



The BEH Mechanism and its Scalar Boson

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by François Englert

Université Libre de Bruxelles, Brussels, Belgium

1 INTRODUCTION: SHORT AND LONG RANGE INTERACTIONS

Physics, as it is conceived today, attempts to interpret diverse phenomena as particular manifestations of testable general laws. Since its inception in the Renaissance, mainly through Galileo's revolutionary concepts, this has been an extraordinarily successful adventure—to the point where after impressive developments in the first half of the twentieth century, one might have even conceived that all phenomena, from the atomic scale to the edge of the visible universe, are governed solely by two fundamental laws, and two known laws. Namely classical general relativity, Einstein's generalisation of Newtonian gravity, and quantum electrodynamics, the quantum version of Maxwell's electromagnetic theory.

Gravitational and electromagnetic interactions are long range interactions, meaning they act on objects no matter how far they are separated from each other. The progress in the understanding of such physics applicable to large scales is certainly attributable to the fact they can be perceived without the mediation of highly sophisticated technical devices. But the discovery of subatomic structures had revealed the existence of other fundamental interactions that are short range, that is, negligible at larger distance scales. In the early 1960s, there was no consistent theoretical interpretation of short range fundamental interactions, nor of the “weak interactions” responsible for radioactive decay, nor of the “strong interactions” responsible for the formation of nuclear structures.

Robert Brout and I [1], and independently Peter Higgs [2], constructed a mechanism to describe short range fundamental interactions. Robert Brout

passed away in 2011 and left me alone to tell our story. I will explain how we were led to propose the mechanism, and how it allows for consistent fundamental theories of short range interactions and for building elementary particle masses. It became a cornerstone of the Standard Model and was recently confirmed by the magnificent discovery at CERN of its predicted scalar boson.

We became convinced that a consistent formulation of short range interactions would require a common origin for both short and long range interactions.

While both classical general relativity and quantum electrodynamics describe long range interactions and are both built upon very large symmetries, labelled “local symmetries,” they have very different structures: in contradistinction to general relativity, long range quantum electrodynamics is fully consistent at the quantum level and was experimentally verified at that level, in particular by the successful inclusion of chemistry in the realm of known physics. As a valid theory of short range interactions clearly required quantum consistency, we were naturally driven to take, as a model of the corresponding long range interactions, the generalisation of quantum electrodynamics, known as Yang-Mills theory.

The quantum constituents of electromagnetic waves are “photons,” massless neutral particles travelling with the velocity of light. Their massless character implies that the corresponding waves are polarised only in directions perpendicular to their propagation. These features are apparently protected by local symmetry, as the latter does not survive the explicit inclusion of a mass term in the theory. Yang-Mills theory is built upon similar local symmetries, enlarged to include several massless interacting quantum constituents, neutral and charged ones. These massless objects are labelled gauge vector bosons (or often simply gauge bosons).

To transmute long range interactions into short range ones in the context of Yang-Mills theory it would suffice to give these generalised photons a mass, a feature that, as we just indicated, is apparently forbidden by local symmetries. Momentarily leaving aside this feature, let us first recall why massive particles transmit in general short-range interactions.

Figure 1 is a Feynman diagram whose intuitive appearance hides a precise mathematical content. Viewing time as running from bottom to top, it describes the scattering of two electrons resulting from the exchange of a massive particle labeled Z of mass m_Z . Classically such process could not occur, as the presence of the Z particle would violate energy conservation. Quantum mechanically it is allowed if the violation takes place within a time span of the order \hbar/mc^2 . This process then describes in lowest order perturbation theory a short-range interaction cut-off at a range $\sim \hbar/mc$.

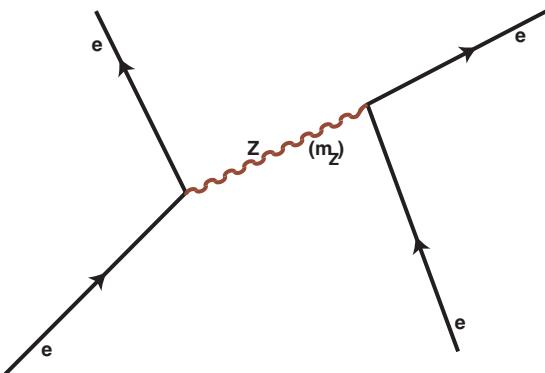


FIGURE 1. Massive particle mediating short-range interactions.

As local symmetries apparently prevent the introduction of massive gauge bosons in the theory, we turn our attention to a class of theories where the state of a system is asymmetric with respect to the symmetry principles that govern its dynamics. This is often the case in the statistical physics of phase transitions [3]. This is not surprising, since more often than not energetic considerations dictate that the ground state or low lying excited states of a many body system become ordered. A collective variable such as magnetisation picks up expectation value, which defines an order parameter that otherwise would vanish by virtue of the symmetry encoded in the formulation of the theory (isotropy in the aforementioned example). This is an example of Spontaneous Symmetry Breaking (SSB) which frequently occurs in the statistical theory of second order phase transitions. Could mass of gauge bosons arise through a similar SSB? This question arises naturally from the seminal work of Yoichiro Nambu, who showed that SSB could be transferred from the statistical theory of phase transitions to the realm of relativistic quantum field theory [4–6], the mathematical framework designed to analyse the world of elementary particles.

This raises a deeper question: could SSB be the agent of the transmutation of long range interactions mediated by massless gauge fields to short range interactions mediated by massive ones, without impairing the validity of the quantum behaviour that characterise the simplest Yang-Mills theory, namely quantum electrodynamics?

As we shall see, the answer is yes to both questions *provided that the notion of SSB is traded for a more subtle one: the BEH mechanism* [1, 2]. To prepare for the discussion of the mechanism, I will first review how SSB can be transferred from the theory of phase transitions to relativistic quantum field theory.

2 SPONTANEOUS SYMMETRY BREAKING

2.1 Spontaneous symmetry breaking in phase transitions

Consider a condensed matter system, whose dynamics is invariant under a continuous symmetry. As the temperature is lowered below a critical one, the symmetry may be reduced by the appearance of an ordered phase. The breakdown of the original symmetry is always a discontinuous event at the phase transition point, but the order parameters may set in continuously as a function of temperature. In the latter case the phase transition is second order. Symmetry breaking by a second order phase transition occurs in particular in ferromagnetism, superfluidity and superconductivity.

I will first discuss the ferromagnetic phase transition, which illustrates three general features of the SSB which set in at the transition point in the low temperature phase: ground state degeneracy, the appearance of a “massless mode” when the dynamics is invariant under a *continuous* symmetry, and the occurrence of a “massive mode” characterising the rigidity of the order parameter.

In the absence of external magnetic fields and of surface effects, a ferromagnetic substance below the Curie point displays a global orientation of the magnetisation, while the dynamics of the system is clearly rotation invariant; namely, the Hamiltonian of the system is invariant under the full rotation group. This is SSB.

A ferromagnetic system is composed of microscopic atomic magnets (in simplified models such as the Heisenberg Model these are spin 1/2 objects) whose interactions tend to orient neighbouring ones parallel to each other. No global orientation appears at high temperature where the disordering thermal motion dominates. Below a critical “Curie temperature” energy considerations dominate and the system picks up a global magnetisation. The parallel

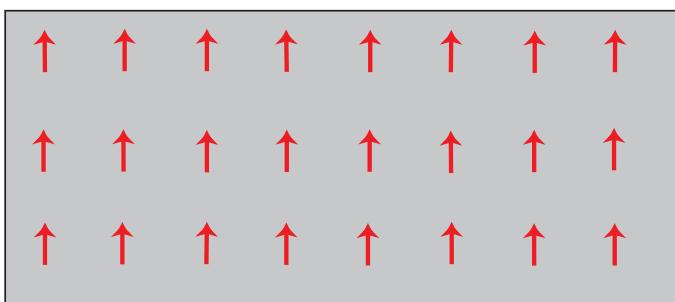


FIGURE 2. Classical representation of a ferromagnet ground state.

orientation of neighbouring magnets propagates, ending up in a macroscopic magnetisation. This selects a direction, which for an infinite isolated ferromagnet is arbitrary. It is easily proven that for an infinite system any pair of possible orientations defines orthogonal ground states and any local excitations on top of these ground states are also orthogonal to each other. Thus the full Hilbert space of the system becomes split into an infinity of disjointed Hilbert spaces. This is ground state degeneracy (Figure 2).

The effective thermodynamical potential V , whose minimum yields magnetisation in the absence of an external magnetic field, is depicted in Figure 3. Above the Curie point T_C the magnetisation \bar{M} vanishes. Below the Curie point the potential develops in a plane $V M_z$ a double minimum which generate a valley in the M_x, M_y directions. Each point of the valley defines one of the degenerate ground states with the same $|\bar{M}|$.

At a given minimum, say, $\bar{M} = M^z \vec{1}_z$, the curvature of the effective potential measures the inverse susceptibility which determines the energy for infinite wavelength fluctuations. This is the analogue of mass in relativistic particle physics. The inverse susceptibility is zero in directions transverse to the order parameter and positive in the longitudinal direction. One thus obtains from the transverse susceptibility a “massless” transverse mode characteristic of broken continuous symmetry: these are the “spin-waves” whose quantum constituents are interacting bosons called “magnons.” The longitudinal susceptibility yields a (possibly unstable) “massive” longitudinal mode which corresponds to fluctuations of the order parameter. In contradistinction to the massless mode which exists only in continuous SSB for which there is a valley, the massive mode is present in any SSB, continuous or discrete, and measures the rigidity of the ordered structure.

The structure of Figure 3 is common to many second order phase transitions and leads to similar consequences. However, in superconductivity a new

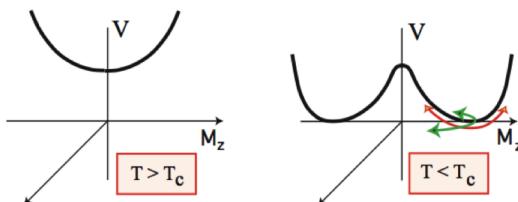


FIGURE 3. Effective thermodynamical potential of a ferromagnet above and below the Curie point.

phenomenon occurs. The quantum phase symmetry is broken by a condensation of electron pairs bounded by an attractive force due to phonon exchange in the vicinity of the Fermi surface. The condensation leads to an energy gap at the Fermi surface. For neutral superconductors, this gap would host a massless mode and one would recover the general features of SSB. But the presence of the long-range Coulomb interactions modifies the picture. The *massless mode disappears*: it is absorbed in electron density oscillations, namely in the “massive” plasma mode. As will be apparent later, this is a precursor of the BEH mechanism [7, 4, 8].

2.2 Spontaneous symmetry breaking in field theory

Spontaneous symmetry breaking was introduced in relativistic quantum field theory by Nambu in analogy with the BCS theory of superconductivity [4]. The problem studied by Nambu [5] and Nambu and Jona-Lasinio [6] is the spontaneous breaking of the $U(1)$ symmetry of massless fermions resulting from the arbitrary relative (chiral) phase between their decoupled right and left constituent neutrinos. Chiral invariant interactions cannot generate a fermion mass in perturbation theory but may do so from a (non-perturbative) fermion condensate: the condensate breaks the chiral symmetry spontaneously. Nambu [5] showed that such spontaneous symmetry breaking is accompanied by a massless pseudoscalar. This is interpreted as the chiral limit of the (tiny on the hadron scale) pion mass. Such an interpretation of the pion constituted a breakthrough in our understanding of strong interaction physics. The massless pseudoscalar is the field-theoretic counterpart of the “massless” spin-wave mode in ferromagnetism. In the model of reference [6], it is shown that SSB also generates a massive scalar boson which is the counterpart of the “massive mode” measuring in phase transitions the rigidity of the order parameter in the spontaneously broken phase.

The significance of the massless boson and of the massive scalar boson occurring in SSB is well illustrated in a simple model devised by Jeffrey Goldstone [9]. The potential $V(\phi_1, \phi_2)$ depicted in Figure 4 has a rotational symmetry in the plane of the real fields (ϕ_1, ϕ_2) , or equivalently is invariant under the $U(1)$ phase of the complex field $\phi = (\phi_1 + i\phi_2)/\sqrt{2}$. This symmetry is spontaneously broken by the expectation value $\langle \phi \rangle$ of the ϕ -field acquired at a minimum of the potential in some direction of the (ϕ_1, ϕ_2) plane, say $\langle \phi_1 \rangle$. Writing $\phi = \langle \phi \rangle + \varphi$

$$\phi_1 = \langle \phi_1 \rangle + \varphi_1, \quad (2.1)$$

$$\phi_2 = \varphi_2. \quad (2.2)$$

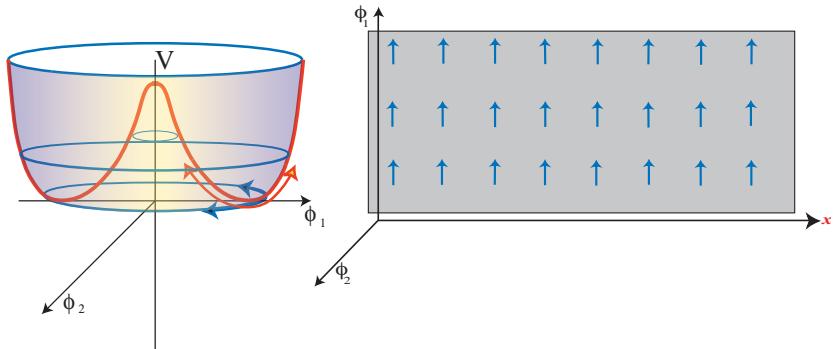


FIGURE 4. Spontaneous symmetry breaking in the Goldstone model

For small ϕ_1 and ϕ_2 we may identify the quantum fluctuation φ_1 climbing the potential as the massive mode measuring the rigidity of the SSB ground state selected by $\langle \phi_1 \rangle$, and the quantum fluctuation φ_2 in the orthogonal valley direction as the massless mode characteristic of a continuous SSB.

Their significance is illustrated in Figure 5 and Figure 6 depicting respectively classical φ_2 and φ_1 wave modes, on the classical background $\langle \phi_1 \rangle$. The corresponding massless and massive bosons are the quantum constituents of these waves.

Figure 5 (a) represents schematically a lowest energy state (a “vacuum”) of the system: a constant non-zero value of the field $\phi_1 = \langle \phi_1 \rangle$ pervades space-time. Figure 5 (b) depicts the excitation resulting from the rotation of half the fields in the (ϕ_1, ϕ_2) plane. This costs only an energy localised near the surface separating

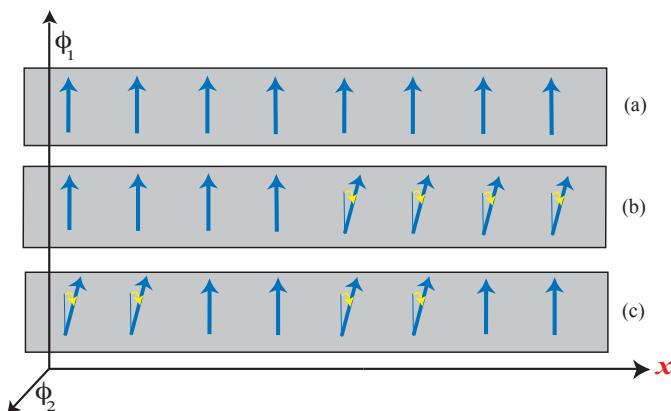


FIGURE 5. Massless Nambu-Goldstone mode φ_2 .

the rotated fields from the chosen vacuum. SSB indeed implies that rotating all the fields would cost no energy at all: one would merely trade the initial chosen vacuum for an equivalent one with the same energy. This is the characteristic *vacuum degeneracy* of SSB. Figure 5 (c) mimics a wave of φ_2 . Comparing 5 (c) with 5 (b), we see that as the wavelength of the wave increases indefinitely, its energy tends to zero, and may be viewed as generating in that limit a motion along the valley of Figure 4. Quantum excitations carried by the wave reach thus zero energy at zero momentum and the mass m_{φ_2} is zero. Figure 5 can easily be generalised to more complex spontaneous symmetry breaking of continuous symmetries. Massless bosons are thus a general feature of such SSB already revealed by Nambu's discovery of the massless pion resulting from spontaneous chiral symmetry breaking [5]. They will be labelled massless Nambu-Goldstone (NG) bosons. Formal proofs corroborating the above simple analysis can be found in the literature [10].

Figure 6 depicts similarly a classical wave corresponding to a stretching of the vacuum fields. These excitations in the φ_1 direction describe fluctuations of the order parameter $\langle \varphi_1 \rangle$. They are volume effects and their energy does not vanish when the wavelength becomes increasingly large. They correspond in Figure 4 to a climbing of the potential. The quantum excitations φ_1 are thus now massive. These considerations can be again extended to more general SSB (even to discrete ones) to account for order parameter fluctuations. Lorentz invariance imposes that such massive excitations are necessarily scalar particles. They were also already present in reference [6] and will be denoted in general as *massive scalar bosons*.

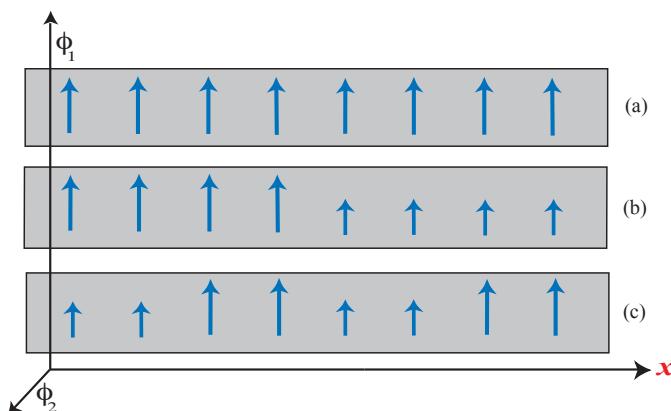


FIGURE 6. Massive scalar mode φ_1 .

To summarise, ϕ_2 describes massless bosons, ϕ_1 massive ones, and the “order parameter” $\langle \phi_1 \rangle$ may be viewed as a condensate of ϕ_1 bosons.

3 THE BEH MECHANISM

The above considerations are restricted to spontaneous symmetry breaking of global continuous symmetries. Global means that the symmetry operations are independent of the space-time point x . For instance in the Goldstone model, the global rotations of the fields in Figure 5 (a) in the (ϕ_1, ϕ_2) plane by angles independent of the space-time point x are symmetries of the theory (they describes motion in the valley of Figure 4): these rotations cost no energy and simply span the degenerate vacua. We will now discuss the fate of SSB when the global symmetry is extended to a local one.

3.1 The fate of the Nambu-Goldstone boson and vector boson masses

We extend the $U(1)$ symmetry of the Goldstone model from global to local. Thus the rotation angle in the (ϕ_1, ϕ_2) plane in Figure 5, or equivalently the rotation in the valley of Figure 4, can now be chosen independently at each space-time point (x) with no cost of energy and no physical effect. To allow such feature, one has to invent a new field whose transformation would cancel the energy that such motion would generate in its absence. This is a “gauge vector field” A_μ . It has to be a vector field to compensate energy in all space directions and it has to transform in a definite way under a rotation in the (ϕ_1, ϕ_2) plane: this is called a gauge transformation and results in a large arbitrariness in the choice of the A_μ field corresponding to arbitrary “internal” rotations at different points of space. The consequence of this gauge symmetry is that the waves are polarised in directions perpendicular to their direction of propagation and that their quantum constituents have to be introduced as massless objects.

Local $U(1)$ symmetry is the simplest gauge field theory and is the symmetry group of quantum electrodynamics. In the local generalisation (the gauging) of the Goldstone model, the introduction of the potential of Figure 4 will deeply affect the “electromagnetic potential” A_μ .

As in the Goldstone model of Section 1.2, the SSB Yang-Mills phase is realised by a non vanishing expectation value for $\phi = (\phi_1 + i\phi_2)/\sqrt{2}$, which we choose to be in the ϕ_1 -direction. Thus

$$\phi = \langle \phi \rangle + \varphi, \quad (3.1)$$

with $\phi_1 = \langle \phi_1 \rangle + \varphi_1$ and $\phi_2 = \varphi_2$. As previously φ_2 and φ_1 appear to describe a NG massless boson and a massive scalar boson.

However a glance on Figure 5 depicting the NG mode immediately shows that Figure 5 (b) and Figure 5 (c) differ from Figure 5 (a) only by local rotations and hence in the local Goldstone model they are just symmetry (or equivalently gauge) transformations. They cost no energy and therefore the NG boson has disappeared: the corresponding fluctuations in the valley are redundant (gauge transformed) descriptions of the same gauge invariant vacuum. It is easy to see that this argument remains valid for any local symmetry and hence *Nambu-Goldstone bosons do not survive the gauging of a global SSB to a local symmetry*. The vacuum is no longer degenerate and strictly speaking there is no spontaneous symmetry breaking of a local symmetry. The reason why the phase with non vanishing scalar expectation value is often labelled SSB is that one uses perturbation theory to select at zero gauge field coupling a scalar field configuration from global SSB; but this *preferred* choice is only a convenient one.

The disappearance of the NG boson is thus an immediate consequence of local symmetry. The above argument [11] was formalised much later [12] but formal proofs not directly based on the gauge invariance of the vacuum were already presented in 1964 [13, 14].

One may now understand in qualitative terms the consequence of the disappearance of the NG boson. Clearly, one does not expect that the degrees of freedom carried by the NG φ_2 field could vanish. As the NG boson disappears because of its coupling to the gauge field, one expects that these degrees of freedom should be transferred to it. This can only occur by adding to the transverse polarisation of the gauge field a longitudinal one. But such polarisation is forbidden as mentioned earlier, for a massless field. Therefore the coupling of the would-be NG boson to the gauge field must render the latter massive! This is the essence of the BEH mechanism.

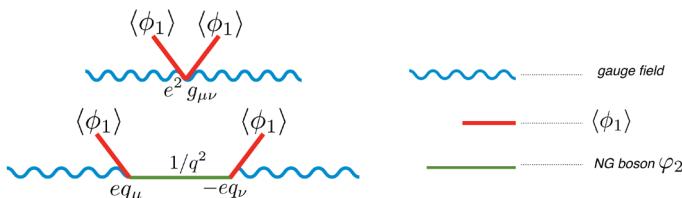


FIGURE 7. Interaction of the gauge field with the condensate.

These qualitative considerations can be made quantitative [1] by considering the Feynman graphs (time runs horizontally) describing the propagation of the A_μ gauge field in the vacuum with non vanishing scalar field expectation value, say $\langle \phi_1 \rangle \neq 0$. This propagation is depicted in lowest order in Figure 7 (time runs horizontally) and the interaction of A_μ with the condensate $\langle \phi_1 \rangle$ amounts to a “polarisation” of the vacuum. The first graph shows the local interaction of the gauge field with the condensate while the second one gives a non-local interaction due to the propagation of a NG boson. Here e is the coupling of the gauge vector to matter, q_μ is a four-momentum ($q_0 = \text{energy}$; $\vec{q} = \text{momentum}$), $q^2 = q_0^2 - \vec{q}^2$ and $g_{\mu\nu}$ has only non-zero values if $\mu \neq \nu : 1, -1, -1, -1$. The two graphs add up to

$$\Pi_{\mu\nu} = (g_{\mu\nu} - q_\mu q_\nu / q^2) \Pi(q^2), \quad (3.2)$$

where

$$\Pi(q^2) = e^2 \langle \phi_1 \rangle^2. \quad (3.3)$$

The second factor of Equation (3.2) does not vanish when $q^2 = 0$. In field theory this means that the gauge field has acquired a mass

$$(M_V^2) = e^2 \langle \phi_1 \rangle^2. \quad (3.4)$$

The first factor describes the projection at $q^2 = m_V^2$ of $g_{\mu\nu}$ on a three-dimensional space of polarisations, which, as explained in qualitative terms above, is required for a massive vector. Its transversality (i.e. its vanishing under multiplication by q^μ) is characteristic of a “Ward Identity” which expresses the fact that the local gauge symmetry has not been broken and is identical to the analogous factor in quantum electrodynamics, a important fact that will be commented on in the following section.

The generalisation of these results to more complicated symmetries yields (for real fields) a mass matrix

$$(M_V^2)^{ab} = -e^2 \langle \phi^B \rangle T^{aBC} T^{bCA} \langle \phi^A \rangle, \quad (3.5)$$

where T^{aBC} is a real anti-symmetric generator coupled to a gauge field A_μ^a and $\langle \phi^A \rangle$ designates a non vanishing expectation value.

In these cases, some gauge fields may remain massless. Consider for instance instead of the invariance of the Goldstone model on a circle in the plane (ϕ_1, ϕ_2) , an invariance on a sphere in a 3-dimensional space (ϕ_1, ϕ_2, ϕ_3) broken by $\langle \phi_1 \rangle \neq$

0. There are now three gauge fields associated to the rotations on the sphere, and while A_μ^2 and A_μ^3 acquire mass A_μ^1 remains massless. This can be understood in the following way: rotation generators around the directions 2 and 3 would move $\langle\phi_1\rangle$ if the symmetry were global and would thus give rise, as in Figure 5, to NG bosons; their degrees of freedom are transferred in local symmetries to the massive gauge vector fields A_μ^2 and A_μ^3 , providing their third degree of polarisation. The expectation value $\langle\phi_1\rangle$ is not affected by rotation generators around the direction 1 and does not generate NG bosons in the global symmetry case and hence the corresponding A_μ^1 remains massless.

Thus the BEH mechanism can unify long and short range interactions in the same theory by leaving unbroken a subgroup of symmetry transformation (e.g. rotation around the direction 1) whose corresponding gauge fields remain massless.

3.2 The fate of the massive scalar boson

A glance at Figure 5 shows that the stretching of (classical) scalar fields is independent of local rotations of the ϕ -field in the (ϕ_1, ϕ_2) plane. This translates the fact that the modulus of the ϕ -field is gauge invariant. Hence the scalar bosons survive the gauging and their classical analysis is identical to the one given for the Goldstone model in Section 1.2. The coupling of the scalar boson φ_1 to the massive gauge bosons follows from the Figure 7, by viewing the Feynman diagrams with time going from top to bottom and using Equation (3.1). One gets the two vertices of Figure 8 where the heavy wiggly lines on the right hand side represent the massive gauge propagators. The vertex couplings follow from Equation (3.4).

3.3 Fermion masses

Let us couple the Yang-Mills fields to massless fermions in a way that respects Yang-Mills symmetry. This coupling preserves the chiral symmetry of the

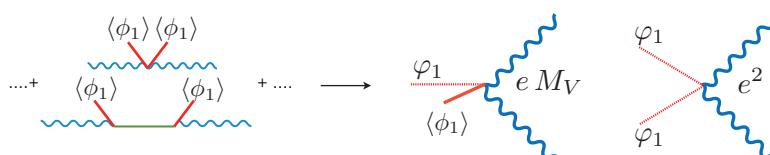


FIGURE 8. Coupling of the scalar boson φ_1 to massive gauge bosons.

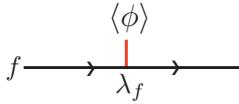


FIGURE 9. Mass generation $m_f = \lambda_f \langle \phi \rangle$ from a coupling λ_f of fermions to the scalar field ϕ .

massless fermions and fermion mass requires SSB. In the Nambu theory of spontaneous breaking of chiral symmetry, this gives rise to NG bosons which are eaten up here by massive gauge fields. This can be done by suitable couplings of the scalar fields whose expectation value breaks the symmetry. Mass generation for fermions is depicted in Figure 9.

3.4 Why is the mechanism needed?

Equation (3.2) expresses the fact that the mass generation from the BEH mechanism does not destroy local symmetry, in contradistinction to a mass term introduced by hand *ab initio*. This equation remains valid at higher orders in perturbation theory and has the same form as the polarisation in quantum electrodynamics. As in the latter case, it implies that in covariant gauges, the gauge vector boson propagator tames the quantum fluctuations, and therefore suggests that the theory is renormalisable [15]. However it is a highly non trivial matter to prove that it does not introduce contributions from unphysical particles, and it is therefore a very difficult problem to prove quantum consistency to all orders. That this is indeed the case has been proven by 't Hooft and Veltman [16] (see also Ref. [17])

The quantum consistency of the BEH mechanism is the basic reason for its success. Precision experiments can be predicted and were indeed verified. Quantum consistency played a critical role in the analysis of the production of the scalar boson at the LHC and of its decay products, leading to the confirmation of the detailed validity of the mechanism.

3.5 Dynamical symmetry breaking

The symmetry breaking giving mass to gauge vector bosons may also arise from a fermion condensate. This is labelled *dynamical symmetry breaking*. If a spontaneously broken global symmetry is extended to a local one by introducing gauge fields, the massless NG bosons disappear as previously from the physical spectrum and their absorption by gauge fields renders these massive. In contradistinction with breaking by scalar field condensate, it is very difficult in this

way to give mass in a renormalisable theory simultaneously to both gauge vector fields and fermions.

3.6 The electroweak theory and the Standard Model

The most impressive success of the BEH mechanism is the electroweak theory for weak and electromagnetic interactions [18] applied to all particles of the Standard Model. These encompass all known particles. These are a) the fermions which are listed in Figure 10, b) γ and W^+, W^-, Z , the gauge vector bosons transmitting the electromagnetic and the weak interactions, c) eight “gluons,” the gauge vectors bosons of the “colour group” $SU(3)$ mediating the strong interactions, and d) last but not least, one massive scalar boson which was recently discovered and identified as the scalar predicted by the BEH mechanism.

The first row in Figure 10 contains the basic constituents of the atom, namely the electron, the three up and down coloured quarks building the proton and neutron bonded by the gluons, to which is added the electron neutrino. The second [19] and third row [20] were completed as predictions in the 1970s and verified afterwards. Colour was also introduced in the 60s. The particles in the first and the second row are called leptons. To all fermions of the table, one must of course also add their antiparticles.

All the fermions are chiral and their chiral components have different group quantum numbers. Hence they are, as the gauge vector bosons, massless in absence of the BEH mechanism, i.e. in absence of the scalar condensate. The condensate $\langle \phi \rangle \neq 0$ gives mass to the W^+, W^-, Z bosons and to all fermions except

particles (charge)				
$e (-1)$	$\nu_e (0)$	$u u u (\frac{2}{3})$	$d d d (-\frac{1}{3})$	
$\mu (-1)$	$\nu_\mu (0)$	$c c c (\frac{2}{3})$	$s s s (-\frac{1}{3})$	+ [antiparticles]
$\tau (-1)$	$\nu_\tau (0)$	$t t t (\frac{2}{3})$	$b b b (-\frac{1}{3})$	

FIGURE 10. Fermion constituents of the Standard Model.

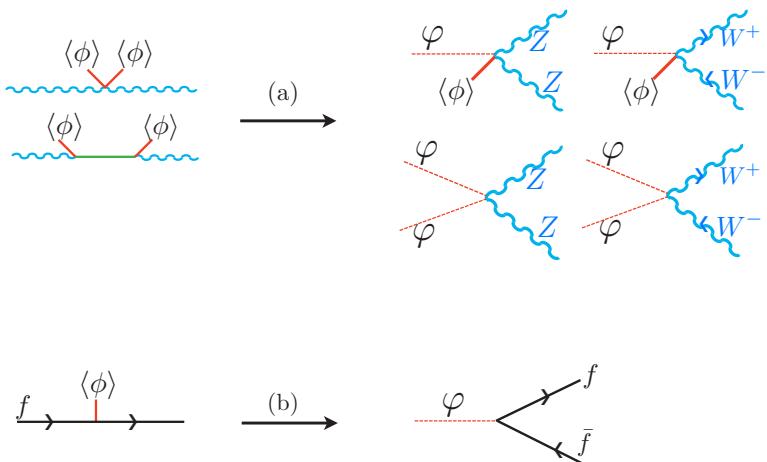


FIGURE 11. Coupling of the scalar boson φ to massive gauge bosons and to elementary fermions.

the three chiral neutrinos which have no opposite chirality counterpart in the conventional Standard Model. The photons and the gluons remain massless but the latter become short range due in the conventional description to a highly non-perturbative vacuum (resulting from a mechanism somehow dual to the BEH mechanism).

The discovery of the Z and W bosons in 1983 and the precision experiments testing the quantum consistency of the Standard Model established the validity of the mechanism, but it was still unclear whether this was the result of a dynamical symmetry breaking or of a particle identifiable as an elementary boson at the energy scale considered.

4 THE DISCOVERY

In the Standard Model, there is one real massive scalar boson φ (also labelled H). It couples to the massive W and Z bosons. This follows from Figure 8 and the couplings are depicted in Figure 11 (a). Its coupling to elementary fermions similarly follows from the couplings in Figure 9 as shown in Figure 11 (b). The coupling to the massless photons is a genuine quantum effect involving loops, even in the lowest order, as indicated in Figure 12.

The LHC site circling under the French-Swiss border is schematically indicated in the picture of Figure 13. The 27 km circular tunnel containing two

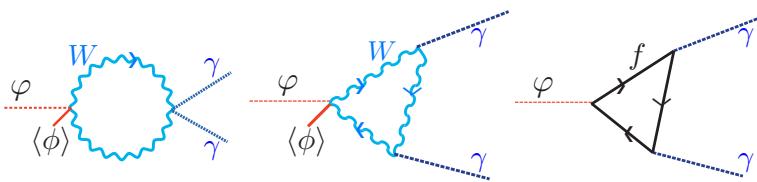


FIGURE 12. Coupling of the scalar boson φ to photons.

opposite beams of protons surrounded by guiding superconducting electromagnets cooled by superfluid helium is pictured in Figure 14. Figures 15 and 16 are pictures of the ATLAS and CMS detectors at diametrical opposite sites of the tunnel. There collisions occur and were used primarily to detect and identify the scalar boson of the Standard Model (and possibly other ones). At the end of 2012 proton-proton collisions occurred at the rate of nearly 10^9s^{-1} and the proton energy reached 8 TeV. At these energies, all quarks of Figure 10 and the gluons connecting them may contribute to the production of the scalar boson. The leading production processes are represented in Figure 17.

As an example of the data gathered by CMS and ATLAS, Figure 18 presents the data obtained by the CMS group of observed decays into 4 leptons at the end



FIGURE 13. Schematic location of the LHC.



FIGURE 14. The LHC dipole magnets.

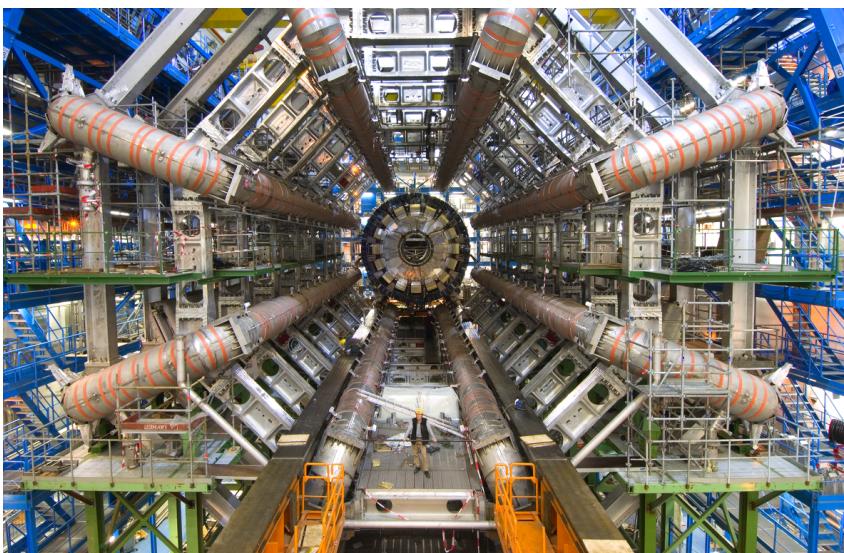


FIGURE 15. The Atlas detector.

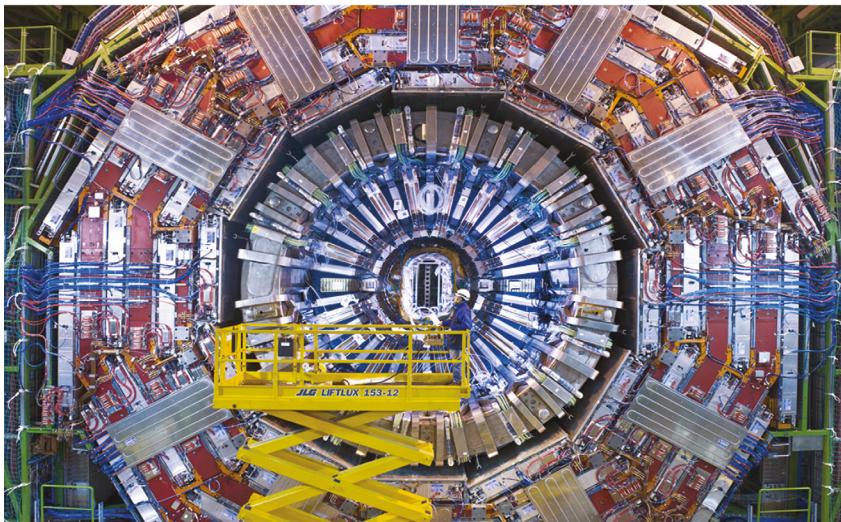


FIGURE 16. The CMS detector.

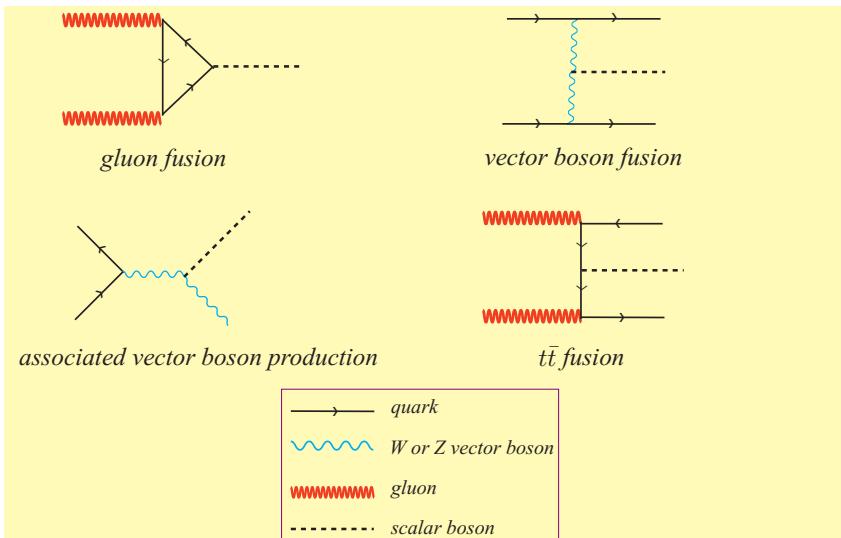


FIGURE 17. Production of the Standard Model scalar.

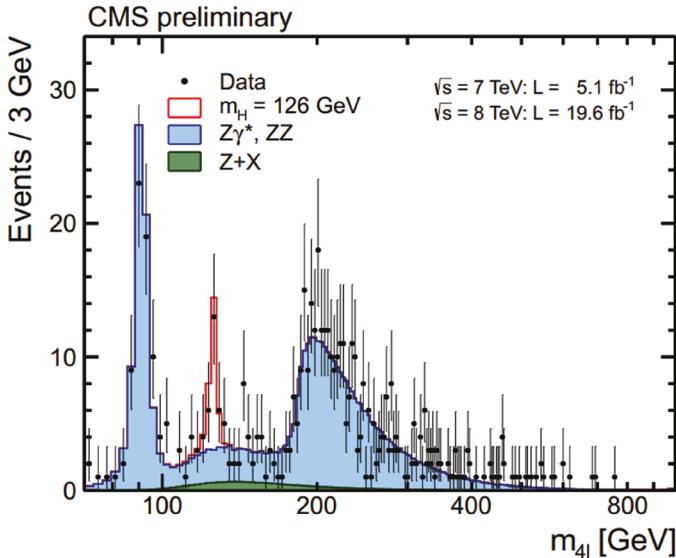


FIGURE 18. Decay of the scalar boson into 4 leptons from two Z s.

the 2012 run. The blue area is the expected background, namely those decays which would follow from the Standard Model if, at given total mass, there would be no contribution from the scalar boson. The red curve measures the contribution that could be due to the scalar decaying into two Z vector bosons which further decay into leptons, as was confirmed by further analysis. Note that one of the Z is real but the other is “virtual,” meaning that this decay is forbidden by energy conservation but may contribute in the quantum theory. Consideration of other decay channels and spin analysis show that the particle detected is consistent with the Standard Model scalar with a mass $m_H \simeq 125$ GeV. The absence of new particles at comparable energies, as well as the success of Feynman graph analysis including loops, points towards an elementary particle, at least up to the energy range considered. This is the first elementary spin zero particle ever detected. It raises the interesting possibility of supersymmetry broken at attainable energies, although there is no indication of it so far.

The elementary character of the scalar already eliminates many dynamical models of symmetry breaking and raises interesting possibilities for extrapolation beyond presently known energies, up to those close to the Planck scale where quantum gravity effects might play a dominant role. The analysis of these speculations is beyond the scope of this talk.

REFERENCES

1. F. Englert and R. Brout, “Broken symmetry and the mass of gauge vector mesons,” *Phys. Rev. Letters* **13** (1964) 321.
2. P.W. Higgs, “Broken symmetries and the masses of gauge bosons,” *Phys. Rev. Lett.* **13** (1964) 508.
3. L.D. Landau, “On the theory of phase transitions I,” *Phys. Z. Sowjet.* **11** (1937) 26 [JETP 7 (1937) 19].
4. Y. Nambu, “Quasi-particles and gauge invariance in the theory of superconductivity,” *Phys. Rev.* **117** (1960) 648.
5. Y. Nambu, “Axial vector current conservation in weak interactions,” *Phys. Rev. Lett.* **4** (1960) 380.
6. Y. Nambu and G. Jona-Lasinio, “Dynamical model of elementary particles based on an analogy with superconductivity I, II,” *Phys. Rev.* **122** (1961) 345; **124** (1961) 246.
7. P.W. Anderson, “Random-phase approximation in the theory of superconductivity,” *Phys. Rev.* **112** (1958) 1900.
8. P.W. Anderson, “Plasmons, gauge invariance, and mass,” *Phys. Rev.* **130** (1963) 439.
9. J. Goldstone, “Field theories with ‘superconductor’ solutions,” *Il Nuovo Cimento* **19** (1961) 154.
10. J. Goldstone, A. Salam and S. Weinberg, “Broken symmetries,” *Phys. Rev.* **127** (1962) 965.
11. F. Englert, “Broken symmetry and Yang-Mills theory,” in *50 years of Yang-Mills Theory*, ed. by G.’t Hooft, World Scientific (2005), pp. 65–95, hep-th/0406162.
12. S. Elitzur, “Impossibility of spontaneously breaking local symmetries,” *Phys. Rev.* **D12** (1975) 3978.
13. P.W. Higgs, “Broken symmetries, massless particles and gauge fields,” *Phys. Lett.* **12** (1964) 132.
14. G.S. Guralnik, C.R. Hagen and T.W.B. Kibble, “Global conservation laws and massless particles,” *Phys. Rev. Lett.* **13** (1964) 585.
15. F. Englert, R. Brout and M. Thiry, “Vector mesons in presence of broken symmetry,” *Il Nuovo Cimento* **43A** (1966) 244; F. Englert, *Proceedings of the 1997 Solvay Conference, Fundamental Problems in Elementary Particle Physics*, Interscience Publishers J. Wiley and Sons, p. 18.
16. G. ’t Hooft, “Renormalizable lagrangians for massive Yang-Mills fields,” *Nucl. Phys.* **B35** (1971) 167; G. ’t Hooft and M. Veltman, Regularization and renormalization of gauge fields, *Nucl. Phys.* **B44** (1972) 189.
17. B.W. Lee and J. Zinn-Justin, “Spontaneously broken gauge symmetries,” *Phys. Rev.* **D5** (1972) 3121; 3137; 3155.
18. S.L. Glashow, “Partial-symmetries of weak interactions,” *Nucl. Phys.* **22** (1961) 579; S. Weinberg, A model of leptons, *Phys. Rev. Lett.* **19** (1967) 1264; A. Salam,

Proceedings of the 8th Nobel Symposium, Elementary Particle Physics, ed. by N. Svartholm, (Almqvist and Wiksell, Stockholm) p. 367.

19. S.L. Glashow, J. Iliopoulos and L. Maiani, “Weak interactions with lepton-hadron symmetry,” *Phys. Rev.* **D2** (1970) 1285.
20. M. Kobayashi and T. Maskawa, “CP Violation in the renormalizable theory of weak interaction,” *Prog. Theor. Phys.* **49** (1973) 652.

Portrait photo of François Englert by photographer Alexander Mahmoud.