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Seminar

The Higgs boson

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Abstract

The Standard Model (SM) of particle physics developed about 30 years ago has been a major success. With the discovery of neutral currents and heavy vector bosons and prediction of their masses many of the major elements of the theory were demonstrated. But that does not mean that SM theory is complete. For instance, the question of why the W^+, W^- and Z particles that mediate the weak force have mass while the other force carriers, the photon and eight gluons, are massless, remains one of the greatest problems of SM.

In the seminar I will pursue the question of the weak boson and fermion masses with the best known answer we today have - spontaneous symmetry breaking and the Higgs mechanism. The electroweak gauge bosons and the fundamental matter particles are supposed to acquire masses through the interaction with a scalar (Higgs) field.

Contents

1 Introduction

Already in 1964 Peter W. Higgs (University of Edinburgh) formulated the Higgs mechanism [1], which allows us to endow some of the fields with mass while retaining their exact gauge symmetry. Using that Glashow, Weinberg and Salam unified electroweak interactions with W's, Z and H (Higgs particle) in 1967 [2]. In 1973 followed the discovery of neutral currents [3](Gargamelle, CERN), one of the first major successes of SM. A year later Iliopoulos completed the formulation of SM with $SU(2)_L \times U(1)_Y$[4]. Almost ten years later (1983) weak bosons W and Z were discovered [6] at UA1 and UA2 experiments at CERN which again confirmed the SM predictions.

In 1989 LEP and SLC started. Collisions of electrons and positrons were used for precision tests of SM and searches of the Higgs boson began in earnest. New machine was proposed in 1980 – LHC. In 1995 top quark was discovered at Fermilab by CDF and D0[5]. Since year 2000, when LEP stopped running, LHC and it's experiments are being built. Their main purpose is to find Higgs particle and to search for new physics in the 100 GeV – 1 TeV range.

As we can see, SM predicted many things however there are still many questions left unanswered. The purpose of this seminar is to pursue the question of masses which the SM claims to answer: the appearance of mass in the nature. The explanation comprises spontaneous symmetry breaking extended by the Higgs mechanism.

2 Lagrange formalism

The connection between symmetries and conservation laws is best discussed in the framework of Lagrangian formulation of quantum field theory. In order to do that we have to extend the (classical mechanics) Lagrange formalism, that is, a discrete system with coordinates $q_i(t)$, to a system with continuously varying coordinates $\phi(\mathbf{x}, t)$. So

$$L(q_i, \dot{q}_i, t) \longrightarrow \mathcal{L}(\phi, \frac{\partial \phi}{\partial x_\mu}, x_\mu), \tag{1}$$

and the well–known Euler–Lagrange equation becomes

$$\frac{\partial}{\partial x_\mu}\left(\frac{\partial \mathcal{L}}{\partial(\partial \phi/\partial x_\mu)}\right) - \frac{\partial \mathcal{L}}{\partial \phi} = 0. \tag{2}$$

Invariance under translations, time displacements, rotations and Lorentz transformations leads to conservation of momentum, energy and angular momentum. For example, an electron, described by the Dirac equation is invariant under phase transformation

$$\psi(x) \longrightarrow e^{i\alpha}\psi(x), \tag{3}$$

where α is a real constant. The family of such phase transformations forms a unitary Abelian group known as $U(1)$. One may think that $U(1)$ invariance of $\mathcal{L}$ is unimportant, however through Noether's theorem invariance implies the existence of a conserved current, in case of electron (QED) conserved charge. From a physicists point of view, the existence of a symmetry implies that some quantity is unmeasurable (i.e. translation symmetry → absolute position in space cannot be determined). Similarly, invariance under transformations of type Eq. 3 implies that phase α is unmeasurable and as such can be chosen arbitrarily. We speak of global gauge invariance. This gauge is not the most general; it would be more satisfactory if α could depend on the space–time position; $\alpha = \alpha(x)$, which we shall call the local gauge.

One can easily show that the Lagrangian corresponding to an electron, described by the Dirac equation, is given by:

$$\mathcal{L} = i\bar{\psi}\gamma^{\mu}\partial_{\mu}\psi - m\bar{\psi}\psi \tag{4}$$

and is not invariant under a local phase transformation:

$$\psi(x) \longrightarrow e^{i\alpha(x)}\psi(x) \tag{5}$$

since the derivative of ψ does not follow Eq. 5. If we insist on imposing invariance on the Lagrangian, we must define a modified derivative, D_{μ}, that transforms like ψ itself, since the $\partial_{\mu}\alpha(x)$ term breaks the invariance. This can be accomplished by

$$D_{\mu} \equiv \partial_{\mu} - igV_{\mu}, \tag{6}$$

where V_{μ} is vector (gauge) field that transforms as

$$V_{\mu} \longrightarrow V_{\mu} + \frac{1}{g}\partial_{\mu}\alpha. \tag{7}$$

Finally we can write a local gauge invariant Lagrangian corresponding to the Dirac equation as:

$$\mathcal{L} = \bar{\psi}\left(i\gamma^{\mu}\partial_{\mu} - m\right)\psi + g\bar{\psi}\gamma^{\mu}\psi V_{\mu} + \ldots \tag{8}$$

To sum up, by demanding local phase invariance, we were forced to introduce a vector (gauge) field V_{μ}. The dots in Eq. 8 correspond to dynamic terms of the field V_{μ} that are not of primary interest. If we substitute $g \to e$ and $V_{\mu} \to A_{\mu}$ we end up with the Lagrangian of quantum electro–dynamics.

If we are to regard the gauge field V_{μ} as the photon field, we have to add the part $-\frac{1}{4}F^{\mu\nu}F_{\mu\nu}$ to Eq. 8, where $F_{\mu\nu}$ is the field tensor. Note that addition of a mass term $\frac{1}{2}m^2V_{\mu}V^{\mu}$ is prohibited by gauge invariance. The gauge particle, in QED example the photon, is massless.

3 Symbolic Lagrange

Instead of writing Lagrange equations in their full form we will connect the Lagrangian formalism with Feynman rules and hence write out the Lagrangian in symbolical form.

We shall associate the various terms in Lagrangian with a set of propagators and vertex factors. The propagators are determined by the terms quadratic in the field, e.g. the terms in Lagrangian containing ϕ^2, $\bar{\psi}\psi$ and so on. The others terms in the Lagrangian are associated with interaction vertices.

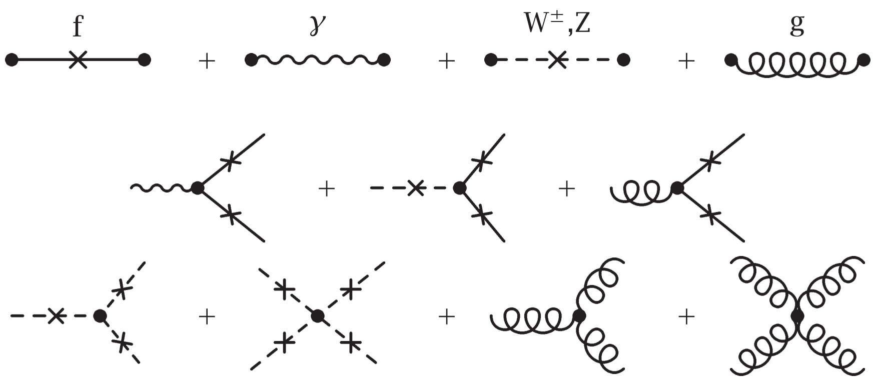

Figure 1: A Lagrangian of the observed interactions.

A Lagrangian corresponding to the observed interactions is presented in Fig. 1. This is the world we live in; actually, the part of it we know how to describe. First line of Fig. **??** symbolically corresponds to kinetic energies of fermions, photons, bosons and gluons (in that order), where crosses on fermion and boson line simply point out that fermions and bosons have mass.

The second line in Fig. 1 corresponds to (massive) fermion interactions with photon, boson and gluon respectively, and the last line in Fig. 1 is showing us third and quartic order self–interaction of (massive) bosons and (massless) gluons. Any possible electroweak and color interactions found in nature can be written down using just these “building blocks”.

Let’s consider just quantum electro–dynamics and weak interactions. The electroweak theory, developed by Glashow, Weinberg and Salam in 1967 [2] describes just a part of the Lagrangian, depictated in Fig. 1. The electroweak ($SU(2)_L \times U(1)_Y$) Lagrangian consists of:

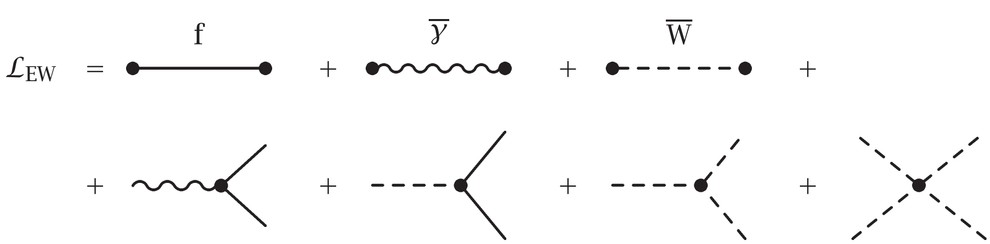

Figure 2: Symbolic $SU(2)_L \times U(1)_Y$ Lagrangian.

where $\overline{\gamma}$ and $\overline{\mathrm{W}}$ are ur–photon and ur–bosons respectively. But Fig. 2 has one problem. It leads to a massless world of fermions and bosons. Both photons and gluons are required to be massless, since the presence of mass terms for gauge fields destroys the gauge invariance of the Lagrangian. So how to apply gauge symmetry to weak interactions, mediated by gauge bosons $W^{\pm}$ and Z, all of them with masses of order of 100 GeV? One could simply "correct" the theory by adding corresponding terms to the Lagrangian and ignoring the symmetry–breaking effect. But that would eventually lead to severe problems which are technically called nonrenormalizable divergences. Instead of that physicists solved a mystery using spontaneous symmetry breaking with the Higgs mechanism.

4 Spontaneous symmetry–breaking

What we shall actually do is redefine our vacuum state. Instead of an "empty" vacuum we want to have a certain kind of field which interacts with fermions and gauge fields and we cannot turn off. It will provide masses to the fermions and will mix ur–photon and ur–bosons in such a way that the photon will remain massless while the bosons will gain mass. This means that we will still have the electroweak symmetry, but it will be hidden (or broken) by our new "non–empty" vacuum.

So rather than to put explicit symmetry–breaking terms into the Lagrangian, which seems artificial and unappealing, we would like to break the electroweak symmetries in a way such that the equations retain their symmetry.

Nature seems to realize this by exploiting the mechanism of spontaneous symmetry breaking; that is, the Lagrangian is invariant under some symmetry, but the symmetry is broken because the vacuum state of the Lagrangian is not invariant. The simplest examples come from solid–state physics, where the phenomenon of spontaneous symmetry breaking is quite common. Consider a ferromagnet, where Lagrangian does not point out any particular direction in space (it is rotationally invariant) the ground (vacuum) state, however, can consist of atoms with spins all aligned along a definite, albeit arbitrary, direction. Thus rotational symmetry can be broken by the vacuum state, even when the Lagrangian remains fully symmetric and there exist infinitely many vacua, i.e. ground states. Another example is the buckling of a rod under axial pressure. The equations are symmetric under rotations about the axis of the rod, yet is buckles in one particular, even though arbitrary, direction. Again there are infinitely many states for the buckled rod.

In both these examples the non–symmetric states correspond to a lower energy than the symmetric ones. The original symmetry of the equations of motion is hidden. It is evident only in our inability to predict in which direction the spins will align or the rod will bend and in the fact that all the non–symmetric solutions are equivalent and can be obtained from one another by a symmetry operation. A critical point exists in both examples, i.e. a critical value of some quantity, either temperature or external force, which will determine whether spontaneous symmetry breaking will occur. Beyond the critical point the vacuum becomes degenerate and the symmetric solution unstable. These properties are typical of all examples of spontaneous symmetry breaking.

Remarkably, we shall find that the above phenomenon allows the construction of a gauge theory in which the underlying symmetry is spontaneously broken, and as a result masses for the weak bosons W and Z as well as for the fermions are "spontaneously generated".

Since the $SU(2)_L \times U(1)_Y$ is relatively complex symmetry it is more illustrative to present the Higgs mechanism on a simpler "toy model" using scalar field and $U(1)$ symmetry. Suppose now that we are dealing with a (complex) scalar field ϕ with the Lagrangian of the form

$$\mathcal{L} = (\partial_\mu \phi)(\partial^\mu \phi^*) - \mu^2 \phi\phi^* - \lambda(\phi\phi^*)^2. \tag{9}$$

Note that real for real field, ϕ, the mass term and the kinetic energy term would each have an extra factor 1/2. In a quantum theory μ^2 would normally be regarded as the (bare) mass of the field quanta and the λ term as a form of self–interaction. $\mathcal{L}$ is invariant under the $U(1)$ group of global transformations (Eq. 3).

The kinetic energy term is positive and can vanish only if ϕ = const. The ground state of the system will be obtained when the value of the constant corresponds to the minimum of the "potential":

$$V(\phi) \equiv \mu^2 \rho + \lambda \rho^2, \tag{10}$$

where $\rho \equiv \phi\phi^*$. Potential V can only have minimum when $\lambda > 0$, which we assume to be the case. Let us however not insist on interpreting μ as a mass and let us consider what happens for distinct cases $\mu^2 > 0$ and $\mu^2 < 0$.

If $\mu^2 > 0$, which is in situation of a massive particle, then V is minimal when $\rho = 0$, i.e. $\phi = 0$, as shown in Fig. 3, and we have a symmetric ground state configuration (the ground state $\phi = 0$ is invariant under Eq. 3). Consequently the option where $\mu^2 < 0$ is the choice we really wish to explore. Now the minimum lies at

$$\rho = -\frac{\mu^2}{2\lambda} \tag{11}$$

which means that there is a whole ring of radius

$$|\phi| = \frac{v}{\sqrt{2}} \equiv \sqrt{\frac{-\mu^2}{2\lambda}} \tag{12}$$

in the complex ϕ plane at each of whose points V is at its minimum value, as shown in Fig. 4.

In this case $\phi = 0$ is an unstable point and any value of ϕ satisfying Eq. 12 will give a true ground state. There are now infinitely many ground states. We see that $\mu^2 = 0$ is the critical transition point between the symmetric solution and the degenerate ground state. From now on we consider only the case $\mu^2 < 0$. Any point on the ring of minima is equivalent since they can all be obtained from any point by applying the transformation from Eq. 3. The lowest energy state has non–zero ϕ and V is now everywhere non–zero constant. We have actually redefined our vacuum, which is not empty anymore, but contains a field that we cannot turn off and which gives us constant non–zero vacuum expectation value v. This procedure is called spontaneous symmetry breaking.

We are concerned with evaluating small perturbations about the energy minimum, so that we should now expand the field variable ϕ not about zero (as we would in "old" vacuum)

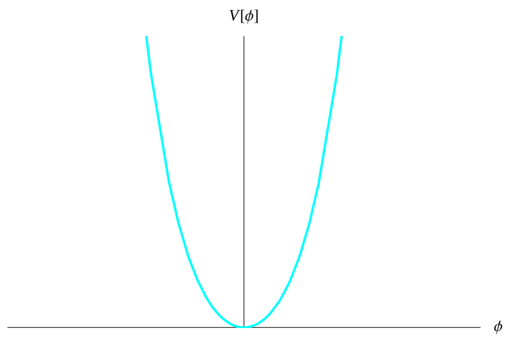

Figure 3: Plot of the potential V as a function of ϕ for $\mu^2 > 0$.

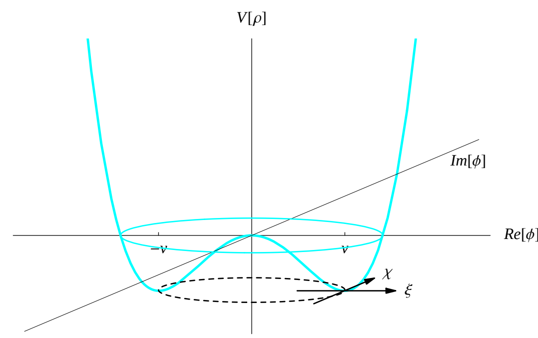

Figure 4: Plot of the potential V as a function of $\phi\phi^*$ for $\mu^2 < 0$.

but around the chosen vacuum v. If we choose this point on the real axis we can write

$$\phi(x) = \sqrt{\frac{1}{2}}\left(v + \xi(x) + i\chi(x)\right), \tag{13}$$

with ξ, χ real and $\xi = \chi = 0$ in the ground state. Substituting into Eq. 9 and ignoring unimportant constant terms, one has

$$\mathcal{L} = \frac{1}{2}\left(\partial_\mu \xi\right)^2 + \frac{1}{2}\left(\partial_\mu \chi\right)^2 - \lambda v^2 \xi^2 - \lambda v \xi\left(\xi^2 + \chi^2\right) - \frac{1}{4}\left(\xi^2 + \chi^2\right)^2. \tag{14}$$

If we were to consider $\mathcal{L}$ as a quantum theory Lagrangian, then it would contain no mass term for the χ field but a normal mass term for the ξ field with

$$m_\xi^2 = 2\lambda v^2 \tag{15}$$

Thus starting with an $\mathcal{L}$ invariant under $U(1)$ constructed from a complex scalar field $\phi(x)$ and the case $\mu^2 < 0$ in the potential we have ended up with a massless field χ and a field ξ whose mass m_ξ has been "spontaneously generated".

Evidently this is not yet what we are searching for. Eventually, we want to give mass to fermions and weak bosons having set out to see whether spontaneous symmetry breaking can cure the disease of unwanted massless vector bosons in $SU(2)_L \times U(1)_Y$ gauge theory, but instead of that we seem to have reached the conclusion that spontaneous symmetry breaking introduces its own massless bosons, so that there appear to be two diseases instead of one.[1] The extraordinary thing is that, taken together, these two problems (massless vector and massless scalar boson) mutually compensate.

5 The Higgs mechanism

In order to demonstrate the "compensation" mechanism mentioned in the previous section it is best to stick to the earlier $U(1)$ toy model. As a further step we replace the global $U(1)$

[1] This is actually a property of a general theorem due to Goldstone, which says that for every broken generator in a spontaneous symmetry breaking there exists a massless scalar (Goldstone) boson.

invariance with a local $U(1)$ gauge invariance. According to the previous discussion (see section 2) we have to replace the derivative ∂_μ by the covariant derivative $D_\mu = \partial_\mu - igV_\mu$ so that $\mathcal{L}$ becomes

$$\mathcal{L} = \left[(\partial_\mu + igV_\mu)\, \phi^* \right] \left[(\partial_\mu - igV_\mu)\, \phi \right] - \mu^2 \phi\phi^* - \lambda \left(\phi\phi^*\right)^2 . \tag{16}$$

It is invariant under the local (Abelian) gauge transformation of Eq. 5 and V_μ is a massless gauge vector boson.

We can once again look for a minimum in the potential, and we find one if $\lambda > 0$. If $\mu^2 < 0$ there is again a ring of degenerate ground states, whereas the symmetric ground state $\phi = 0$ can be obtained if $\mu^2 > 0$. The interesting case is of course the former, and, proceeding as before, we evaluate small perturbations around chosen ground state (Eq. 13). We find

$$\mathcal{L} = \frac{1}{2} g^2 v^2 V_\mu V^\mu + \frac{1}{2} \left(\partial_\mu \xi\right)^2 + \frac{1}{2} \left(\partial_\mu \chi\right)^2 - \lambda v^2 \xi^2 - ev V_\mu \partial^\mu \chi . \tag{17}$$

The term involving V_μ^2 is a great surprise since in a quantum picture it looks as if the gauge field V_μ has acquired mass. Gauge invariance is, of course, still there since Eq. 17 must be equivalent to Eq. 16.

If we look at the structure of $\mathcal{L}$ in Eq. 16 it now seems to describe the interaction of a massive vector field V_μ and two scalars, the massive ξ field and massless χ field. It is instructive to count the degrees of freedom in the two versions of $\mathcal{L}$, Eq. 16 and Eq. 17.

In Eq. 16 there is one massless vector field with two degrees of freedom (corresponding to two transverse modes) and one complex scalar field with two degrees, Fig. 5. In Eq. 17 we have one massive vector boson with three degrees of freedom (longitudinal mode is now allowed) and two real scalar fields, which sums up to five degrees of freedom. We seem to have gained an extra degree, but this is only apparent since we can utilize the gauge invariance to choose a particular gauge in which χ simply does not appear.

Since the theory does not change with any choice of the transformation function $\theta(x)$ in Eq. 5, let us choose $\theta(x)$ at each space–time point to equal the phase of $\phi(x)$. Then in this gauge we write expansion of ϕ around the minima of the form

$$\phi(x) = \sqrt{\frac{1}{2}} \left(v + h(x) \right) , \tag{18}$$

where $h(x)$ is real. Now we can finally write our Lagrangian

$$\mathcal{L} = \frac{1}{2} \left(\partial_\mu h\right)^2 - \lambda v^2 h^2 - \lambda v h^3 - \frac{1}{4} \lambda h^4 + v e^2 V_\mu^2 h + \frac{1}{2} e^2 v^2 V_\mu^2 + \frac{1}{2} e^2 V_\mu^2 h^2 . \tag{19}$$

In this form Lagrangian $\mathcal{L}$ describes the interaction of the massive vector boson V_μ with the massive, real, scalar field h (called the “Higgs boson”), whose squared mass is given by

$$2\lambda v^2 = -2\mu^2 \tag{20}$$

All massless particles have completely disappeared and the number of degrees of freedom is back to four, as it should be. What has happened is that in the spontaneously broken symmetry the gauge boson has acquired mass at the expense of the would–be Goldstone boson,

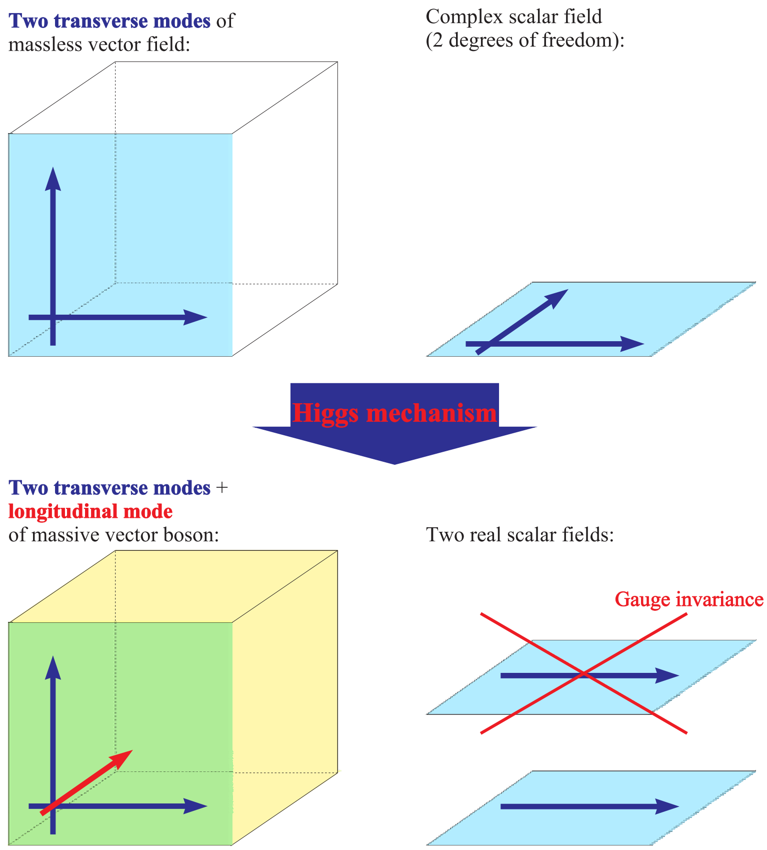

Figure 5: Counting degrees of freedom.

which simply disappears. This is called the Higgs mechanism. For each vector gauge field that gets massive we need one complex scalar field, one piece of which becomes unphysical and disappears (it reappears as the longitudinal mode of the vector field) leaving one real scalar physical field, the Higgs boson.

Three things should be emphasized at this point: First there is a real boson, h, that should occur as a physical boson – the Higgs boson. Secondly, its mass depends on λ and on v. The gauge boson mass determines v, but λ is a parameter, characteristic of the scalar potential and it can not be calculated but has to be determined by the experiment. Therefore, the Higgs mass is unknown. Thirdly, the interaction terms (plus those that occur when fermions are given mass) determine the production mechanisms and decays of the Higgs boson. The self–interaction terms depend on λ but the terms describing the interaction of h with V_μ do not depend on λ, so their strength is known.

Of course the simple $U(1)$ model was used for demonstration purposes only, since it clearly gives an unphysical massive photon. The much more complex $SU(2)_L \times U(1)_Y$ symmetry however gives us a more realistic outcome. It is not so tedious to write down the electroweak model spontaneously broken and corrected with Higgs mechanism if one uses the symbolical Lagrange introduced in 3. Diagrammatic form of Lagrangian is then

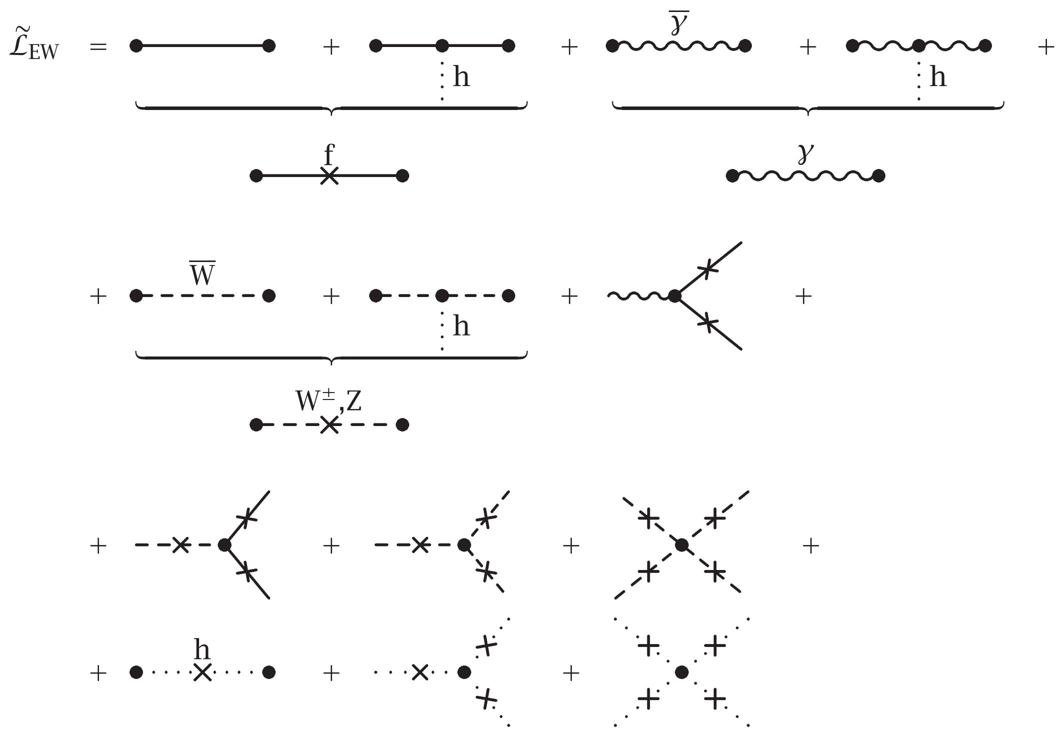

Figure 6: The diagrammatic form of the Lagrangian after the symmetry breaking and Higgs mechanism.

Comparison of Fig. 5 and Fig. 2 shows some additional developments. First of all, in electroweak lagrangian (Fig. 2) we had ur–photon and ur–bosons (three $\overline{\mathrm{W}}$'s and $\overline{\gamma}$), which were massless. After the symmetry breaking they gain mass and since they are massive, they mix, and after mixing we end up with massless photon, γ, and three massive vector bosons, W^+ W^- and Z, which is in agreement with the observed state of Fig. 1. At this point it should be stressed that coupling of fermions and weak bosons to the Higgs field is proportional to their masses.

6 Higgs mass

From Lagrangian Eq. 19 we can write for the weak boson ($\mathrm{W}^\pm$) mass

$$M_{\mathrm{W}} = \frac{1}{2} g v \tag{21}$$

which gives us $v = 246$ GeV at $M_{\mathrm{W}} = 80.4$ GeV. Consequently, as previously stated, v is fixed, while λ remains uncertain. Hence we cannot predict Higgs mass, which is $M_{\mathrm{H}} = 2\sqrt{-\lambda v^2}$.

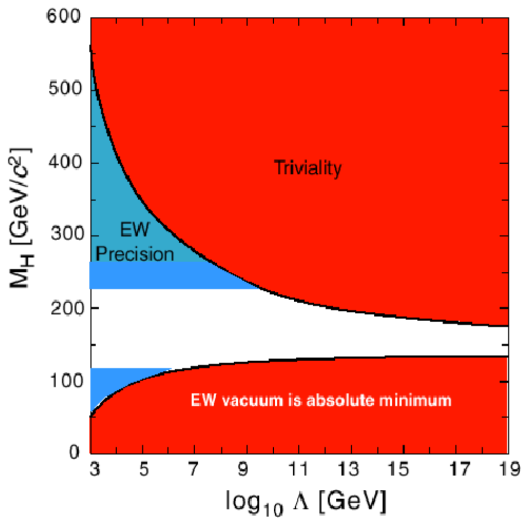

Figure 7: Higgs mass constraints from vacuum stability and Higgs self–coupling.

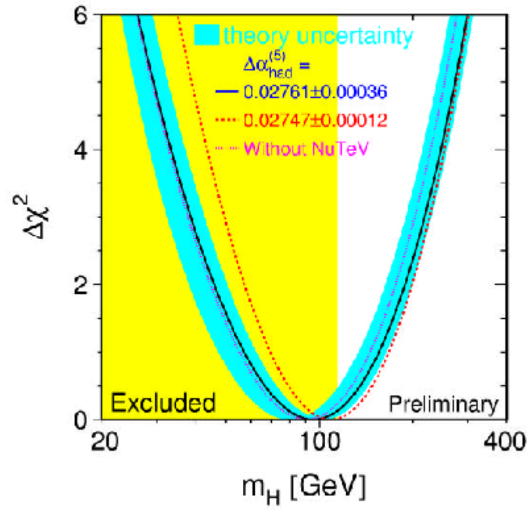

Figure 8: Higgs mass bounds from LEP.

Other more accurate (but still indirect) bounds are obtained from different experiments, LEP (Large Electron Positron collider [15], CERN) among them. The only direct bound is the center–of–mass energy of LEP, because Higgs was not observed at LEP. Very rare production of a Higgs particle could happen in a process similar to diagram c) in 7.1, with the exception of electrons replacing quarks. Since Z mass is known (91.2 GeV) and highest CMS energy achieved at LEP was around 205 GeV, Higgs mass has to be larger that the difference (∼ 114 GeV). Current bounds on Higgs mass at 95% confidence level today are

$$114.4 \text{ GeV} < M_{\mathrm{H}} < 211 \text{ GeV}, \tag{22}$$

where the lower mass bound is defined by LEP measurements and upper limit by calculations from higher order corrections. At this point is should however be stressed that the SM is only one of possible SSM scenario and that other possible generally have one ore more Higgs bosons with mass up to 1 TeV.

7 Higgs production and decay at LHC

7.1 Production

The most realistic option of discovering a Higgs particle will be at the LHC (Large Hadron Collider [14], CERN), so it might be instructive to look at the production and decay relations in this environment. The dominant Higgs production mechanism at the LHC for all possible Higgs masses will be the gluon fusion process. Other processes with their Feynman diagrams in Table 7.1 are also of interest because of the special signatures they can provide for the identification of the Higgs. In Fig. 9 the cross section is shown as a function of the Standard model Higgs mass. At the highest masses a significant part of the cross section is from vector boson fusion.

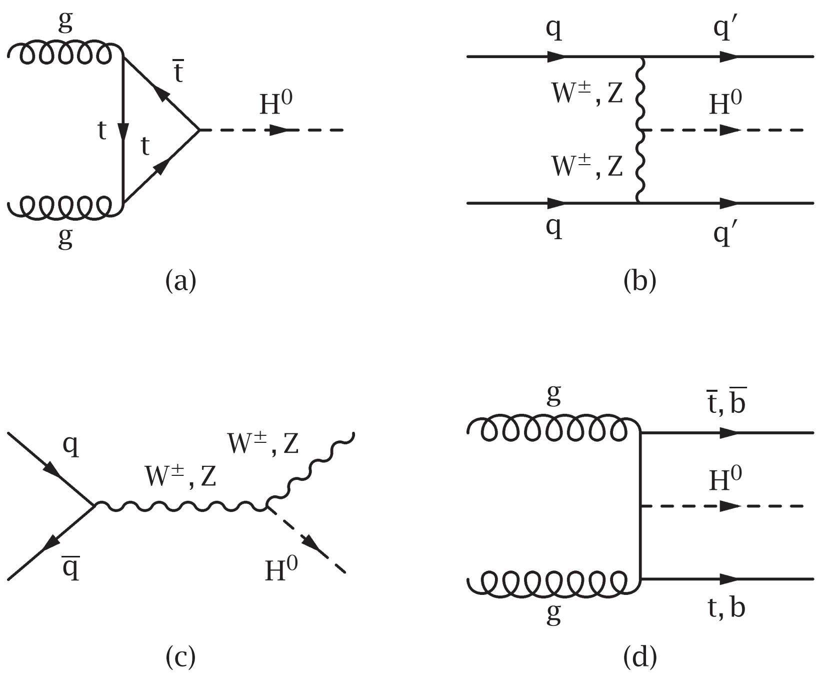

Table 1: The most important processes for Higgs production at hadron colliders. Gluon fusion (a), vector boson fusion (b), associated production with weak bosons (c) and an example of the diagram having associated production with a t or b pair (d).

7.2 Decay

The total decay width and lifetime, as well as branching ratios for specific decay channel, are determined by the strength of the couplings of the Higgs boson to fermions and weak bosons. The measurements of the decay characteristics can therefore experimentally show that the Higgs couplings grow with the masses of the particles, which is a direct consequence of the Higgs mechanism.

To find the so called intermediate Higgs particle which mass is below the $H \to ZZ$ decay threshold and above the limit set by the searches at LEP (Eq. 22), will be difficult. The obvious way to detect a Higgs would be in the dominant $H \to b\bar{b}$ channel but with the b–quarks fragmenting into jets this will be overwhelmed by the di–jet rate. Also the $H \to b\bar{b}$ lacks any trigger as it neither has high jet energy nor isolated leptons in the final state. A more favorable situation can be obtained either by looking at associative production or at one of the rarer decays.

With the Higgs produced together with either a top quark or a vector boson the problem

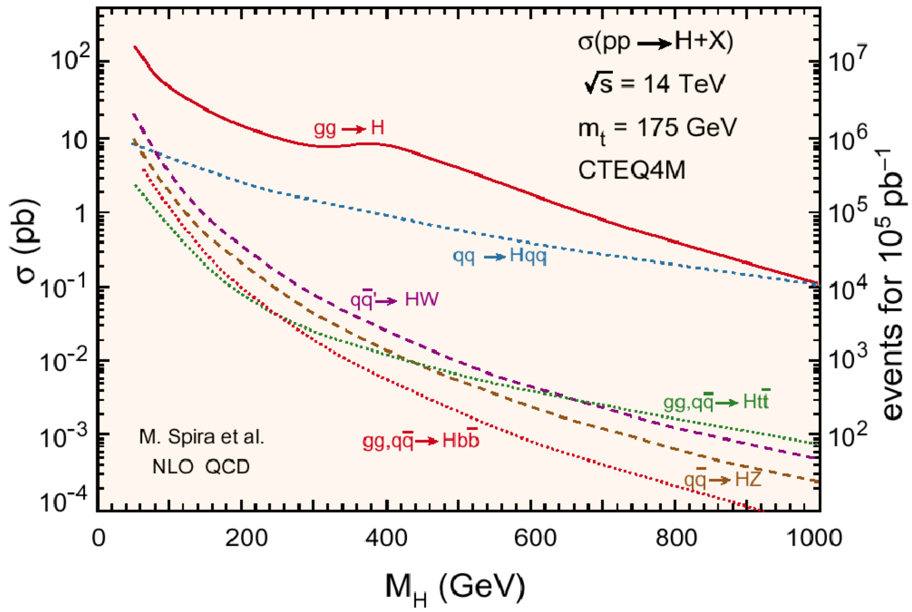

Figure 9: The production cross section of the standard model Higgs boson. Across the complete possible mass range the gluon fusion process is dominating.

of getting a trigger for the Higgs event are solved by requiring a high energy lepton from one of the top quarks or the vector boson decay. The HZ production is already suppressed in comparison with the HW and taking the factor three lower branching ratio for leptonic decays into account the HZ production mode is of limited interest as the rate will be low.

The next handle for the decay is to identify the jets with b–quarks. The method called b–tagging is based either on the long lifetime if the b-quarks which causes secondary vertices or high amount of leptons in B–meson decays. While the HW mode will in general have two b–quarks in the final state the $Ht\bar{t}$ will have four as $t \to b$. The use of multiple b–tagging can provide a larger rejection of jets and thus counteract the lower cross section for $Ht\bar{t}$ production in the interesting mass range.

The $H \to b\bar{b}$ decay gives further problems in the reconstruction. The Higgs mass has to be reconstructed from two jets giving trouble with invisible energy from escaping neutrinos and energy lost outside the jet cone. As a result the reconstructed mass peak will be wide. The $Ht\bar{t}$ channel also suffers from the combinatorial problem of selecting the correct combination of b–jets.

The other way of identifying a Higgs in the intermediate mass region is so select an exotic (= very rare) decay as the $H \to \gamma\gamma$ decay. The trigger in this case is two isolated electromagnetic clusters. While the channel suffers from a branching ratio around 10^{-3} the backgrounds are also much lower than in the case of the $H \to b\bar{b}$ decay due to clear signature of two isolated photons in the final state. The main backgrounds are from direct photon production and jets faking photons.

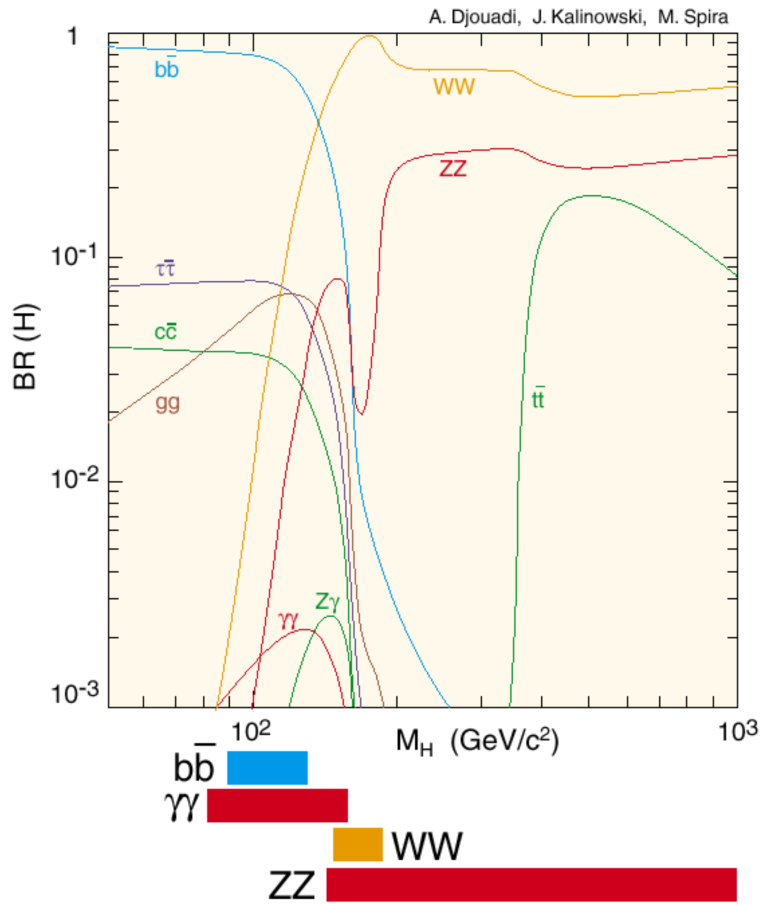

Figure 10: Branching ratios of the dominant decay modes of the standard model Higgs particle. All relevant high–order corrections are taken into account.

In the upper limit of the mass interval with a Higgs mass above 130 GeV the branching ratio to vector bosons reaches significant levels. However, at least one of the vector bosons will not be on the mass shell. The obvious decay channel is $H \to ZZ^* \to 4\mu(e)$ where a mass constraint can be made on one of the lepton pairs while the other pair is given to have an invariant mass below m_Z. Looking on Fig. 10 a funny shape of the branching ratio to Z–bosons can be seen for Higgs masses in the interval 150 GeV $< m_H <$ 190 GeV. It is caused by a threshold effect where the decay to two W bosons on the mass shell turns possible but still at least one of the Z bosons needs to be below the mass shell.

The main irreducible background for the $H \to ZZ^*$ decay is direct ZZ^* and $Z\gamma^*$ production with decay to four leptons. A good mass resolution is required to reduce this continuum background. The most important reducible backgrounds are $t\bar{t}$ and $Zb\bar{b}$ production again with four leptons in the final state. The main cuts to reduce the background is isolated electrons, a mass cut on one of the lepton pairs to the Z mass, and a requirement for the other lepton pair to have an invariant mass above 20 GeV. The last cut mainly reduces the background from the $Z\gamma^*$ and $Zb\overline{b}$ processes.

If a SM Higgs is having a mass above twice the Z mass the discovery will be easy through the decay channel $H \to ZZ \to 4\mu(e)$. This is called the gold plated channel for Higgs decays. Both lepton pairs will have an on–shell mass making it possible to reduce many types of backgrounds.

The main irreducible background is direct ZZ production, but a requirement for at least one of the Z bosons to have a transverse momentum above half the Higgs mass will strongly suppress this background. The upper mass limit for detecting the Higgs in this decay channel is given by the reduced production rate and the increased width of the Higgs. A larger width of the signal increases the irreducible continuum background, but the reduced rate is the most serious problem. As an example fewer the 200 Higgs particles with $m_H = 700$ GeV decay in the $H \to ZZ \to 4\mu(e)$ channel in a year at high luminosity. Taking into account the kinematic cuts and detecting efficiencies hardly any particles are left for detection.

With the fixed collision energy of the LHC the production cross section of a Higgs particle falls with an increasing Higgs mass. The rate in a selective decay channel like the four lepton channel is thus no longer high enough for the highest Higgs masses. With the decay to vector bosons totally dominating, the only possible detection channels lest are the ones with at least on of the vector bosons decaying to neutrinos or jets.

8 Conclusions

So we have seen the spontaneous symmetry breaking and Higgs mechanism on work. They are the cornerstone in the electroweak sector of the standard model, where the electroweak gauge bosons and the fundamental matter particles are supposed to acquire masses through the interaction with a scalar (Higgs) field.

The Higgs I have shown is not the only one, just as our choice of the Higgs field is not the only one. It is merely the simplest one that already solves our problems of generating masses. There are other theories that incorporate different Higgs fields and end up with more Higgs particles. Supersymmetric standard model (SUSY) is one of them and Minimal supersymmetric standard model predicts as many as five Higgs particles.

What will new experiments show? Nobody knows the answer. But one thing for sure: if not Higgs, then some other physics will be employed at those energies. And physicists quest for The Theory will continue.

And the Lord looked upon Her world, and She marveled at its beauty - for so much beauty there was that She wept. It was a world of one kind of particle and one force carried by one messenger who was, with divine simplicity, also the one particle.

- The Very New Testament 2:1

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