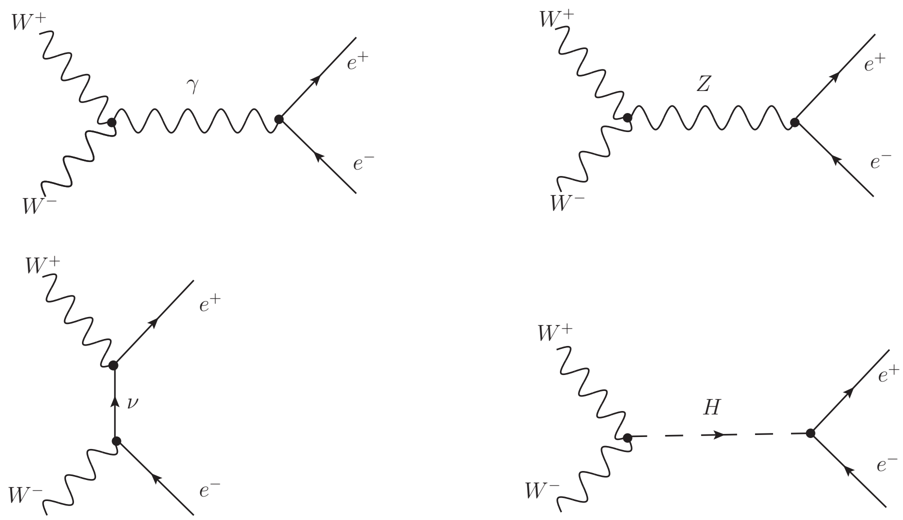

Figure 2: Tree-level diagrams contributing to $W^+W^- \to e^+e^-$ scattering

2.4 Additional motivation for a 'Higgs'

We have introduced Higgs mechanism as a means of breaking the electroweak symmetry (and consequently generating mass), but there are in fact other reasons to believe that an additional field should be there in the theory. Consider an example of W bosons scattering to a pair e^+e^-, the relevant tree-level diagrams are depicted in the figure below:

We can in perturbation theory calculate the cross-sections for such a process, including all or only some of the channels as shown in figure 2. Experiments have measured the total cross-section for such scattering, hence we can compare our predictions. The results of this comparison are in the next figure. It turns out that if we do not include the gauge boson exchange channel the theoretical prediction diverges with centre of mass (CM) energy of the incoming particles s. Upon including this contribution the agreement is much better as shown in figure 3, but still the discrepancy grows as $\sqrt{s}$. It is therefore necessary to have additional channel (Higgs exchange) to match the experimental results.

3 Higgs problems

The discussion we had so far could have conveyed the impression, that elementary Higgs model is the only possibility and answers all unresolved questions about

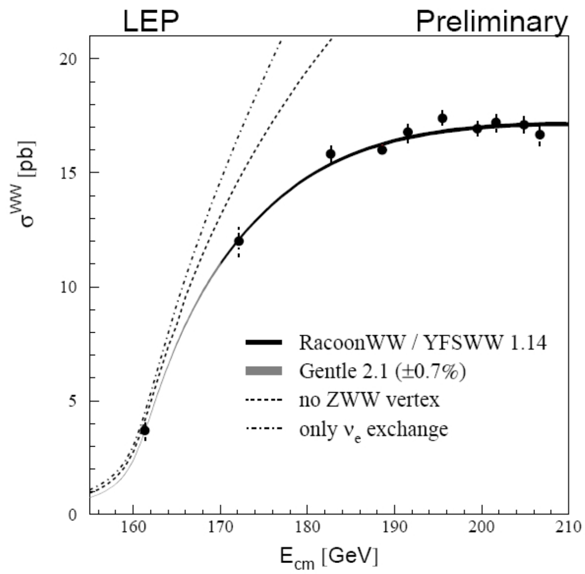

Figure 3: $W^+W^- \rightarrow e^+e^-$ scattering: theoretical predictions vs. experimental results. Taken from [26].

Standard Model. This is not the case: the minimal Higgs model, which introduces only one doublet of complex scalar fields suffers from some important problems, which limit its validity. We shall describe some of the issues below. A thorough discussion can be found in reference [14].

3.1 Higgs bounds and the TeV scale

We can consider the process of gauge boson scattering at tree-level again. If we apply the partial-wave expansion known from classical scattering theory to the amplitude M of this process we obtain the following expression:

$$M \equiv \sum_J (2J+1) a_J(s) P(\cos\theta), \tag{15}$$

where P are the Legendre polynomials, $a_J(s)$ are functions depending on the centre of mass energy of the scattering particles and the sum runs over angular momentum. We can then calculate the *total* cross-section noting that due to the orthogonality of the polynomials P it also can be written as $\sum_J \sigma_J$, which is a statement of separation of the process into channels corresponding to exchange of particles of different angular momentum J.

Partial wave-unitarity is a requirement that the total probability for scattering not exceed unity, which applied to the $J = 0$ channel (i.e. Higgs exchange) brings about the condition for Higgs mass:

$$M_H \leq \left(\frac{8\pi\sqrt{2}}{3G_F}\right)^{1/2} = 1\,TeV/c^2 \tag{16}$$

What we learn from this *perturbative* analysis is that:

- if this bound is violated the perturbation theory breaks down and weak interactions become 'strong'
- this is a signal that new phenomena assosciated with electroweak interactions are to be expected at 1 TeV scale.

3.2 Triviality

There is yet another consistency problem of the minimal Higgs model, associated with renormalisation properties of the theory. It turns out that we can also extract an upper bound on Higgs mass M_H, or, equivalently an upper scale Λ to which the theory is perturbatively valid. Let us make a few remarks on that:

Higgs is a scalar particle, but it is well-known that only non-interacting scalar field theories are valid on an arbitrarily high energy scales. To illustrate this fact let us consider a simplified result of $\lambda\phi^4$ theory. We can write the equation for the running of the coupling constant (i.e. relating the strength of the coupling at different energy scales). We have:

$$\frac{1}{\lambda(\mu)} = \frac{1}{\lambda(\Lambda)} + \frac{3}{2\pi^2}\log\frac{\Lambda}{\mu}, \tag{17}$$

Where Λ is some high reference scale at which we fix the value of the coupling constant and μ is the low energy physical scale

- if we want our theory to be valid at all energies then we take the limit of the reference scale $\Lambda \to \infty$ while keeping the physical scale μ fixed. Since $\lambda(\Lambda)$ is finite it follows that $\lambda(\mu) = 0$ i.e. the theory at low energy is non-interacting or *trivial*.
- we can also rewrite equation 17:

  $$\lambda(\Lambda) = \frac{\lambda(\mu)}{1 - (3\lambda(\mu)/(4\pi^2))\log(\Lambda^2/\mu^2)}, \tag{18}$$

  this provides us with an alternative interpretation: no matter how small the coupling at physical scale $\lambda(\mu)$ is, the coupling at the reference point in infinity explodes i.e. we run into Landau pole.

Using the same equation for running of the coupling constant we can demand that $\lambda \geq 0$ at all scales, so that Higgs potential is bounded from below. We can therfore neglect the $\frac{1}{\lambda(\Lambda)}$ term and obtain the inequality:

$$\lambda(\mu) \leq \frac{2\pi^2}{3\log\frac{\Lambda}{\mu}}, \tag{19}$$

which we can rewrite as:

$$\Lambda \leq \mu \exp\left(\frac{2\pi^2}{3\lambda(\mu)}\right) \tag{20}$$

if we set our physical scale to $\mu = M_H$ and use the tree-level relation $M_H^2 = 2\lambda(M_H)v^2$ (i.e. no quantum corrections) we get:

$$\Lambda \leq M_H \exp\left(\frac{4\pi^2 v^2}{3M_H^2}\right) \tag{21}$$

- this tells us that for a given M_H there is a scale at which the theory ceases to make sense.
- lattice calculations show that in order to have some predicitve power at few TeV we need $M_h \leq 710 \pm 60\, TeV/c^2$

There is also a lower bound on the value of Higgs mass - it is provided by requiring vacuum stability, i.e. that upon including 1-loop corrections the minimum of the potential still satisfies $\langle\phi\rangle_0 \neq 0$ up to a scale Λ (it is intuitively clear that the value of the *vacuum expectation value* receives quantum corrections as it is related to the Higgs mass). This is crucial for our theory, since we necessarily need a non-zero vacuum expectation value of the Higgs potential to achieve the electroweak symmetry breaking. It is interesting to note, that this inequality (can be found in reference [14])

$$M_H^2 \geq \frac{3G_F\sqrt{2}}{8\pi^2}\left(2M_W^4 + M_Z^2 - 4m_t^4\right)\log\frac{\Lambda^2}{v^2}, \tag{22}$$

was derived when the mass of the top quark was not yet measured, but was thought not to exceed $\approx 80\, GeV/c^2$. It is in fact twice as big, hence this particular estimate gives a trivial bound.

A complete 2-loop calculation results give bounds as in the figure 4. If the theory is to make sense up to $\Lambda = 10^{16}\, GeV/c^2$ then $134 \leq M_H \leq 177\, GeV/c^2$.

We would like to emphasize that all the above results were derived in the framework of standard perturbation theory. One should not take them therefore at face

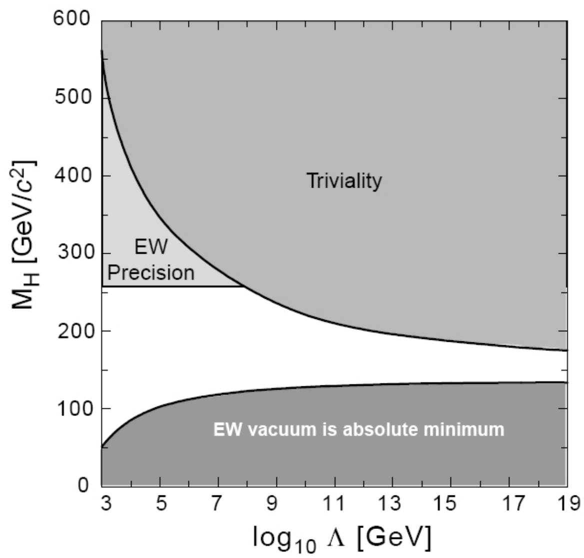

Figure 4: Higgs mass constrained by triviality and vacuum stability. Taken from [26].

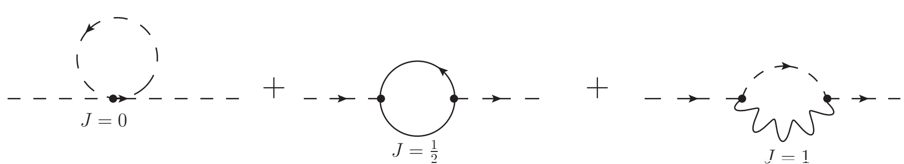

Figure 5: 1-loop corrections to Higgs mass

value, since some of them may just signal the breakdown of perturbative approach, i.e. that the features of electroweak theory at $1\ TeV$ scale are described by non-perturbative effects or in yet another words, that the electroweak theory becomes strongly coupled, much as QCD is in certain regimes. It is then a priori possible that the theory is valid up to a scale Λ, but its perturbative approximation breaks down at $\Lambda_p < \Lambda$. That situation still allows us to probe the region above Λ_p, using, for example, lattice calculations. A more thorough discussion can be found in [14].

Let us now mention the problem of Higgs mass divergencies in theory. The relation $M_H^2 = \lambda v^2$ is valid only at tree-level i.e. it is the *bare mass*. As we well know the mass receives quantum corrections. The figure 5 shows the 1-loop corrections due to virtual particles of spin $J = 0, \frac{1}{2}, 1$ - those corrections for elementary SM Higgs model are quadratically divergent! This is related to the fact that the loop momenta are a priori unrestricted i.e. they can take arbitrarily high values and the radiative corrections blow the Higgs mass up. A typical treatment of this conceptual problem is to assume a cut-off scale Λ, which restricts the momenta running in the loops. This is a statement of the fact that our theory is only valid up to a certain energy scale, hence calculations with unconstrained loop integrals invariably bring about trouble. The question therefore is what the value of Λ should be. If we demand that the radiative corrections be small then we have only two ways out of the problem:

- either Λ is to be kept small,
- or 'new physics' cuts off the integrals

Since the scale of EW symmetry breaking given by the value of the vacuum expectation value of the Higgs potential is $246\ GeV$, the relevant scale is $\Lambda = 1\ TeV$. In Technicolor models, Higgs is not an elementary particle, it is composite and at its binding energy scale new physics intervenes.

We should note here that in supersymmetric theories this problem is solved in a very elegant way: the extended matter content of SUSY gives cancellations between the fermionic and bosonic loops, hence mass of the Higgs is protected from radiative corrections.

All of this yet again emphasizes the importance of the TeV energy scale and justifies expectations of new experimental discoveries at LHC.

4 Higgs @LHC

The subject of experimental Higgs searches is a vast one, and mostly beyond the scope of this report. There is a wealth of literature on Higgs searches, we have mainly benefitted from [14] and [8].

With the LHC beginning operation in June 2009, the expectation is to 'find' Higgs within the mas range $115 - 200\ GeV/c^2$ in the next few years. Earlier searches, culminating with the LEP2 experiment have already exlcuded lower values. Below we will briefly describe some of the Higgs production and decay processes relevant for this energy range and comment on their respective properties. Some of them are potential *discovery channels*. It is however very important to keep in mind that even successful discovery of 'Higgs' in one of them will not be sufficient to draw conclusions as to the validity of the Higgs mechanism of symmetry breaking we described in previous sections. To achieve that a whole range of measurements will have to be performed to measure the various Yukawa couplings (ratios, actually). Therefore the experiments will study numerous processes

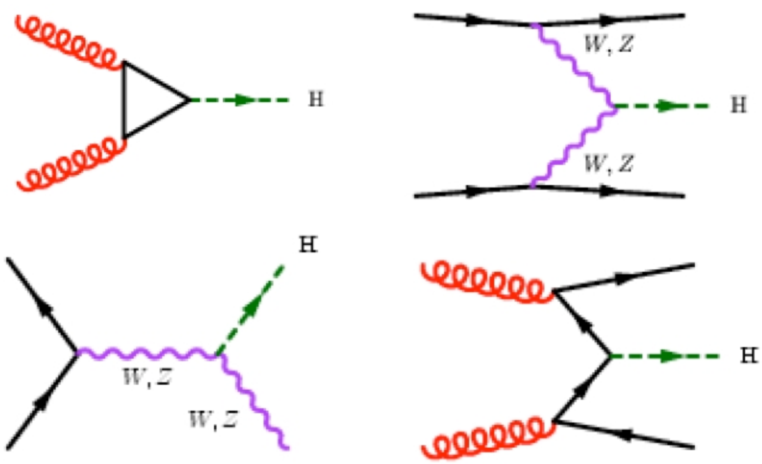

Figure 6: Higgs production channels. Taken from [8].

involving Higgs, exploiting the fact that their respective cross-sections are proportional to different Yukawa paramers. This is still more important in the context of beyond SM investigations. To be able to distinguish between their predictions (i.e. 'new physics') and the signatures of elementary Higgs one needs a complete experimental programme - this is the real task of the LHC. Some of the most important modes of Higgs production are depicted in figure 6 and descibed below.

The main production channels are as follows:

- gluon fusion $gg \to H$
  - largest rate for all Higgs masses M_H
  - proportional to the Yukawa-top quark coupling y_t
- weak boson fusion (WBF) $qq \to qqH$
  - second larest production rate
  - proportional to the WWH coupling
- Higgs-strahlung $q\bar{q} \to W(Z)H$
  - third largest production rate
  - the same coupling as in weak boson fusion
- $t\bar{t}(b\bar{b})H$ associated production

– proportional to Yukawa-quark couplings

A typical calculation of a process of the type $q\bar{q} \to ZH$, for which the vertices are shown in figure 7, would go as follows. The cross section is given as:

$$\sigma = \int dPS(2) \frac{1}{\mathcal{F}} \sum_{r,t,j,a} |\mathcal{M}|^2, \tag{23}$$

where $\mathcal{F}$ is the flux – in the massless quarks limit we have $\mathcal{F} = 2s = 2(p_1 + p_2)^2$ –, the sum is over quark spins r,t and colors a and vector boson polarisations j and contains averaging factor $1/36$ for the incoming quarks. We also integrate over 2-particle phase space of Z and Higgs boson.

The expression for amplitude $\mathcal{M}$ (in the conventions of [31]) reads:

$$\mathcal{M} = \frac{(2\pi)^4}{(p_1 + p_2)^2 + M_0^2 + i\varepsilon} \frac{ig^2}{\sqrt{2k_1^0}\sqrt{2k_2^0}} \frac{M_0}{c_w^2} \bar{e}^j_\mu(k_2) \bar{u}_r(p_2) \gamma^\mu \left[1 - \frac{8}{3} s_w^2 + \gamma^5\right] u_t(p_1), \tag{24}$$

where M_0 is the Z boson mass, s_w, c_w are the sine and cosine of the weak mixing angle θ_W; u, $\bar{u}$ are the spinors for external fermions and antifermions and $\bar{e}_\mu$ is the polarisation vector for external massive gauge boson. The amplitude needs to be squared and spimplified with help of the following identities:

$$\sum_{t=1}^{2} u_t(p) \bar{u}_t(p) = \frac{1}{2p^0}(-i\not{p}) \quad \text{for massless fermions} \tag{25}$$

$$\sum_{r=1}^{2} u_r(p) \bar{u}_r(p) = \frac{1}{2p^0}(-i\not{p}) \quad \text{for massless antifermions} \tag{26}$$

$$\sum_{j=1}^{3} e^j_\mu(k) \bar{e}^j_\nu(k) = \delta_{\mu\nu} + \frac{k_\mu k_\nu}{M_0^2} \quad \text{for massive vector bosons} \tag{27}$$

Note that there is no difference in fermion/antifermion spin-sums in the massless limit. We would then obtain:

$$|\mathcal{M}|^2 \sim A \left(C_1 \cdot Tr[\not{p}_2 \not{k}_2 [...] \not{p}_1 \not{k}_2] + C_2 \cdot Tr[\not{p}_2 \gamma^\mu [...] \not{p}_1 \gamma_\mu]\right), \tag{28}$$

where $A = \frac{g^4 M_0^2}{4k_1^0 k_2^0 c_w^4} \frac{1}{\left[(p_1+p_2)^2 + M_0^2 + i\varepsilon\right]^2}$, $C_1 = \frac{-1}{4p_1^0 p_2^0}$, $C_2 = \frac{-1}{4p_1^- p_2^0 M_0^2}$ and [...] stands for $[1 - \frac{8}{3}s_w^2 + \gamma^5]$.

Traces may be evaluated in the usual way:

$$Tr[\not{p}_2 \not{k}_2 [...] \not{p}_1 \not{k}_2] = (1 - \frac{8}{3} s_w^2) Tr[\not{p}_2 \not{k}_2 \not{p}_1 \not{k}_2] + Tr[\not{p}_2 \not{k}_2 \gamma^5 \not{p}_1 \not{k}_2] =$$

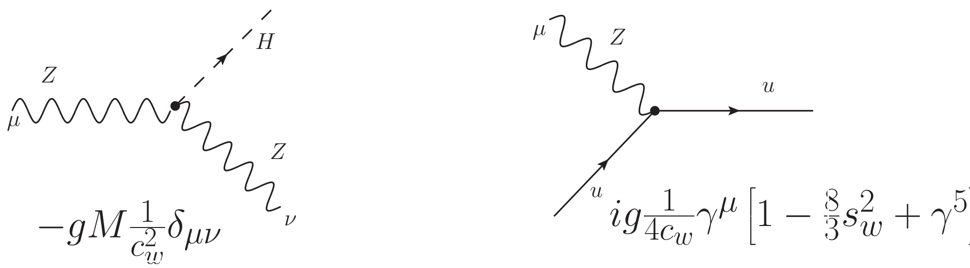

Figure 7: SM vertices relevant for the process under consideration

$$= 4(1-\frac{8}{3}s_w^2)[2(p_2\cdot k_2)(p_1\cdot k_2)-(p_2\cdot p_1)(k_2\cdot k_2)], \tag{29}$$

where we used that fact that $Tr[\gamma^\mu\gamma^\nu\gamma^\rho\gamma^\sigma\gamma^5]\sim\varepsilon^{\mu\nu\rho\sigma}$ hence this term will vanish, as it is mulitplied by an expression $k_{2\mu}k_{2\nu}p_{1\rho}p_{2\sigma}$ symmetric in μ, ν.

$$Tr[\not{p}_2\gamma^\mu[...]\not{p}_1\gamma_\mu]=(1-\frac{8}{3}s_w^2)Tr[\not{p}_2\gamma^\mu\not{p}_1\gamma_\mu]+Tr[\not{p}_2\gamma^\mu\gamma^5\not{p}_1\gamma_\mu]=$$

$$-2(1-\frac{8}{3}s_w^2)Tr[\not{p}_2\not{p}_1]-2Tr[\not{p}_2\not{p}_1\gamma^5]=-8(1-\frac{8}{3}s_w^2)p_1\cdot p_2, \tag{30}$$

where we used the reduction formula in $n=4$ dimensions: $\gamma^\mu\gamma^\alpha\gamma_\mu=-2\gamma^\alpha$. Plugging this in we obtain an expression purely in terms of kinematical variables p_1, p_2, k_1, k_2 and constants. To obtain a numerical value (albeit of little use, since the Higgs will decay immediately anyways), we would still need to parametrise the momenta (usually in the incoming particles' centre of mass frame) and perform the phase space integral, which in this case would involve only integration over the azimuthal angle.

The Standard Model vertices in the Feynman gauge have been taken from [31].

Higgs, being a short-lived, neutral scalar particle will not, of course, be directly visible to the experiment. The only accessible information comes through the various decay products. The same remark about the importance of studying different channels applies in this case – some of them are proportional to Higgs-quark couplings (mostly top), while others are dominated by the EW coupling. Some of the principal modes are shown in the figures 8 and 9.

The branching ratios of Higgs decays are M_H-dependent, hence channels will have varying significance in different energy ranges. The branching ratios relevant for Higgs search at the LHC are shown in the figure 10.

Higgs, as we remarked before, is only accesibile through the decay products. However, similiar products appear of course in various other processes which do

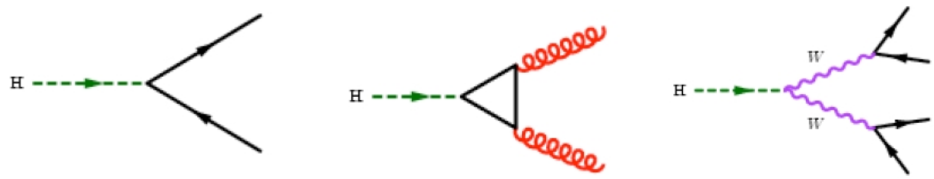

Figure 8: Higgs decay channels. Taken from [8].

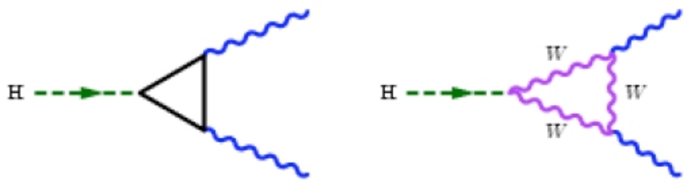

Figure 9: Higgs decay channels. Taken from [8].

not involve Higgs i.e. we have a strong QCD background. Experimental investigation it therefore conducted by means of *inclusive searches*, wherein an interesting process is singled out, all other (QCD) processes contributing to the same final decay products are identified. Monte Carlo simulations are then performed for the background and the invariant mass spectrum is contructed. In a real experiment the invariant mass spectrum for the interesting final state is measured - the Higgs should appear as additional contribution on top of the background. Inclusive Higgs searches at the LHC involve the processes depicted in figure 11 and described below:

- $H \to \gamma\gamma$
  - has a very small branching ratio, i.e. Higgs boson decays via this channel with small probability compared to other possible channels,
  - and suffers from a strong $pp \to \gamma\gamma$ background, i.e. there are many other QCD processes which result in the same final state particles ($\gamma\gamma$).
  - but ATLAS and CMS detectors have a very good photon energy resolution and should be able to see a Higgs peak.
- $H \to ZZ \to l^+l^-l^+l^-$
  - small branching ratio
  - but least background

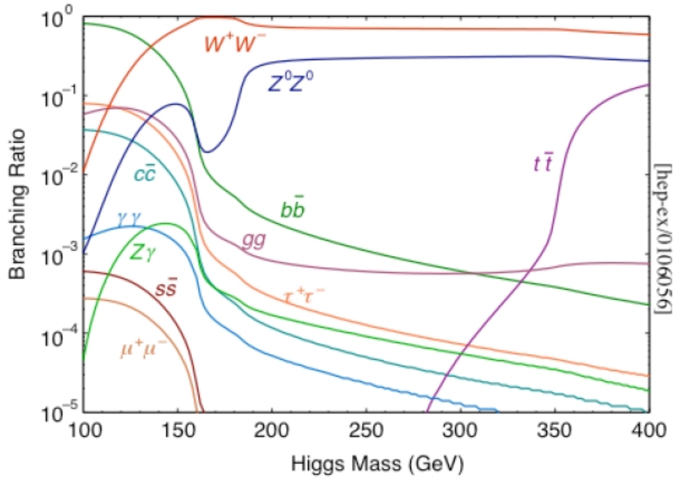

Figure 10: Higgs decay branching ratios as a function of Higgs mass. Taken from [8].

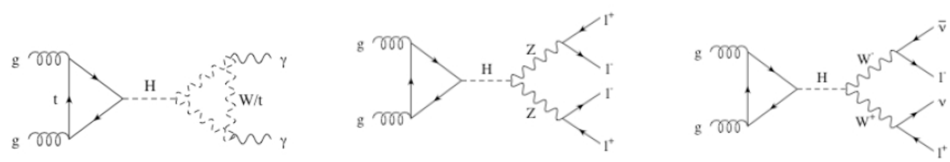

Figure 11: Processes used in inclusive Higgs searches. Taken from [8].

  - discovery mode for heavy Higgs (0.8 -1.0 TeV)

- $H \to WW \to l^+\nu l^-\bar{\nu}$

Also a simulation of Higgs $H \to \gamma\gamma$ peak on top of the background in the invariant mass spectrum is shown in figue 12.

5 Technicolor and extensions

5.1 Motivation

Having described the standard Higgs mechanism of the electroweak symmetry breaking, we can move on to one of the possible alternatives, falling into the category of the 'beyond the standard model' theories. Before we start describing the

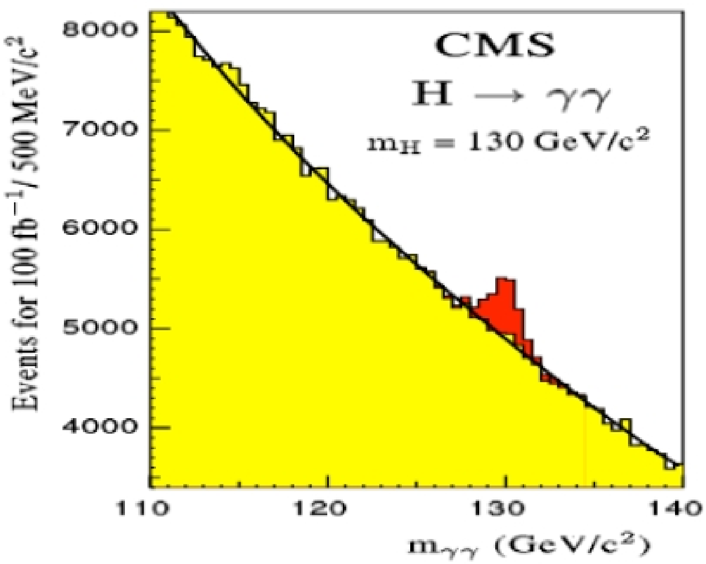

Figure 12: A simulated invariant mass plot with a peak corresponding to a Higgs boson. Taken from [8].

features of this model let us provide some motivation (or in other words reasons for dissatisfaction with the elementary Higgs model) - some of which we have already become acquainted with in section 3. We shall follow the discussion in reviews [20, 22, 19], and also [5, 6]

- elementary Higgs gives no *dynamical* explanation for EW symmetry breaking - the Higgs potential we've used to spontaneously break the electroweak symmetry is put there by hand. We have an experimental constraint on the vacuum expectation value $v = 246\ GeV/c^2$ (from the vector gauge boson masses), but we do not have any dynamical reason for the shape of the potential - it is therefore a phenomenological explanation, much like Ginzburg-Landau theory.

- elementary Higgs models are *unnatural* - require fine tuning to great precision. This is related to our discussions of radiative corrections of Higgs mass. Since the corrections may be very high (depending on the cut-off scale) and we want a rather low physical higgs mass, then to balance it out we must have a very high bare mass parameter, but more then that, it has to be set precisely to such a value that the difference between the bare mass and corrections (both, say, of order $10^{16}\ GeV$ if the relevant scale is the unification scale) gives the physical mass of order $10^2 - 10^3\ GeV$. This is known as *fine tuning*.

- elementary Higgs models are *trivial*. We have discussed the questions of validity of the model and the perturbative approach in section 3.

- elementary Higgs models do not tell us much about flavor physics. The number of generations is completely independent of the Higgs mechanism and the Yukawa couplings are free parameters of the theory. It is a valid approach to search for a theory that has more to say about the origin of flavour at $1\ TeV$ scale.

5.2 The power of analogies

With this motivation we can set off to describe one of the possible solutions. The Technicolor model was first introduced in 1979 [32], [30] and its extensions soon after [10], [9]. It is modelled on a mechanism of dynamical mass generation and symmetry breaking already present in the pure QCD, but with extended gauge group (the reasons for that we shall describe in some detail). The main difference to the Higgs mechanism in particular can be nicely described:

- if elementary Higgs mechanism of spontaneous symmetry breaking is modelled on Ginzburg-Landau phase transition, then Technicolor can be thought of as being based on the BCS theory of superconductivity. The phenomenological potential in Ginzburg-Landau theory provided a correct description of the phase transition in terms of the wave-function of the superconducting state, but microscopic reasons for this were not know. In the framework of BCS theory this is explained by Cooper pairs of electrons forming due to attractive interaction via exchange of phonons. Below the critical temperature T_c this attractive interaction becomes strong enough for the electronic pairs to Bose-Einstein condense.

- Analogously, dynamics of technifermion gauge interactions generates scalar bound states as in BCS model of superconductivity. The extended gauge group has a bigger matter content, in particular new technifermions interacting through technigluons exchange. At certain energy scale those interactions become strong enough for the condensates to appear and break the symmetry. This is modelled on the creation of hadronic bound states in QCD.

- Higgs is therefore a (techni)fermion-antifermion bound state.

We have invoked an analogy with QCD - it is an interesting and very useful excercise to pursue this direction and examine the mass generation in $SU(3)_c \otimes SU(2)_L \otimes U(1)_Y$ theory.

Let us for simplicity consider theory of massless u, d quarks with $SU(3)_c \otimes SU(2)_L \otimes U(1)_Y$ gauge symmetry. There is the standard QCD lagrangian obtained by replacing the usual derivative in the kinetic terms by a covariant one. The Higgs is of course absent in this discussion. This lagrangian is massless (explicit mass terms forbidden by symmetry), hence is posseses an exact $SU(2)_L \otimes SU(2)_R$ (global) chiral symmetry. We know that this non-abelian gauge theory is asymptotically free, i.e. the interaction (which is attractive) is strong at low energy, hence, at a sufficiently low scale Λ_{QCD} fermion condensates appear $\langle \bar{q}q \rangle \neq 0$. But since these couple left and right-handed quarks the chiral symmetry is broken. Thus we have achieved spontaneous chiral symmetry breaking:

$$SU(2)_L \otimes SU(2)_R \to SU(2)_V.$$

By the Goldstone theorem there will be massless bosons corresponding to the broken symmetries - $SU(2)_L \otimes SU(2)_R$ has 6 generators, while $SU(2)_V$ only 3, hence three Goldstone boson have to appear (Nambu identified them with massless pions). Since the symmetries were local, massless pions become longitudinal components of EW gauge bosons, which therefore acquire mass.

So we have discovered that the electroweak symmetry is spontaneously broken in a natural, *dynamical* way in QCD, and mass for vector gauge bosons is generated. Can QCD alone provide solution? The answer is negative:

- while QCD condensates generate mass the value would be $M_W \approx 30\ MeV$ - which is by far too small.
- but the tree-level relation between M_W and M_Z is correct.
- the vacuum expectation value of the condensates $\frac{1}{2}\langle\Omega|\bar{q}q|\Omega\rangle = 4\pi f_\pi^3$ is related to vector gauge boson mass by: $M_W = \frac{g f_\pi}{2}$

The lesson therefore is, that QCD spontaneously breaks the EW symmetry and generates gauge boson masses in right proportions, but the scale is incorrect. It is therefore natural to investigate the possibility of QCD-like mechanism, extended non-abelian gauge symmetry which is broken at a higher scale, which retains the qualitative features of QCD, but provides scaled-up *vev* and therefore gauge boson masses compatible with the experiment. This is precisely the main idea of Technocolor.

5.3 Technicolor toy model

Let us introduce a simplified model of Technicolor where the gauge group is extended by $SU(N)_{TC}$, usually $N = 4$, hence we have extended matter content corresponding to (in simplest version) new particles in fundemental representation of the additional gauge group interacting by exchange of the 'technigluons' (in adjoint representation). We need to further specify the representations of the whole gauge group to set appropriate interactions. Let us do it methodically:

- we construct a theory with a local $SU(N)_{TC} \otimes SU(3)_c \otimes SU(2)_L \otimes U(1)_Y$ gauge symmetry,
- technifermions are chiral doublets of massless color singlets:

$$\left(\begin{array}{c} U \\ D \end{array} \right)_L , U_R, D_R,$$

  i.e. they transform under the weak $SU(2)$ symmetry, they are massless and they do not interact via the usual QCD strong interactions. Also, they are assigned electric charge so that $Q(U) = 1/2$ and $Q(D) = -1/2$.
- ordinary fermions are technicolor singlets, so they do not interact with technigluons.

The theory, as we have contructed it, has two strongly interacting sectors (the usual QCD sector an the technisector), which are only coupled by the EW interactions. Let us now repeat the qualitative analysis we made for the QCD:

- the lagrangian contains the usual kinetic terms with appropriate covariant derivatives:

$$\bar{U}_L i \not{D} U_L + \bar{U}_R i \not{D} U_R + \bar{D}_L i \not{D} D_L + \bar{D}_R i \not{D} D_R, \tag{31}$$

  and posses a global chiral $SU(2)_L \otimes SU(2)_R$ symmetry.
- in analogy with QCD (techni)gluon exchange is attractive and at sufficiently low scale Λ_{TC} (but much higher then Λ_{QCD}) condensates form: $\langle \bar{U}_L U_R \rangle \neq 0$, $\langle \bar{D}_L D_R \rangle \neq 0$ - hence the chiral technicolor symmetry is spontaneously broken:

$$SU(2)_L \otimes SU(2)_R \rightarrow SU(2)_V$$

- three *technipions* appear: $\pi_T^0, \pi_T^\pm$ - the usual Goldstone bosons associated with broken symmetries.

What is crucial is that the spontaneous breaking of chiral symmetry breaks also the electroweak symmetry to EM. The would-be Goldstone bosons become the longitudinal components of W,$Z^{\pm}$, which therefore acquire mass:

$$M_W = \frac{gF_\pi}{2}, \tag{32}$$

where, as mentioned in the section on QCD, F_π is proportional to the vacuum expectation value of the condensate. In order to obtain observable masses we need $F_\pi = 246\, GeV$.

There is in fact a general pattern in what we described above:

- if we take a strongly interacting gauge theory with chiral symmetry breaking $G \to G'$ such that the gauge group $SU(2)_W \otimes U(1)_Y \subset G$ and $U(1)_{EM} \subset G'$ but $SU(2)_W \otimes U(1)_Y \not\subset G'$.
- then breaking chiral symmetry will automatically break electroweak symmetry to electromagnetism.

There is another technical point, we also need that $SU(2)_V \subset G'$, which ensures that the value of the F associated with Z and $W^{\pm}$ is the same and therefore the tree-level relation 13:

$$\frac{M_W^2}{M_Z^2 \cos^2\theta_W} = 1,$$

holds in theory. If the chiral symmetry group is bigger then $SU(2)_L \otimes SU(2)_R$, then breaking it to $SU(2)_V$ will generate additional Goldstone bosons that will not become longitudinal components of Z and $W^{\pm}$ gauge boson. This, of course, is a problem, since we have not seen any massless particles but photon, so those Goldstone bosons must be made massive, in fact heavy above the scales of previous experiments, since we would have seen them otherwise.

Does Technicolor provide the answers we were looking for? Let us sum up our eforts: we have a dynamical mechanism of EW symmetry breaking, where the 'Higgs' is a technifermion condensate. Since Technicolor, like QCD, is asympototically free some of the problems of elementary Higgs model we mentioned before are solved automatically:

- naturalness: masses of all bound technihadrons are of order $\leq \Lambda_{TC}$ and they receive no big corrections.
- triviality: all asymptotically free theories are nontrivial.

However, there are still some weak points:

1. TC introduces more technipions that need to be given mass by some other mechanism,
2. fermions are still massless,
3. we have not approached the question of flavor at all.

6 ETC summed up

To solve the problems mentioned in the previous section yet another model was (very quickly) introduced, known as Extended Technicolor. The idea is to embed the Technicolor gauge group a larger one: $G_{TC} \subset G_{ETC}$ that couples quarks and leptons to technifermions. When the G_{ETC} is spontaneously broken to G_{TC} at scale Λ_{ETC} then quarks and leptons can acquire masses:

$$m \sim \frac{g_{ETC}^2 F_\pi^3}{\Lambda_{ETC}^2}. \tag{33}$$

This seems a very nice solution, but a few problems are immediately visible: the matter content of the ETC is huge. All of those new particles can appear running in loops, thereby introducing quantum corrections. This is however extremely tightly constrained after two decades of precision electroweak measurements - any significant contribution to EW observables is at odds with measured values, so this somehow has to be controlled. Furthermore Flavor Changing Neutral Currents are generated at unacceptably high levels - this is also excluded by precise measurements. Let us see how this comes about.

An exact model of the ETC is not yet available, hence the necessity of working with effective field theory description, i.e. one where the heavy ETC gauge bosons are integrated out (much like Fermi theory was an effective description of electroweak theory). The effective interactions look like this:

$$g_{ETC}^2 \left(\alpha_{ab} \frac{(\bar{T}\gamma_\mu t^a T)(\bar{T}\gamma^\mu t^b T)}{\Lambda_{ETC}^2} + \beta_{ab} \frac{(\bar{T}\gamma_\mu t^a T)(\bar{q}\gamma^\mu t^b q)}{\Lambda_{ETC}^2} + \gamma_{ab} \frac{(\bar{q}\gamma_\mu t^a q)(\bar{q}\gamma^\mu t^b q)}{\Lambda_{ETC}^2} \right) \tag{34}$$

Let us have a look at this expression: the first term involves only technifermions, the second term provides quark masses:

$$\beta g_{ETC}^2 \frac{(\bar{T}T)(\bar{q}q)}{\Lambda_{ETC}^2} \to \beta g_{ETC}^2 \left(\frac{\langle \bar{T}T \rangle}{\Lambda_{ETC}^2} \right) \bar{q}q, \tag{35}$$

the third one is responsible for FCNCs! The kaon system provides constraints on FCNCs, the ETC contribution in this case reads:

$$\frac{g_{ETC}^2 \theta_{sd}^2}{\Lambda_{ETC}^2} (\bar{s}\Gamma^\mu d)(\bar{s}\Gamma_\mu d) + h.c. \tag{36}$$

and gives numerically too high values.

This problem has led in fact to most people rejecting TC/ETC as a viable candidate theory. This judgement might be too fast, the conclusion is however, that simply scaling up QCD in *not enough*. Different dynamics is needed - i.e. different 'running' of the coupling constant. Appropriate modification, called Walking Technicolor, has been constructed: [17], [18], [3]. It is based on an observation, that in most ETC models TC is precociously asymptotically free, i.e. that the Technigluon exchange interaction becomes weak too fast. If TC remains strong from F_π (i.e. the scale where TC becomes strong enough for condensates to appear) up to Λ_{ETC} we will have a different relationship for fermion masses (compare the powers of the energy scales in the equation 33!):

$$m \sim \frac{g_{ETC}^2 F_\pi^2}{\Lambda_{ETC}} \tag{37}$$

This makes a huge difference: since $\Lambda_{ETC} > F_\pi$, to obtain the same quark masses Λ_{ETC} can be higher than before, which in turn supresses FCNCs, as Λ_{ETC}^2 appears in the denominator in front of the FCNC term in 36! Walking Technicolor has the necessary dynamics, i.e. a coupling that stays constant over a large range of energies between F_π and Λ_{ETC} - this is the origin of the name, the coupling 'walks' instead of running. Still the model is not perfect, the problem is the top quark mass, which is extremely heavy and cannot be obtained within the model. A solution has been proposed, 'top assisted technicolor', but we shall not get into the detail of that.

We conclude this discussion with a statement, that the main idea of technicolor is certainly very elegant and much progress has been made in the last two decades in this field (especially taking into account small number of active research in this field compared to, say, supersymmetry), but some problems still remain. There is a large number of models with varying details (much like in supersymmetry), hence exact predictions are rather difficult. Since LHC is coming online, the need for predcitions capturing the qualitative features fo these models became pressing and a toy model dubbed 'Technicolor Straw Man' has been put forward [21], to test it against the experiments (much like MSSM in supersymmetry). We shall see an example in the next section:

7 ETC @LHC

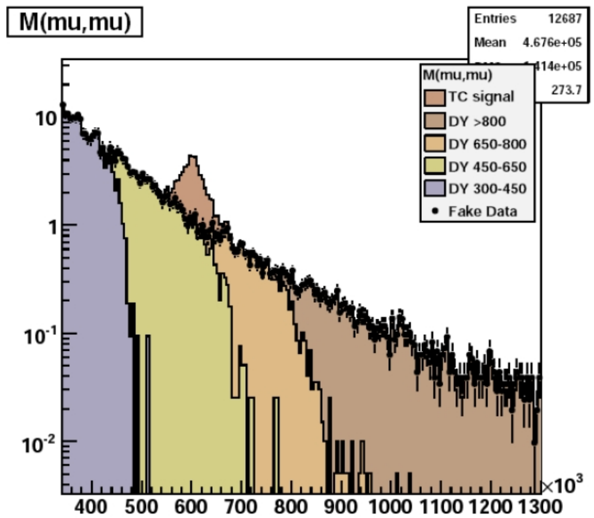

Figure 13: Simulated invariant mass plot with a peak due to technimeson decay. Taken from [23].

Any model needs testable predictions and clear signatures that would allow to distinguish it from various alternatives. One such signature for Technicolor comes from technimeson (i.e. bound states of technifermions) decays [23]:

$$\rho_{TC}, \omega_{TC} \rightarrow \mu\mu,$$

where $\mu\mu$ is a final state. The figure 13 shows result of a simulation using the TC Straw Man model, where a invariant mass plot for decays into this final state was produced. We can see a prominent peak on top of simulated QCD Drell-Yan background i.e. this is a potentially a clean signal.

Other processes have been proposed as good Technicolor signatures [4]:

- $$q\bar{q}' \rightarrow \rho_T^{\pm} \rightarrow V_1 V_2,$$

  where $V_1 V_2 = W^{\pm} Z$, $W^{\pm} \pi_T^0$, $\pi_T^{\pm} Z$, $\pi_T^{\pm} \pi_T^0$

- $$q\bar{q} \rightarrow \rho_T^0 \rightarrow V_1 V_2,$$

  where $V_1 V_2 = W^+ W^-$, $W^{\pm} \pi_T^{\mp}$, $\pi_T^+ \pi_T^-$

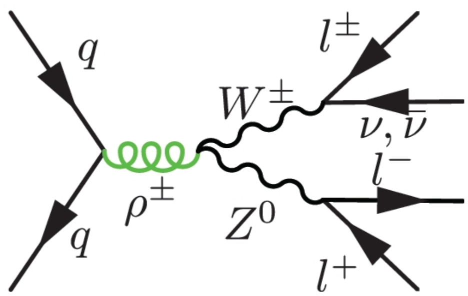

Figure 14: A possible technicolor discovery channel involving a technimeson decay. Taken from [4].

- subsequently $\pi_T^0 \to b\bar{b}$ and $\pi_T^{\pm} \to c\bar{b}$
  - analysis of those dijet technipions decays is harder, but still possible at LHC.

7.1 TC extras

There are additional subjects of interest associated with the Technicolor model, that we will not discuss beyond mentioning them:

- TC particles as Dark Matter candidates - Technicolor vast matter content and heavy masses, as well as interactions with QCD sector only via EW suggest that a DM candidates should not be hard to find and in fact some proposal have been put forward [13, 28, 24, 12]:
- a new source of CP violation - as discussed in spontaneous breaking of CP is possible in Technicolor models, though there are still some unresolved issues. A discussion may be found in [19].
- a possible first order electroweak phase transition in certain technicolor models: [7].

8 Summary

We have reviewed the $SU(2) \otimes U(1)$ electroweak symmetry model and discussed the spontaneous breaking of this symmetry to the $U(1)_{EM}$ and associacted mass generation. We have examined two mechanisms i.e. the elementary Higgs model and a SM extension - Technicolor. After discussions of section 3 it should be clear that the elementary model suffers from serious theoretical deficiencies which limit its validity to a certain energy scale Λ, which - there are good reasons to believe - may be of order as low as $1\ TeV$, and is therefore open to experimental investigations at the LHC. We have been careful to underline the perturbative character of this analysis and the fact that some of our conclusions tell us more about the applicability of perturbation theory to EW theory at high energy scales than the validity of the theory itself. Our review of the technicolor was focused on main ideas and motivations rather then precise phenomenological implications. We have emphasized the importance of the *dynamical* mechanism of symmetry breaking and the relation of technicolor to mass generation built in QCD. Our discussion of problems and shortcomings of this model has led us to various extensions, we tried to convey that many of the original problems of the model have been solved, though open questions still remain. Finally we have mentioned some of the aspects of detection of Higgs or Technicolor at the LHC, emphasizing the need for a detailed analysis beyond the Higgs discovery channel. Interested reader may find the (mostly review) sources we have used as well us some of the original papers in the bibliography.

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