COLLOQUIUM ON THE 2013 NOBEL PRIZE IN PHYSICS AWARDED TO FRANCOIS ENGLERT AND PETER HIGGS

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Presentation at University of Connecticut

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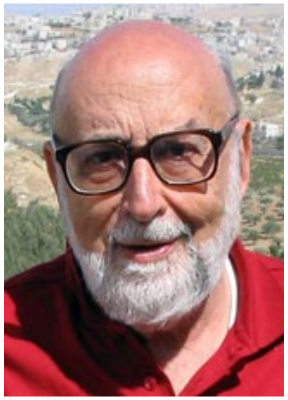
The Nobel Prize in Physics 2013
François Englert, Peter Higgs

The Nobel Prize in Physics 2013



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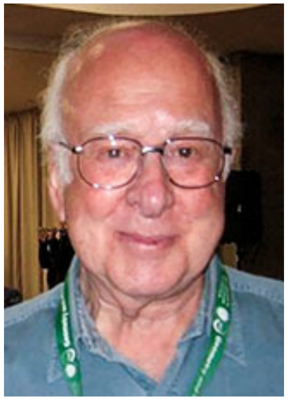
François Englert



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Peter W. Higgs

The Nobel Prize in Physics 2013 was awarded jointly to François Englert and Peter W. Higgs ***"for the theoretical discovery of a mechanism that contributes to our understanding of the origin of mass of subatomic particles, and which recently was confirmed through the discovery of the predicted fundamental particle, by the ATLAS and CMS experiments at CERN's Large Hadron Collider"***

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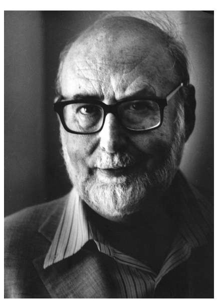
Brout-Englert-Higgs mechanism



BROKEN SYMMETRY AND THE MASS OF GAUGE VECTOR MESONS*

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(Received 26 June 1964)

See here (EN) or here (FR) for information (news, presentations, pics,...) on the Brout-Englert-Higgs mechanism and the Standard Model Scalar particle.

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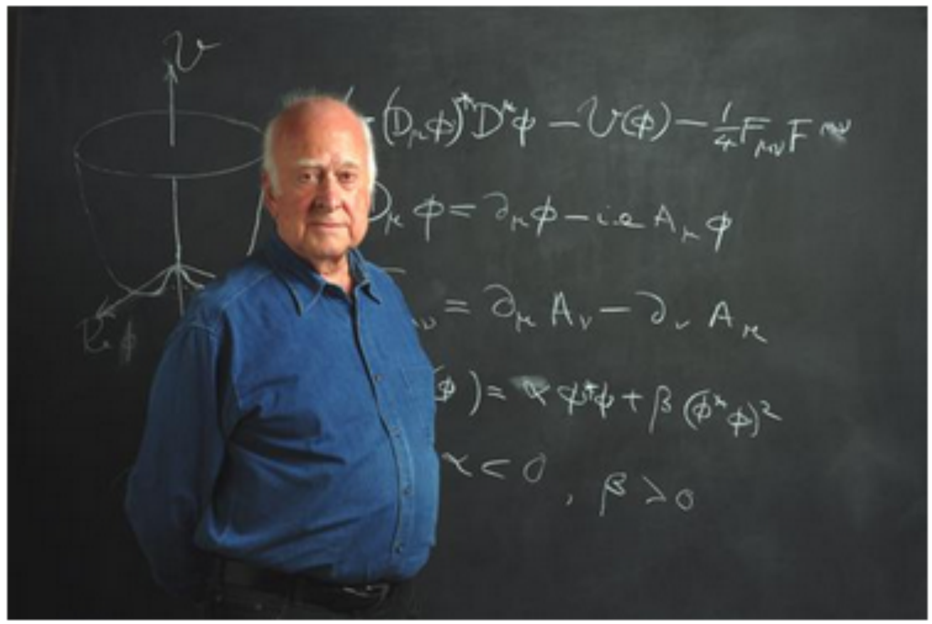

Welcome to the Higgs site at the University of Edinburgh

8th October 2013: The Nobel Prize in Physics 2013 was awarded jointly to François Englert and Peter W. Higgs *"for the theoretical discovery of a mechanism that contributes to our understanding of the origin of mass of subatomic particles, and which recently was confirmed through the discovery of the predicted fundamental particle, by the ATLAS and CMS experiments at CERN's Large Hadron Collider"* Nobel Prize announcement

1 Introduction

1.1 Background

The 2013 Nobel Prize in Physics was awarded to Francois Englert and Peter Higgs for their work in 1964 that lead in 2012 to the discovery at the CERN Large Hadron Collider (LHC) of the Higgs Boson after a 50 year search. It is great tragedy that Robert Brout, the joint author with Francois Englert of one of the papers that led to the 2013 Nobel Prize, died in 2011, just one year before the discovery of the Higgs boson and two years before the awarding of the Nobel prize for it. (For me personally this is keenly felt since my first post-doc was with Robert and Francois in Brussels 1970 - 1972.)

The significance of the Higgs Boson is that it is tied in with the theory of the origin of mass and of the way that mass can arise through collective effects (known as broken symmetry) that only occur in systems with a large number of degrees of freedom. Such collective effects are properties that a system of many objects collectively possess that each one individually does not – the whole being greater than the sum of its parts. A typical example is temperature. A single molecule of H_2O does not have a temperature, and you cannot tell if it was taken from ice, water or steam. These different phases are collective properties of large numbers of H_2O molecules acting in unison. Moreover, as one changes the temperature one can change the phase (freezing water into ice for instance), with it being the existence of such phase changes that is central to broken symmetry.

I counted at least 18 times that Nobel Prizes in Physics have in one way or another been given for aspects of this problem:

Dirac (1933); Anderson (1936); Lamb (1955); Landau (1962); Tomonaga, Schwinger, Feynman (1965); Gell-Mann (1969); Bardeen, Cooper, Schrieffer (1972); Richter, Ting (1976); Glashow, Salam, Weinberg (1979); Wilson (1982); Rubbia, van der Meer (1984); Friedman, Kendall, Taylor (1990); Lee, Osheroff, Richardson (1996); 't Hooft,Veltman (1999); Abrikosov, Ginzburg, Leggett (2003); Gross, Politzer, Wilczek (2004); Nambu, Kobayashi, Maskawa (2008); Englert, Higgs (2013).

And this leaves out Philip Anderson who made major contributions to collective aspects of mass generation and C.N. Yang who with Mills developed non-Abelian Yang-Mills gauge theories but got Nobel prizes (1977, 1957) for something else.

So now to the paper by Englert and Brout and the two papers by Higgs. The paper by Englert and Brout appeared in Physical Review Letters 13, 321 (1964) after being submitted on June 26, 1964 and was two and one half pages long. The first of Higgs' two papers appeared in Physics Letters 12, 132 (1964) after being submitted on July 27, 1964 and was one and one half pages long, and the second paper appeared in Physical Review Letters 13, 508 (1964) after being submitted on August 31, 1964 and was also one and one half pagers long. Thus a grand total of just five and one half pages.

*Work supported in part by the U. S. Atomic Energy Commission and in part by the Graduate School from funds supplied by the Wisconsin Alumni Research Foundation.

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BROKEN SYMMETRY AND THE MASS OF GAUGE VECTOR MESONS*

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(Received 26 June 1964)

It is of interest to inquire whether gauge vector mesons acquire mass through interaction[1]; by a gauge vector meson we mean a Yang-Mills field[2] associated with the extension of a Lie group from global to local symmetry. The importance of this problem resides in the possibility that strong-interaction physics originates from massive gauge fields related to a system of conserved currents.[3] In this note, we shall show that in certain cases vector mesons do indeed acquire mass when the vacuum is degenerate with respect to a compact Lie group.

Theories with degenerate vacuum (broken symmetry) have been the subject of intensive study since their inception by Nambu.[4-6] A characteristic feature of such theories is the possible existence of zero-mass bosons which tend to restore the symmetry.[7,8] We shall show that it is precisely these singularities which maintain the gauge invariance of the theory, despite the fact that the vector meson acquires mass.

We shall first treat the case where the original fields are a set of bosons φ_A which transform as a basis for a representation of a compact Lie group. This example should be considered as a rather general phenomenological model. As such, we shall not study the particular mechanism by which the symmetry is broken but simply assume that such a mechanism exists. A calculation performed in lowest order perturbation theory indicates that those vector mesons which are coupled to currents that "rotate" the original vacuum are the ones which acquire mass [see Eq. (6)].

We shall then examine a particular model based on chirality invariance which may have a more fundamental significance. Here we begin with a chirality-invariant Lagrangian and introduce both vector and pseudovector gauge fields, thereby guaranteeing invariance under both local phase and local γ_5-phase transformations. In this model the gauge fields themselves may break the γ_5 invariance leading to a mass for the original Fermi field. We shall show in this case that the pseudovector field acquires mass.

In the last paragraph we sketch a simple argument which renders these results reasonable.

(1) Lest the simplicity of the argument be shrouded in a cloud of indices, we first consider a one-parameter Abelian group, representing, for example, the phase transformation of a charged boson; we then present the generalization to an arbitrary compact Lie group.

The interaction between the φ and the A_μ fields is

$$H_{\text{int}} = ieA_\mu \varphi^* \overleftrightarrow{\partial}_\mu \varphi - e^2 \varphi^* \varphi A_\mu A_\mu, \qquad (1)$$

where $\varphi = (\varphi_1 + i\varphi_2)/\sqrt{2}$. We shall break the symmetry by fixing $\langle\varphi\rangle \neq 0$ in the vacuum, with the phase chosen for convenience such that $\langle\varphi\rangle = \langle\varphi^*\rangle = \langle\varphi_1\rangle/\sqrt{2}$.

We shall assume that the application of the theorem of Goldstone, Salam, and Weinberg[7] is straightforward and thus that the propagator of the field φ_2, which is "orthogonal" to φ_1, has a pole at $q=0$ which is not isolated.

We calculate the vacuum polarization loop $\Pi_{\mu\nu}$ for the field A_μ in lowest order perturbation theory about the self-consistent vacuum. We take into consideration only the broken-symmetry diagrams (Fig. 1). The conventional terms do not lead to a mass in this approximation if gauge invariance is carefully maintained. One evaluates directly

$$\Pi_{\mu\nu}(q) = (2\pi)^4 ie^2[g_{\mu\nu}\langle\varphi_1\rangle^2 - (q_\mu q_\nu/q^2)\langle\varphi_1\rangle^2]. \quad (2)$$

Here we have used for the propagator of φ_2 the value $[i/(2\pi)^4]/q^2$; the fact that the renormalization constant is 1 is consistent with our approximation.[9] We then note that Eq. (2) both maintains gauge invariance ($\Pi_{\mu\nu}q_\nu = 0$) and causes the A_μ field to acquire a mass

$$\mu^2 = e^2\langle\varphi_1\rangle^2. \quad (3)$$

We have not yet constructed a proof in arbitrary order; however, the similar appearance of higher order graphs leads one to surmise the general truth of the theorem.

Consider now, in general, a set of boson-field operators φ_A (which we may always choose to be Hermitian) and the associated Yang-Mills field $A_{a,\mu}$. The Lagrangian is invariant under the transformation[10]

$$\delta\varphi_A = \sum\nolimits_{a,A}\epsilon_a(x)T_{a,AB}\varphi_B,$$
$$\delta A_{a,\mu} = \sum\nolimits_{c,b}\epsilon_c(x)c_{acb}A_{b,\mu} + \partial_\mu\epsilon_a(x), \quad (4)$$

where c_{abc} are the structure constants of a compact Lie group and $T_{a,AB}$ the antisymmetric generators of the group in the representation defined by the φ_B.

Suppose that in the vacuum $\langle\varphi_{B'}\rangle \neq 0$ for some B'. Then the propagator of $\sum_{A,B'}T_{a,AB'}\varphi_A$ $\times\langle\varphi_{B'}\rangle$ is, in the lowest order,

$$\left[\frac{i}{(2\pi)^4}\right]\sum_{A,B',C'}\frac{T_{a,AB'}\langle\varphi_{B'}\rangle T_{a,AC'}\langle\varphi_{C'}\rangle}{q^2}$$
$$\equiv\left[\frac{-i}{(2\pi)^4}\right]\frac{(\langle\varphi\rangle T_aT_a\langle\varphi\rangle)}{q^2}.$$

With λ the coupling constant of the Yang-Mills field, the same calculation as before yields

$$\Pi_{\mu\nu}{}^a(q) = -i(2\pi)^4\lambda^2(\langle\varphi\rangle T_aT_a\langle\varphi\rangle)$$
$$\times[g_{\mu\nu} - q_\mu q_\nu/q^2],$$

giving a value for the mass

$$\mu_a{}^2 = -(\langle\varphi\rangle T_aT_a\langle\varphi\rangle). \quad (6)$$

(2) Consider the interaction Hamiltonian

$$H_{\text{int}} = -\eta\bar\psi\gamma_\mu\gamma_5\psi B_\mu - \epsilon\bar\psi\gamma_\mu\psi A_\mu, \quad (7)$$

where A_μ and B_μ are vector and pseudovector gauge fields. The vector field causes attraction whereas the pseudovector leads to repulsion between particle and antiparticle. For a suitable choice of ϵ and η there exists, as in Johnson's model,[11] a broken-symmetry solution corresponding to an arbitrary mass m for the ψ field fixing the scale of the problem. Thus the fermion propagator $S(p)$ is

$$S^{-1}(p) = \gamma p - \Sigma(p) = \gamma p[1-\Sigma_2(p^2)] - \Sigma_1(p^2), \quad (8)$$

with

$$\Sigma_1(p^2) \neq 0$$

and

$$m[1-\Sigma_2(m^2)] - \Sigma_1(m^2) = 0.$$

We define the gauage-invariant current $J_\mu{}^5$ by using Johnson's method[12]:

$$J_\mu{}^5 = -\eta\lim_{\xi\to 0}\bar\psi'(x+\xi)\gamma_\mu\gamma_5\psi'(x),$$

$$\psi'(x) = \exp[-i\textstyle\int_{-\infty}^x \eta B_\mu(y)dy^\mu\gamma_5]\psi(x). \quad (9)$$

This gives for the polarization tensor of the

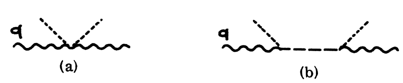

FIG. 1. Broken-symmetry diagram leading to a mass for the gauge field. Short-dashed line, $\langle\varphi_1\rangle$; long-dashed line, φ_2 propagator; wavy line, A_μ propagator. (a) $\rightarrow (2\pi)^4 ie^2 g_{\mu\nu}\langle\varphi_1\rangle^2$, (b) $\rightarrow -(2\pi)^4 ie^2(q_\mu q_\nu/q^2)$ $\times\langle\varphi_1\rangle^2$.

pseudovector field

$$\Pi_{\mu\nu}{}^5(q) = \eta^2 \frac{i}{(2\pi)^4} \int \mathrm{Tr}\{S(p-\tfrac{1}{2}q)\Gamma_{\nu 5}(p-\tfrac{1}{2}q; p+\tfrac{1}{2}q) \times S(p+\tfrac{1}{2}q)\gamma_\mu\gamma_5 - S(p)[\partial S^{-1}(p)/\partial p_\nu]S(p)\gamma_\mu\} d^4p, \quad (10)$$

where the vertex function $\Gamma_{\nu 5} = \gamma_\nu\gamma_5 + \Lambda_{\nu 5}$ satisfies the Ward identity[5]

$$q_\nu \Lambda_{\nu 5}(p-\tfrac{1}{2}q; p+\tfrac{1}{2}q) = \Sigma(p-\tfrac{1}{2}q)\gamma_5 + \gamma_5\Sigma(p+\tfrac{1}{2}q), \quad (11)$$

which for low q reads

$$q_\nu\Gamma_{\nu 5} = q_\nu\gamma_\nu\gamma_5[1-\Sigma_2] + 2\Sigma_1\gamma_5 - 2(q_\nu p_\nu)(\gamma_\lambda p_\lambda)(\partial\Sigma_2/\partial p^2)\gamma_5. \quad (12)$$

The singularity in the longitudinal $\Gamma_{\nu 5}$ vertex due to the broken-symmetry term $2\Sigma_1\gamma_5$ in the Ward identity leads to a nonvanishing gauge-invariant $\Pi_{\mu\nu}{}^5(q)$ in the limit $q \to 0$, while the usual spurious "photon mass" drops because of the second term in (10). The mass of the pseudovector field is roughly $\eta^2 m^2$ as can be checked by inserting into (10) the lowest approximation for $\Gamma_{\nu 5}$ consistant with the Ward identity.

Thus, in this case the general feature of the phenomenological boson system survives. We would like to emphasize that here the symmetry is broken through the gauge fields themselves. One might hope that such a feature is quite general and is possibly instrumental in the realization of Sakurai's program.[3]

(3) We present below a simple argument which indicates why the gauge vector field need not have zero mass in the presence of broken symmetry. Let us recall that these fields were introduced in the first place in order to extend the symmetry group to transformations which were different at various space-time points. Thus one expects that when the group transformations become homogeneous in space-time, that is $q \to 0$, no dynamical manifestation of these fields should appear. This means that it should cost no energy to create a Yang-Mills quantum at $q = 0$ and thus the mass is zero. However, if we break gauge invariance of the first kind and still maintain gauge invariance of the second kind this reasoning is obviously incorrect. Indeed, in Fig. 1, one sees that the A_μ propagator connects to intermediate states, which are "rotated" vacua. This is seen most clearly by writing $\langle\varphi_1\rangle = \langle[Q\varphi_2]\rangle$ where Q is the group generator. This effect cannot vanish in the limit $q \to 0$.

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*This work has been supported in part by the U. S. Air Force under grant No. AFEOAR 63-51 and monitored by the European Office of Aerospace Research.

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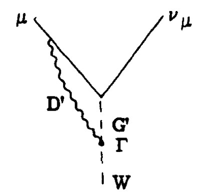

Fig. 1.

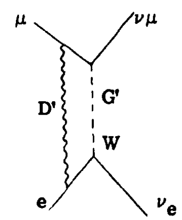

Fig. 2.

well into account the radiation correction to the β-decay constant found by Berman [3] and Kinoshita and Sirlin [4] we obtain for the muon life time

$$\frac{\tau_\mu}{\tau_\mu{}^0} = 1 - \frac{3e^2}{4\pi}\ln\frac{\Lambda^2}{\mu^2} + \frac{3e^2}{2\pi}\ln\frac{\Lambda_\beta}{2E} - \frac{3}{5}\frac{M_\mu{}^2}{\mu^2}\,, \qquad (1)$$

where $\tau_\mu{}^0$ is the muon life time calculated by means of universal theory of four fermion interaction with a constant taken from β-decay without any corrections, Λ_β is the cut off momentum due to the strong interactions, $\Lambda_\beta \sim M$, E is the energy of β-transition. According to experimental data $\tau_\mu/\tau_\mu{}^0 = 0.988 \pm 0.004$.

Substituting the numbers into (1) we obtain $\tau_\mu/\tau_\mu{}^0 = 1.003$ and the disagreement between the theory and experiment will be in our case $1.5 \pm 0.4\%$. When discussing this result one should take into consideration that in (1) only the terms $\sim e^2 \ln e^{-2}$ were correctly taken into account but the terms $\sim e^2$ were discarded.

It seems to us that the conclusion that in the theory of weak interaction with intermediate W-meson β- and μ-constants must be with good accuracy the same (taking into account the corrections due to the electromagnetic and weak interactions), is in favour of the weak interaction theory with W-meson unlike the four-fermion theory.

More detailed paper will be published elsewhere.

The author is indebted to B. V. Geshkenbein, I. Yu. Kobsarev, L. B. Okun, A. M. Perelomov, I. Ya. Pomeranchuk, V. S. Popov, A. P. Rudik and M. V. Terentyev for valuable discussions.

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BROKEN SYMMETRIES, MASSLESS PARTICLES AND GAUGE FIELDS

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Recently a number of people have discussed the Goldstone theorem [1,2]: that any solution of a Lorentz-invariant theory which violates an internal symmetry operation of that theory must contain a massless scalar particle. Klein and Lee [3] showed that this theorem does not necessarily apply in non-relativistic theories and implied that their considerations would apply equally well to Lorentz-invariant field theories. Gilbert [4], however, gave a proof that the failure of the Goldstone theorem in the nonrelativistic case is of a type which cannot exist when Lorentz invariance is imposed on a theory. The purpose of this note is to show that Gilbert's argument fails for an important class of field theories, that in which the conserved currents are coupled to gauge fields.

Following the procedure used by Gilbert [4], let us consider a theory of two hermitian scalar fields

$\varphi_1(x)$, $\varphi_2(x)$ which is invariant under the phase transformation

$$\begin{aligned} \varphi_1 &\to \varphi_1 \cos\alpha + \varphi_2 \sin\alpha \,, \\ \varphi_2 &\to -\varphi_1 \sin\alpha + \varphi_2 \cos\alpha \,. \end{aligned} \tag{1}$$

Then there is a conserved current j_μ such that

$$\mathrm{i}[\int \mathrm{d}^3x\, j_0(x),\ \varphi_1(y)] = \varphi_2(y). \tag{2}$$

We assume that the Lagrangian is such that symmetry is broken by the nonvanishing of the vacuum expectation value of φ_2. Goldstone's theorem is proved by showing that the Fourier transform of $\mathrm{i}\langle[j_\mu(x),\ \varphi_1(y)]\rangle$ contains a term $2\pi\langle\varphi_2\rangle\epsilon(k_0)k_\mu\delta(k^2)$, where k_μ is the momentum, as a consequence of Lorentz-covariance, the conservation law and eq. (2).

Klein and Lee [3] avoided this result in the nonrelativistic case by showing that the most general form of this Fourier transform is now, in Gilbert's notation,

$$\mathrm{F.T.} = k_\mu \rho_1(k^2, nk) + n_\mu \rho_2(k^2, nk) + C_3 n_\mu \delta^4(k)\,, \tag{3}$$

where n_μ, which may be taken as (1, 0, 0, 0), picks out a special Lorentz frame. The conversation law then reduces eq. (3) to the less general form

$$\mathrm{F.T.} = k_\mu \delta(k^2)\rho_4(nk) + [k^2 n_\mu - k_\mu(nk)]\rho_5(k^2, nk) + C_3 n_\mu \delta^4(k)\,. \tag{4}$$

It turns out, on applying eq. (2), that all three terms in eq. (4) can contribute to $\langle\varphi_2\rangle$. Thus the Goldstone theorem fails if $\rho_4 = 0$, which is possible only if the other terms exist. Gilbert's remark that no special timelike vector n_μ is available in a Lorentz-covariant theory appears to rule out this possibility in such a theory.

There is however a class of relativistic field theories in which a vector n_μ does indeed play a part. This is the class of gauge theories, where an auxiliary unit timelike vector n_μ must be introduced in order to define a radiation gauge in which the vector gauge fields are well defined operators. Such theories are nevertheless Lorentz-covariant, as has been shown by Schwinger [5]. (This has, of course, long been known of the simplest such theory, quantum electrodynamics.) There seems to be no reason why the vector n_μ should not appear in the Fourier transform under consideration.

It is characteristic of gauge theories that the conservation laws hold in the strong sense, as a consequence of field equations of the form

$$\begin{aligned} j^\mu &= \partial_\nu F'^{\mu\nu}, \\ F_{\mu\nu}' &= \partial_\mu A_\nu' - \partial_\nu A_\mu'\,. \end{aligned} \tag{5}$$

Except in the case of abelian gauge theories, the fields A_μ', $F_{\mu\nu}'$ are not simply the gauge field variables A_μ, $F_{\mu\nu}$, but contain additional terms with combinations of the structure constants of the group as coefficients. Now the structure of the Fourier transform of $\mathrm{i}\langle[A_\mu'(x),\ \varphi_1(y)]\rangle$ must be given by eq. (3). Applying eq. (5) to this commutator gives us as the Fourier transform of $\mathrm{i}\langle[j_\mu(x),\ \varphi_1(y)]\rangle$ the single term $[k^2 n_\mu - k_\mu(nk)]\rho(k^2, nk)$. We have thus exorcised both Goldstone's zero-mass bosons and the "spurion" state (at $k_\mu = 0$) proposed by Klein and Lee.

In a subsequent note it will be shown, by considering some classical field theories which display broken symmetries, that the introduction of gauge fields may be expected to produce qualitative changes in the nature of the particles described by such theories after quantization.

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BROKEN SYMMETRIES AND THE MASSES OF GAUGE BOSONS

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(Received 31 August 1964)

In a recent note[1] it was shown that the Goldstone theorem,[2] that Lorentz-covariant field theories in which spontaneous breakdown of symmetry under an internal Lie group occurs contain zero-mass particles, fails if and only if the conserved currents associated with the internal group are coupled to gauge fields. The purpose of the present note is to report that, as a consequence of this coupling, the spin-one quanta of some of the gauge fields acquire mass; the longitudinal degrees of freedom of these particles (which would be absent if their mass were zero) go over into the Goldstone bosons when the coupling tends to zero. This phenomenon is just the relativistic analog of the plasmon phenomenon to which Anderson[3] has drawn attention: that the scalar zero-mass excitations of a superconducting neutral Fermi gas become longitudinal plasmon modes of finite mass when the gas is charged.

The simplest theory which exhibits this behavior is a gauge-invariant version of a model used by Goldstone[2] himself: Two real[4] scalar fields φ_1, φ_2 and a real vector field A_μ interact through the Lagrangian density

$$L=-\tfrac{1}{2}(\nabla\varphi_1)^2-\tfrac{1}{2}(\nabla\varphi_2)^2 -V(\varphi_1^{\,2}+\varphi_2^{\,2})-\tfrac{1}{4}F_{\mu\nu}F^{\mu\nu}, \qquad (1)$$

where

$$\nabla_\mu\varphi_1=\partial_\mu\varphi_1-eA_\mu\varphi_2,$$

$$\nabla_\mu\varphi_2=\partial_\mu\varphi_2+eA_\mu\varphi_1,$$

$$F_{\mu\nu}=\partial_\mu A_\nu-\partial_\nu A_\mu,$$

e is a dimensionless coupling constant, and the metric is taken as $-+++$. L is invariant under simultaneous gauge transformations of the first kind on $\varphi_1\pm i\varphi_2$ and of the second kind on A_μ. Let us suppose that $V'(\varphi_0^{\,2})=0$, $V''(\varphi_0^{\,2})>0$; then spontaneous breakdown of U(1) symmetry occurs. Consider the equations [derived from (1) by treating $\Delta\varphi_1$, $\Delta\varphi_2$, and A_μ as small quantities] governing the propagation of small oscillations about the "vacuum" solution $\varphi_1(x)=0$, $\varphi_2(x)=\varphi_0$:

$$\partial^\mu\{\partial_\mu(\Delta\varphi_1)-e\varphi_0A_\mu\}=0, \qquad (2a)$$

$$\{\partial^2-4\varphi_0^{\,2}V''(\varphi_0^{\,2})\}(\Delta\varphi_2)=0, \qquad (2b)$$

$$\partial_\nu F^{\mu\nu}=e\varphi_0\{\partial^\mu(\Delta\varphi_1)-e\varphi_0A^\mu\}. \qquad (2c)$$

Equation (2b) describes waves whose quanta have (bare) mass $2\varphi_0\{V''(\varphi_0^{\,2})\}^{1/2}$; Eqs. (2a) and (2c) may be transformed, by the introduction of new variables

$$B_\mu=A_\mu-(e\varphi_0)^{-1}\partial_\mu(\Delta\varphi_1),$$
$$G_{\mu\nu}=\partial_\mu B_\nu-\partial_\nu B_\mu=F_{\mu\nu}, \qquad (3)$$

into the form

$$\partial_\mu B^\mu=0,\quad \partial_\nu G^{\mu\nu}+e^2\varphi_0^{\,2}B^\mu=0. \qquad (4)$$

Equation (4) describes vector waves whose quanta have (bare) mass $e\varphi_0$. In the absence of the gauge field coupling ($e=0$) the situation is quite different: Equations (2a) and (2c) describe zero-mass scalar and vector bosons, respectively. In passing, we note that the right-hand side of (2c) is just the linear approximation to the conserved current: It is linear in the vector potential, gauge invariance being maintained by the presence of the gradient term.[5]

When one considers theoretical models in which spontaneous breakdown of symmetry under a semisimple group occurs, one encounters a variety of possible situations corresponding to the various distinct irreducible representations to which the scalar fields may belong; the gauge field always belongs to the adjoint representation.[6] The model of the most immediate interest is that in which the scalar fields form an octet under SU(3): Here one finds the possibility of two nonvanishing vacuum expectation values, which may be chosen to be the two $Y=0$, $I_3=0$ members of the octet.[7] There are two massive scalar bosons with just these quantum numbers; the remaining six components of the scalar octet combine with the corresponding components of the gauge-field octet to describe massive vector bosons. There are two $I=\frac{1}{2}$ vector doublets, degenerate in mass between $Y=\pm1$ but with an electromagnetic mass splitting between $I_3=\pm\frac{1}{2}$, and the $I_3=\pm1$ components of a $Y=0$, $I=1$ triplet whose mass is entirely electromagnetic. The two $Y=0$, $I=0$ gauge fields remain massless: This is associated with the residual unbroken symmetry under the Abelian group generated by Y and I_3. It may be expected that when a further mechanism (presumably related to the weak interactions) is introduced in order to break Y conservation, one of these gauge fields will acquire mass, leaving the photon as the only massless vector particle. A detailed discussion of these questions will be presented elsewhere.

It is worth noting that an essential feature of the type of theory which has been described in this note is the prediction of incomplete multiplets of scalar and vector bosons.[8] It is to be expected that this feature will appear also in theories in which the symmetry-breaking scalar fields are not elementary dynamic variables but bilinear combinations of Fermi fields.[9]

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[1]P. W. Higgs, to be published.

[2]J. Goldstone, Nuovo Cimento 19, 154 (1961); J. Goldstone, A. Salam, and S. Weinberg, Phys. Rev. 127, 965 (1962).

[3]P. W. Anderson, Phys. Rev. 130, 439 (1963).

[4]In the present note the model is discussed mainly in classical terms; nothing is proved about the quantized theory. It should be understood, therefore, that the conclusions which are presented concerning the masses of particles are conjectures based on the quantization of linearized classical field equations. However, essentially the same conclusions have been reached independently by F. Englert and R. Brout, Phys. Rev. Letters 13, 321 (1964): These authors discuss the same model quantum mechanically in lowest order perturbation theory about the self-consistent vacuum.

[5]In the theory of superconductivity such a term arises from collective excitations of the Fermi gas.

[6]See, for example, S. L. Glashow and M. Gell-Mann, Ann. Phys. (N.Y.) 15, 437 (1961).

[7]These are just the parameters which, if the scalar octet interacts with baryons and mesons, lead to the Gell-Mann–Okubo and electromagnetic mass splittings: See S. Coleman and S. L. Glashow, Phys. Rev. 134, B671 (1964).

[8]Tentative proposals that incomplete SU(3) octets of scalar particles exist have been made by a number of people. Such a rôle, as an isolated $Y=\pm1$, $I=\frac{1}{2}$ state, was proposed for the κ meson (725 MeV) by Y. Nambu and J. J. Sakurai, Phys. Rev. Letters 11, 42 (1963). More recently the possibility that the σ meson (385 MeV) may be the $Y=I=0$ member of an incomplete octet has been considered by L. M. Brown, Phys. Rev. Letters 13, 42 (1964).

[9]In the theory of superconductivity the scalar fields are associated with fermion pairs; the doubly charged excitation responsible for the quantization of magnetic flux is then the surviving member of a U(1) doublet.

SPLITTING OF THE 70-PLET OF SU(6)

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(Received 18 September 1964)

1. In a previous note,[1] hereafter called I, we proposed an expression for the mass operator responsible for lifting the degeneracies of spin-unitary spin supermultiplets [Eq. (31)-I]. The purpose of the present note is to apply this expression to the 70-dimensional representation of SU(6).

The importance of the 70-dimensional representation has already been underlined by Pais.[2] Since

$$\underline{35}\otimes\underline{56}=\underline{56}\oplus\underline{70}\oplus\underline{700}\oplus\underline{1134}, \qquad (1)$$

it follows that $\underline{70}$ is the natural candidate for accommodating the higher meson-baryon resonances. Furthermore, since the SU(3)⊗SU(2) content is

$$\underline{70}=(\underline{1},\underline{2})+(\underline{8},\underline{2})+(\underline{10},\underline{2})+(\underline{8},\underline{4}), \qquad (2)$$

we may assume that partial occupancy of the $\underline{70}$ representation has already been established through the so-called γ octet[2] $(\frac{3}{2})^-$. Recent experiments appear to indicate that some $(\frac{1}{2})^-$ states may also be at hand.[3] With six masses at one's disposal, our formulas can predict the masses of all the other occupants of $\underline{70}$ and also provide a consistency check on the input. Our discussion of the $\underline{70}$ representation thus appears to be of immediate physical interest.

1.2 Ideas About Mass

As introduced by Newton mass was mechanical. The first ideas on dynamical mass were due to Poincare (Poincare Stresses). However all classical so far.

With quantum field theory mass could change through self interactions (radiative corrections to the self-energy - Lamb shift) to give $m = m_0 + \delta m$, or through a change in vacuum (BCS) $E = p^2/2m + \Delta$ where Δ is self-consistent gap parameter.

Then through Nambu (1960) and Goldstone (1961) the possibility arose that all mass could come from self interaction, and especially so for gauge bosons (Anderson 1957, 1963), culminating in the Weinberg (1967), Salam (1968) and Glashow (1961, 1970) renormalizable $SU(2) \times U(1)$ theory of electroweak interactions and the discovery first of weak neutral currents (1973), then charmed particles (1974), then the intermediate vector bosons of the weak interactions (1983), and finally the Higgs boson (2012). All of this is possible because of Dirac's Hilbert space formulation of quantum mechanics in which one sets $\psi(x) = \langle x|\psi\rangle$, with the physics being in the properties of the states $|\psi\rangle$. We thus live in Hilbert space and not in coordinate space, and not only that, there is altogether more in Hilbert space than one could imagine, such as half-integer spin and collective macroscopic quantum systems such as superconductors and superfluids. In this Hilbert space we find an infinite Dirac sea of negative energy particles. This large number of degrees of freedom can act collectively to provide the dynamics needed to produce mass generation and the Higgs boson.

Outline of the talk:

(1) Introduction

(2) The Higgs Boson Discovery

(3) Broken Symmetry

(4) Theoretical Background Leading to the Higgs Boson Papers in 1964

(5) What Exactly is the Higgs Field?

(6) What Comes Next?

(7) The Moral of the Story

2 The Higgs Boson Discovery

So now to the Higgs boson discovery. First the announcement from CERN on July 4, 2012, then the announcements by the experimental groups ATLAS and CMS at the Large Hadron Collider. From amongst a set of

(a) 10^{15} proton-proton collisions produced at the Large Hadron Collider

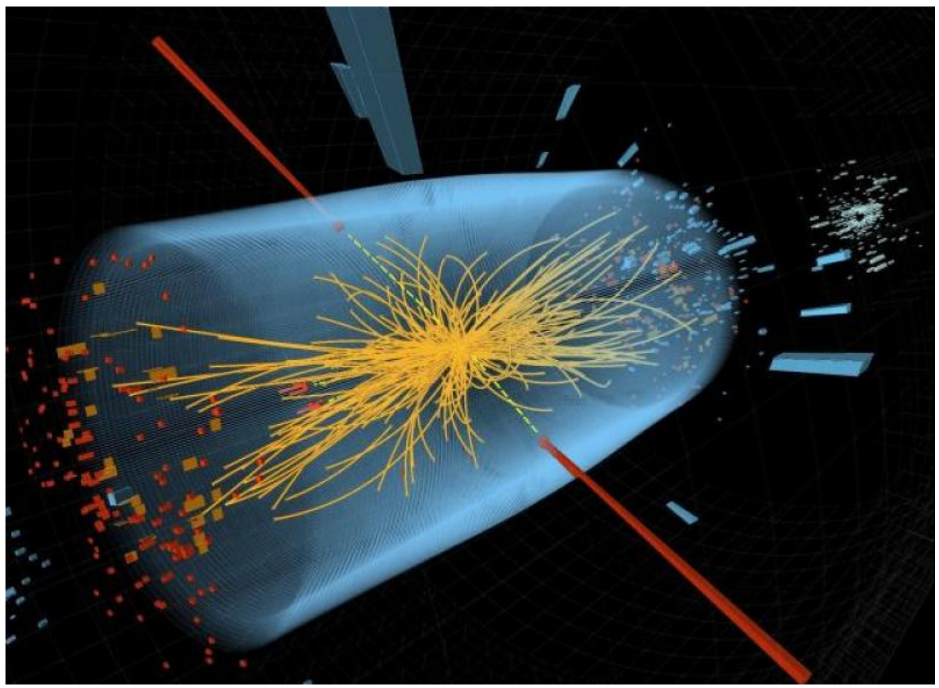
(b) of the order of 240,000 collisions produce a Higgs boson

(c) of which just 350 decay into pairs of gamma rays and

(d) just 8 decay into a pair of leptons.

The search for the Higgs boson is thus a search for some very rare events.

Higgs within reach



The ATLAS (http://www.atlas.ch/news/2012/latest-results-from-higgs-search.html) and CMS (http://cms.web.cern.ch/news/observation-new-particle-mass-125-gev) experiments at CERN today presented their latest results in the search for the long-sought Higgs boson (http://press.web.cern.ch/press-releases/2012/07/cern-experiments-observe-particle-consistent-long-sought-higgs-boson) . Both experiments see strong indications for the presence of a new particle, which could be the Higgs boson, in the mass region around 126 gigaelectronvolts (GeV).

The experiments found hints of the new particle by analysing trillions of proton-proton collisions from the Large Hadron Collider (LHC) (/about/accelerators/large-hadron-collider) in 2011 and 2012. The Standard Model (/about/physics/standard-model) of particle physics predicts that a Higgs boson would decay into different particles – which the LHC experiments then detect.

Both ATLAS and CMS gave the level of significance of the result as 5 sigma on the scale that particle physicists use to describe the certainty of a discovery. One sigma means the results could be random fluctuations in the data, 3 sigma counts as an observation and a 5-sigma result is a discovery. The results presented today are preliminary, as the data from 2012 is still under analysis. The complete analysis is expected to be published around the end of July.

Read the CERN press release→ (http://press.web.cern.ch/press/PressReleases/Releases2012/PR17.12E.html)

Find out more

- What is the Higgs boson? (http://cdsweb.cern.ch/record/1458922)
- How do physicists look for it? (http://cdsweb.cern.ch/record/1458015)
- What comes next? (http://press.web.cern.ch/backgrounders/higgs-update-4-july)
- Multimedia (http://cern.ch/go/higgs-1207-media)
- ICHEP2012 (http://www.ichep2012.com.au/)

Posted by Cian O'Luanaigh (http://profiles.web.cern.ch/717206) on 4 Jul 2012. Last updated 7 Oct 2013, 14.58.

More like this

- (/about/updates/2013/12/highlights-cern-2013) Highlights from CERN in 2013 (/about/updates/2013/12/highlights-cern-2013) 19 Dec 2013

ATLAS and the Higgs

Published October 2012 (Updated April 2013)

At a seminar held on 4 July 2012, the ATLAS experiment announced that it had observed a new particle: a boson consistent with the Higgs boson. The excess of signal over background was observed at a mass of around 126 GeV, and the level of confidence in the results was calculated to be 5 sigma. (You can find an explanation of GeV here and more information about sigma here).

At the same seminar, ATLAS' sister experiment on the LHC (Large Hadron Collider), CMS, announced very similar results. The similarity acts as verification: if one experiment saw something very different to the other, there would be doubts about the results.

In March 2013, in the light of the updated ATLAS and CMS results, CERN announced that the new particle was indeed a Higgs boson. Having analyzed two and a half times more data than was available for the discovery announcement in July, the confidence of observation has risen to 10 sigma. The experiments were also able to show that the properties of the particle as well as the ways it interacts with other particles were well-matched with those of a Higgs boson, which is expected to have spin 0 and parity +. Physicists have now to pursue their measurements to determine if this Higgs particle corresponds to the Standard Model Higgs boson or if it is part of a new physics scenario.

ATLAS' role

ATLAS is located at Point 1 of the LHC, which accelerates proton beams to high energy and then collides them head-on at four different points along its 27-km ring. As a "general purpose" detector, it is designed to identify and measure many different types of particles produced in these collisions. From the data captured, ATLAS physicists are able to study a broad variety of interesting physics topics and to search for new phenomena, such as the Higgs boson. (You can find a description of the ATLAS detector here).

Thanks to the particularly impressive performance of the LHC in producing collisions during 2012, and the detector's very high data-taking efficiency (nearly 96%), ATLAS was able to record nearly 22 inverse femtobarns of data during 2012 to add to the 4.8 inverse femtobarns it recorded in 2011.

To obtain the high quantity of data, the LHC attained very high instantaneous luminosities. This means that there were many more proton collisions occurring at essentially the same time in the detector, an effect known as "event pile-up", making it more complex to process and analyse the data. Fortunately, the quality of data taken during that time was excellent, so ATLAS physicists were able to take advantage of the additional data to make the discovery after little more than two years of LHC operation.

Why is this important to mankind?

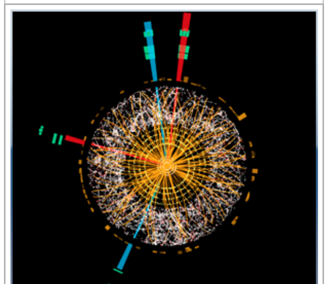
This result is an important advance in our understanding of the basic forces holding the universe together. In particular this new boson provides support for the existence of the proposed Higgs field, which explains how some particles come to have mass and others don't. Without mass, all particles would fly around freely and matter as we know it would not exist.

Physicists work to a theory of fundamental particles and their interactions called the Standard Model, which was first proposed in the 1970s. So far experiments have been able to confirm the existence of nearly all its elements with a high degree of precision. The Higgs boson, however, had eluded detection until now, prompting speculation that the theory could be incomplete. The findings so far suggest a Higgs boson compatible with the Standard Model, but further studies are needed to confirm this.

Videos

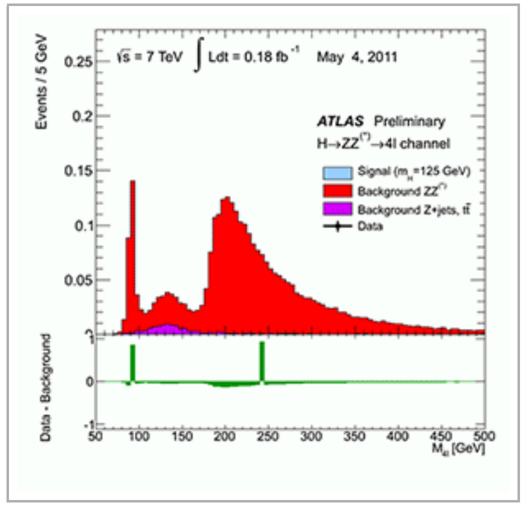
Photos

Proton Collisions



Higgs Plots

There is a more philosophical reason for the importance of this observation. Human beings have a capacity for abstract thinking and reasoning that goes beyond solving only our immediate needs. This scientific investigation and the large, complex apparatus needed to make it happen are examples of our unique human ability and drive to find out "why?". This drive forms the basis of our civilisation, producing knowledge and tools for future generations.



Pursuit of new physics

Up to now, the more detailed studies of the newly discovered particle's properties reveal it to be compatible with the Standard Model Higgs boson. However, scientists are looking for more Higgs particles which, according to almost all high energy extensions of the Standard Model, should exist. Some of the most popular new models of physics are the so-called supersymmetry theories, which could potentially solve a number of problems in theoretical physics. The most minimalist supersymmetry theory predicts at least five (!) Higgs bosons: three neutral and two charged. So in the future if we detect more than one, we will know that we are looking at new physics!

LINKS

ATLAS and the Higgs: Resources

Compact Muon Solenoid experiment at CERN's LHC

Observation of a New Particle with a Mass of 125 GeV

CMS Experiment, CERN
4 July 2012

The statement below is also available as PDF files in: English | French | Catalan | Chinese (traditional) | Chinese (simplified) | Croatian | Dutch | Finnish | Galician | German | Greek | Hindi | Hungarian | Italian | Korean | Persian | Polish | Portuguese | Romanian | Russian | Serbian | Spanish | Mexican Spanish | Turkish | Urdu

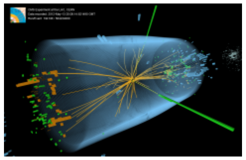


Figure 1. Event recorded with the CMS detector in 2012 at a proton-proton centre-of-mass energy of 8 TeV. The event shows characteristics expected from the decay of the SM Higgs boson to a pair of photons (dashed yellow lines and green towers). The event could also be due to known standard model background processes.
DOWNLOAD high-resolution images

SUMMARY

In a joint seminar today at CERN and the "ICHEP 2012" conference[1] in Melbourne, researchers of the Compact Muon Solenoid (CMS) experiment at the Large Hadron Collider (LHC) presented their preliminary results on the search for the standard model (SM) Higgs boson in their data recorded up to June 2012. CMS observes an excess of events at a mass of approximately 125 GeV[2] with a statistical significance of five standard deviations (5 sigma)[3] above background expectations. The probability of the background alone fluctuating up by this amount or more is about one in three million. The evidence is strongest in the two final states with the best mass resolution: first the two-photon final state and second the final state with two pairs of charged leptons (electrons or muons). We interpret this to be due to the production of a previously unobserved particle with a mass of around 125 GeV. The CMS data also rule out the existence of the SM Higgs boson in the ranges 110–122.5 GeV and 127–600 GeV with 95% confidence level[4] – lower masses were already excluded by CERN's LEP collider at the same confidence level. Within the statistical and systematic uncertainties, results obtained in the various search channels are consistent with the expectations for the SM Higgs boson. However, more data are needed to establish whether this new particle has all the properties of the SM Higgs boson or whether some do not match, implying new physics beyond the standard model. The LHC continues to deliver new data at an impressive rate. By the end of 2012, CMS hopes to have more than triple its total current data sample. These data will enable CMS to elucidate further the nature of this newly observed particle. They will also allow CMS to extend the reach of their many other searches for new physics.

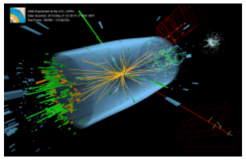


Figure 2. Event recorded with the CMS detector in 2012 at a proton-proton centre-of-mass energy of 8 TeV. The event shows characteristics expected from the decay of the SM Higgs boson to a pair of Z bosons, one of which subsequently decays to a pair of electrons (green lines and green towers) and the other Z decays to a pair of muons (red lines). The event could also be due to known standard model background processes.
DOWNLOAD high-resolution images

CMS SEARCH STRATEGY

CMS analysed the full data sample of proton-proton collisions collected in all of 2011 and in 2012, up until June 18. These data amount to up to 5.1 fb^{-1} of integrated luminosity[5], at a centre-of-mass energy of 7 TeV in 2011 and up to 5.3 fb^{-1} at 8 TeV in 2012. The standard model predicts that the Higgs boson lasts for only a very short time before it breaks up, or "decays", into other well-known particles. CMS studied five main Higgs boson decay channels. Three channels result in pairs of bosonic particles (γγ, ZZ or WW) and two channels result in pairs of fermionic particles (bb or ττ), where γ denotes a photon, Z and W denote the force carriers of the weak interaction, b denotes a bottom quark, and τ denotes a tau lepton. The γγ, ZZ and WW channels are equally sensitive in the search for a Higgs boson around 125 GeV and all are more sensitive than the bb and ττ channels. The γγ and ZZ channels are especially important as they both allow the mass of the new particle to be measured with precision. In the γγ channel the mass is determined from the energies and directions of two high-energy photons measured by the CMS crystal electromagnetic calorimeter (ECAL, Figure 1). In the ZZ channel the mass is determined from the decays of the two Zs to two pairs of electrons, or two pairs of muons, or a pair of electrons and a pair of muons (Figure 2). These are measured in the ECAL, inner tracking and muon detectors. The WW channel is more complex. Each W is identified through its decay to an electron and a neutrino or a muon and a neutrino. The neutrinos pass through the CMS detectors undetected, so the SM Higgs boson in the WW channel would manifest itself as a broad excess in the mass distribution, rather than a narrow peak. The bb channel has large backgrounds from standard model processes, so the analysis searches for events in which a Higgs boson is produced in association with a W or Z, which then decays to electron(s) or muon(s). The ττ channel is measured by observing τ decays to electrons, muons and hadrons.

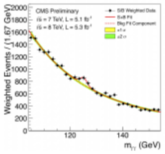


Figure 3. Di-photon (γγ) invariant mass distribution for the CMS data of 2011 and 2012 (black points with error bars). The data are weighted by the signal to background ratio for each sub-category of events. The solid red line shows the fit result for signal plus background; the dashed red line shows only the background.
DOWNLOAD this plot

CMS SEARCH RESULTS

The CMS data sample should be sensitive enough to completely exclude the mass range 110–600 GeV at 95% confidence level, if the SM Higgs does not exist. In fact, the CMS data do rule out the existence of the SM Higgs boson in two broad mass ranges of 110–122.5 GeV and 127–600 GeV with 95% confidence level. The range of 122.5–127 GeV cannot be excluded because we see an excess of events in three of the five channels analysed:

- **γγ channel:** the γγ mass distribution is shown in Figure 3. There is an excess of events above background with a significance of 4.1 sigma, at a mass near 125 GeV. The observation of the two-photon final state implies that the new particle is a boson, not a fermion, and that it cannot be a "spin 1" particle.
- **ZZ channel:** Figure 4 shows the mass distribution for the four leptons (two pairs of electrons, or two pairs of muons, or the pair of electrons and the

pair of muons). Accounting also for the decay angle characteristics, it yields an excess of 3.2 sigma above background at a mass near 125 GeV.
- **WW channel:** a broad excess in the mass distribution of 1.5 sigma is observed.
- **bb and ττ channels:** no excess is observed.

The statistical significance of the signal, from a combined fit to all five channels (Figure 5), is 4.9 sigma above background. A combined fit to just the two most sensitive and high-resolution channels (γγ and ZZ) yields a statistical significance of 5.0 sigma. The probability of the background alone fluctuating up by this amount or more is about one in three million. The mass of the new particle is determined to be 125.3 ± 0.6 GeV, independent of any assumptions about the expected relative yields of the decay channels. The measured production rate (σ_{DAT}) of this new particle is consistent with the predicted rate (σ_{SM}) for the SM Higgs boson: $\sigma_{DAT}/\sigma_{SM} = 0.80 \pm 0.22$. Great care has also been taken to understand numerous details of the detector performance, the event selection, background determinations and other possible sources of systematic and statistical uncertainties. The 2011 analysis[6] showed an excess of events at about 125 GeV. Therefore, to avoid a potential bias in the choice of selection criteria for the 2012 data that might artificially enhance this excess, the 2012 data analysis was performed "blind"[7], meaning that the region of interest was not examined until after all the analysis criteria had been fully scrutinized and approved. As a general cross-check, the analyses were performed by at least two independent teams. A number of other features reinforce confidence in the results:

- The excess is seen at around 125 GeV in both the 2011 data sample (7 TeV) and the 2012 data sample (8 TeV);
- The excess is seen at the same mass in both the high-resolution channels (γγ and ZZ);
- The excess seen in the WW is consistent with one that would arise from a particle at 125 GeV;
- The excess is seen in a range of final states involving photons, electrons, muons and hadrons.

The preliminary results presented today will be refined, with the aim of submitting them for publication towards the end of the summer.

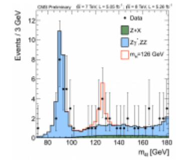


Figure 4. Distribution of the four-lepton reconstructed mass for the sum of the 4e, 4μ, and 2e2μ channels. Points represent the data, shaded histograms represent the background and un-shaded histogram the signal expectations. The distributions are presented as stacked histograms. The measurements are presented for the sum of the data collected at centre-of-mass energies of 7 TeV and 8 TeV.
DOWNLOAD this plot

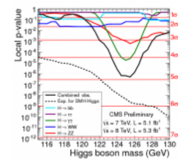


Figure 5. The observed probability (local p-value) that the background-only hypothesis would yield the same or more events as are seen in the CMS data, as a function of the SM Higgs boson mass for the five channels considered. The solid black line shows the combined local p-value for all channels.
DOWNLOAD this plot

FUTURE PLANS

The new particle observed at about 125 GeV is compatible, within the limited statistical accuracy, with being the SM Higgs boson. However, more data are required to measure its properties such as decay rates in the various channels (γγ, ZZ, WW, bb and ττ) and ultimately its spin and parity, and hence ascertain whether it is indeed the SM Higgs boson or the result of new physics beyond the standard model. The LHC continues to perform extremely well. By the end of 2012, CMS expects to more than triple its total data sample, and hence to probe further the nature of this new particle. If this particle is indeed the SM Higgs boson, its properties and implications for the standard model will be studied in detail. If it is not the SM Higgs boson, CMS will explore the nature of the new physics that it implies, which may include additional particles that are observable at the LHC. In either case, searches will also continue for other new particles or forces that can be observed in future runs of the LHC at higher beam energies and intensities.

More information

- **This statement is available online at:** http://cms.web.cern.ch/news/observation-new-particle-mass-125-gev
- CMS public website: http://cern.ch/cms
- CMS Higgs Seminar at CERN (Prof. Joseph Incandela)
- CERN Press Release, 4 July 2012 [Also available in French]
- ATLAS Experiment Higgs search results
- ATLAS Higgs Seminar at CERN (Prof. Fabiola Gianotti)

- Full video of the seminar
- Full video of the press conference
- Photographs from the seminar and press conference
- Photographs from the webcast at ICHEP

EVENT IMAGES AND ANIMATION OF REAL CMS COLLISIONS

Images:

- Two photons i.e. γγ (8 TeV) event display
- ZZ to two electrons and two muons (8 TeV) event display
- All images from this statement (including plots)

Animations:

- Two photons i.e. γγ (8 TeV) event display animation: CDS | YouTube
- ZZ to two electrons and two muons (8 TeV) event display animation: CDS | YouTube
- ZZ to 4 muons (7 TeV) event display animation: CDS | YouTube

SHORT MOVIES ABOUT THE HIGGS

- Animation of the Higgs mechanism [with subtitles] See also [without subtitles]
- What is the Higgs? by Don Lincoln
- Higgs boson: How do we search for it? by Don Lincoln

REFERENCES

3 Broken Symmetry

3.1 Global Discrete Symmetry – Real scalar field – Goldstone (1961)

Consider a real scalar field (just one degree of freedom) with a potential energy in the form of a double-well potential, viz. a potential shaped like a letter W with two wells:

$$V(\phi) = \frac{1}{4}\lambda^2\phi^4 - \frac{1}{2}\mu^2\phi^2. \tag{1}$$

This potential has a discrete symmetry under $\phi \to -\phi$, and its derivative is given by

$$\frac{dV(\phi)}{d\phi} = \lambda^2\phi^3 - \mu^2\phi. \tag{2}$$

The potential has a local maximum at $\phi = 0$ where $V(\phi = 0)$ is zero and $d^2V(\phi = 0)/d\phi^2$ is negative, and two-fold **degenerate** global minima at $\phi = +\mu/\lambda$ and $\phi = -\mu/\lambda$ where $V(\phi = \pm\mu/\lambda)$ is equal to $-\mu^4/4\lambda^2$ and $d^2V(\phi = \pm\lambda/\mu)/d\phi^2$ is positive. Since $\phi = 0$ is a local maximum, if we consider small oscillations around $\phi = 0$ of the form $\phi = 0 + \chi$ we generate a negative quadratic term $-(1/2)\mu^2\chi^2$ and thus negative mass squared, viz. $m^2 = -\mu^2$, a so-called tachyon. The tachyon signals an instability of the configuration with $\phi = 0$ (i.e. we roll away from the top of the hill).

However if we fluctuate around either global minimum (i.e. we oscillate in the vertical around the bottom of either hill) by setting $\phi = \pm\mu/\lambda + \chi$ we get

$$V(\phi) = -\frac{\mu^4}{4\lambda^2} + \mu^2\chi^2 \pm \mu\lambda\chi^3 + \frac{1}{4}\lambda^2\chi^4. \tag{3}$$

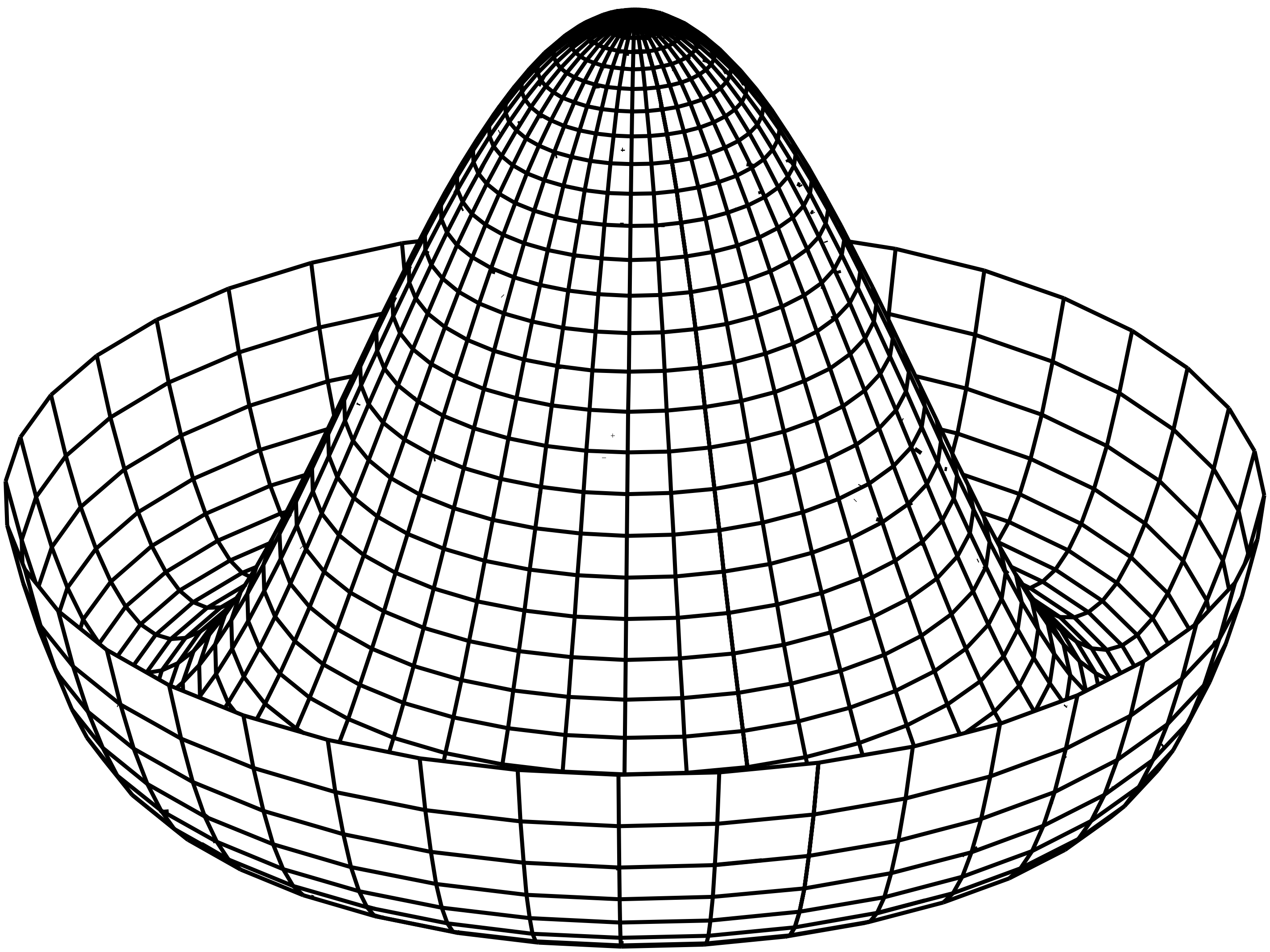
The field χ now has the positive $m^2 = +2\mu^2$ $(= -2 \times m^2(\text{tachyon}))$. Thus (Goldstone 1961) the would-be tachyonic particle becomes a massive particle, and with one scalar field we obtain one particle. This particle is the **Higgs boson** in embryo.

The $-\mu^4/4\lambda^2$ potential term contributes to the **cosmological constant**, and would be unacceptably large (10^{60} times too large) if the boson has the 126 GeV mass that the Higgs boson has now been found to have.

All this arises because the minimum is two-fold degenerate, and picking either one breaks the symmetry spontaneously. Just like people at dinner. Each one can take the cup to their left or their right, but once one person has done so, the rest have no choice. However a person at the opposite end of the table may not know what choice was made at the other end of the table and may make the opposite choice of cup, and thus persons in the middle could finish up with no cup. To ensure that this does not happen we need long range correlations – hence massless Goldstone boson.

Figure 1: Double Well, Mexican Hat Potential

http://upload.wikimedia.org/wikipedia/commons/7/7b/Mexican...



3.2 Global Continuous Symmetry – Complex scalar field – Goldstone (1961)

Consider a complex scalar field (two degrees of freedom) $\phi = \phi_1 + i\phi_2 = re^{i\theta}$, $\phi^*\phi = \phi_1^2 + \phi_2^2 = r^2$, with a potential energy in the shape of a rotated letter W or a broad-brimmed, high-crowned Mexican Hat (as shown in the figure):

$$V(\phi) = \frac{1}{4}\lambda^2(\phi^*\phi)^2 - \frac{1}{2}\mu^2\phi^*\phi = \frac{1}{4}\lambda^2(\phi_1^2 + \phi_2^2)^2 - \frac{1}{2}\mu^2(\phi_1^2 + \phi_2^2). \tag{4}$$

The potential has a continuous global symmetry of the form $\phi \to e^{i\alpha}\phi$ with constant α, and its derivative is given by

$$\frac{dV(\phi)}{d\phi_1} = \lambda^2\phi_1^3 + \lambda^2\phi_2^2\phi_1 - \mu^2\phi_1, \qquad \frac{dV(\phi)}{d\phi_2} = \lambda^2\phi_2^3 + \lambda^2\phi_1^2\phi_2 - \mu^2\phi_2. \tag{5}$$

This potential has a local maximum at $\phi_1 = 0, \phi_2 = 0$ where $V(\phi = 0)$ is zero, and infinitely **degenerate** global minima at $\phi_1^2 + \phi_2^2 = \mu^2/\lambda^2$ (the entire 360 degree valley or trough between the brim and the crown of the Mexican hat). Again we would have a tachyon if we expand around the local maximum, only this time we would get two. So fluctuate around any one of the global minima by setting $\phi_1 = \mu/\lambda + \chi_1$, $\phi_2 = \chi_2$, to get

$$V(\phi) = -\frac{\mu^4}{4\lambda^2} + \mu^2\chi_1^2 + \mu\lambda\chi_1^3 + \frac{1}{4}\lambda^2\chi_1^4 + \mu\lambda\chi_1\chi_2^2 + \frac{\lambda^2}{2}\chi_1^2\chi_2^2. \tag{6}$$

The (embryonic) Higgs boson field is now χ_1 with $m^2 = +2\mu^2$. However, the field χ_2 no has **no mass at all** (it corresponds to horizontal oscillations along the valley floor), and is known as a **Goldstone boson** (Goldstone 1961). Thus from a complex scalar field we obtain two particles. Since the Goldstone boson is massless, it travels at the speed of light. It is thus intrinsically relativistic, and being massless can provide for long range correlations. Moreover, if such particles exist then they could generate fermion masses entirely dynamically (Nambu 1960, Nambu and Jona-Lasinio 1961), with the pion actually serving this purpose.
The $-\mu^4/4\lambda^2$ potential term remains and the cosmological constant problem is just as severe as before.

Now if have massless particles, would get long range forces (just like photons), and yet nuclear and weak forces are short range. So what can we do about such Goldstone bosons. Two possibilities - they could get some mass because symmetry is not exact (pion mass), or we could get rid of them altogether (the Higgs mechanism).

3.3 Local Continuous Symmetry – Complex Scalar Field and Gauge Field – Higgs (1964)

Consider a complex scalar field (two degrees of freedom) coupled to a massless vector gauge boson (another two degrees of freedom) for a total of four degrees of freedom in all, viz. $\phi = re^{i\theta}$, $\phi^*\phi = r^2$, $F_{\mu\nu} = \partial_\mu A_\nu - \partial_\nu A_\mu$. The system has kinetic energy K and potential energy V:

$$\begin{aligned} K &= \frac{1}{2}(-i\partial_\mu + eA_\mu)(re^{-i\theta})(i\partial^\mu + eA^\mu)(re^{i\theta}) - \frac{1}{4}F_{\mu\nu}F^{\mu\nu} = \frac{1}{2}\partial_\mu r\partial^\mu r + \frac{1}{2}r^2(eA_\mu - \partial_\mu\theta)(eA^\mu - \partial^\mu\theta) - \frac{1}{4}F_{\mu\nu}F^{\mu\nu}, \\ V(\phi) &= \frac{1}{4}\lambda^2(\phi^*\phi)^2 - \frac{1}{2}\mu^2\phi^*\phi = \frac{1}{4}\lambda^2 r^4 - \frac{1}{2}\mu^2 r^2, \end{aligned} \tag{7}$$

and because of the gauge boson the system is now invariant under continuous **local** gauge transformations of the form $\phi \to e^{i\alpha(x)}\phi$, $eA_\mu \to eA_\mu + \partial_\mu\alpha(x)$ with spacetime dependent $\alpha(x)$. With derivative

$$dV(\phi)/dr = \lambda^2 r^3 - \mu^2 r \tag{8}$$

the potential has a local maximum at $r = 0$ where $V(r = 0)$ is zero and **degenerate** global minima at $r = \mu/\lambda$ (infinitely degenerate since independent of θ). Again we would have two tachyons if we expand around the local maximum. So fluctuate around the global minimum by setting $r = \mu/\lambda + \chi_1$, $\theta_2 = \chi_2$. On defining $B_\mu = A_\mu - (1/e)\partial_\mu\chi_2$ we obtain

$$\begin{aligned} K &= \frac{1}{2}\partial_\mu\chi_1\partial^\mu\chi_1 + \frac{e^2}{2}\left(\frac{\mu^2}{\lambda^2} + \frac{2\mu}{\lambda}\chi_1 + \chi_1^2\right)B_\mu B^\mu - \frac{1}{4}(\partial_\mu B_\nu - \partial_\nu B_\mu)(\partial^\mu B^\nu - \partial^\nu B^\mu), \\ V(\phi) &= -\frac{\mu^4}{4\lambda^2} + \mu^2\chi_1^2 + \mu\lambda\chi_1^3 + \frac{1}{4}\lambda^2\chi_1^4. \end{aligned} \tag{9}$$

There is again a Higgs boson field χ_1 with $m^2 = +2\mu^2$. However, the field χ_2 has disappeared completely. Instead the vector boson now has a **non-zero mass** given by $m = e\mu/\lambda$. Since a massive gauge boson has three degrees of freedom (two transverse and one longitudinal) while a massless gauge boson such as the photon only has two transverse degrees of freedom, the would-be massless Goldstone boson is absorbed into the now massive gauge boson to provide its needed longitudinal degree of freedom. Hence a massless Goldstone boson and a massless gauge boson are replaced by one massive gauge boson, with two long-range interactions being replaced by one short range interaction. This is known as the Higgs mechanism (1964) though it is due initially to Anderson (1957). The fourth of the original four degrees of freedom becomes the massive Higgs boson, and its presence is an indicator that the Higgs mechanism has taken place. However, the cosmological constant problem remains as severe as before.

4 Theoretical Background Leading to the Higgs Boson Papers in 1964

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For the interested reader a nice discussion of the history and development of the Higgs mechanism may be found in Frank Close's book "The Infinity Puzzle: The personalities, politics, and extraordinary science behind the Higgs boson".

4.2 Theoretical Background

In order to characterize macroscopic ordered phases in a general way Landau introduced the concept of a macroscopic order parameter ϕ. For a ferromagnet for instance ϕ would represent the spontaneous magnetization M and would thus be a matrix element of a field operator in an ordered quantum state that described the ordered magnetic phase. Building on this approach Landau and Ginzburg (1950) wrote down a Lagrangian for such a ϕ for a superconductor, with its potential being of the generic form $V(\phi) = (T-T_C)^2\phi^4+(T-T_C)\phi^2$ where T_C is the critical temperature. For temperatures above the critical temperature the potential would have the shape of a single well, viz. like the letter U, with the coefficient of the ϕ^2 term being positive. For temperatures below the critical temperature the potential would have the shape of a double well, viz. like the letter W, with the coefficient of the ϕ^2 term being negative. At the minimum of the potential the order parameter would be zero above the critical temperature (normal phase with state vector $|N\rangle$ in which $\langle N|\phi|N\rangle$ is zero). Below the critical temperature the order parameter would be non-zero (superconducting state $|S\rangle$ in which $\langle S|\phi|S\rangle$ is non-zero).

In 1957 Bardeen, Cooper and Schrieffer (BCS) developed a macroscopic theory of superconductivity based on Cooper pairing of electrons in the presence of a filled Fermi sea of electrons, and explicitly constructed the state $|S\rangle$. In this state the matrix element $\langle S|\psi(x)\psi(x)|S\rangle$ was equal to a spacetime independent function Δ, the gap parameter, which led to a mass shift to electrons propagating in the superconductor of the form $E = p^2/2m + \Delta$. The gap parameter Δ would be temperature dependent and would only be non-zero below the critical temperature. In 1959 Gorkov was able to derive the Ginzburg-Landau Lagrangian starting from the BCS theory and identify the order parameter as $\phi(x) = \langle C|\psi(x)\psi(x)|C\rangle$ where $|C\rangle$ is a coherent state in the Hilbert space based on $|S\rangle$. In the superconducting case then ϕ is not itself a quantum-field-theoretic operator (viz. a q-number operator that would have a canonical conjugate with which it would not commute) but is instead a c-number matrix element of a q-number field operator in a macroscopic coherent quantum state.

In 1957 Anderson used the BCS theory to explain the Meissner effect, an effect in which electromagnetism becomes short range inside a superconductor, with photons propagating in it becoming massive. The effect was one of spontaneous breakdown of local gauge invariance, and was explored in detail by Anderson (1957, 1963) and Nambu (1960).

In parallel with these studies Nambu (1960), Goldstone (1961), and Nambu and Jona-Lasinio (1961) explored the spontaneous breakdown of some continuous global symmetries and showed that collective massless excitations (Goldstone bosons) were generated and that the analog gap parameter would provide for dynamically induced fermion masses. In 1962 Goldstone, Salam and Weinberg showed that there would always be massless Goldstone bosons in any Lorentz invariant theory in which a continuous global symmetry was spontaneously broken. While one could avoid this outcome if the symmetry was also broken in the Lagrangian, as was thought to be the case for the pion, a non-massless but near Goldstone particle (i.e. one with broken symmetry suppressed couplings to matter at low energies), in general the possible presence of massless Goldstone bosons was a quite problematic outcome because it would imply the existence of non-observed long range forces.

In 1962 Schwinger raised the question of whether gauge invariance actually required that photons be massless, and noted for the inverse photon propagator $D^{-1}(q^2) = q^2 - q^2\Pi(q^2)$ that if the vacuum polarization $\Pi(q^2)$ had a massless pole of the form $\Pi(q^2) = m^2/q^2$, then $D^{-1}(q^2)$ would behave as the massive particle $D^{-1}(q^2) = q^2 - m^2$. A massless Goldstone boson could thus produce a massive vector boson.

With Anderson having shown that a photon would become become massive in a superconductor, there was a spirited discussion in the literature (Anderson 1963, Klein and Lee 1964, Gilbert 1964) as to whether an effect such as this might hold in a relativistic theory as well or whether it might just have been an artifact of the fact that the BCS theory was non-relativistic. With the work of Englert and Brout (1964) and Higgs (1964), and then Guralnik, Hagen and Kibble (1964) (reproduced below) the issue was finally resolved, with it being established that in the relativistic case the Goldstone theorem did not in fact hold if there was a spontaneous breakdown of a continuous local theory, with the would-be Goldstone boson no longer being an observable massless particle but instead combining with the initially massless vector boson to produce a massive vector boson. Technically, this mechanism should be known as the Anderson, Englert, Brout, Higgs, Guralnik, Hagen, Kibble mechanism, and while it has undergone many name variations over time, it is now commonly called the Higgs mechanism. What set Higgs' work apart from the others was that in his 1964 Physical Review Letter paper Higgs noted that as well as a massive gauge boson there should also be an observable massive scalar boson, this being the Higgs boson.

At the time of its development in 1964 there was not much interest in the Higgs mechanism, with all of the Englert and Brout, Higgs, and Guralnik, Hagen and Kibble papers getting hardly any citations during the 1960s at all. The primary reason for this was that at the time there was little interest in Yang-Mills theories in general, broken or unbroken, and not only was there no experimental indication at all that one should consider them, it was not clear if Yang-Mills theories were even quantum-mechanically viable. All this changed in the early 1970s when 't Hooft and Veltman showed that these theories were renormalizable, and large amounts of data started to point in the direction of the relevance of non-Abelian gauge theories to physics, leading to the $SU(3)\times SU(2)\times U(1)$ picture of strong, electromagnetic and weak interactions, which culminated in the discoveries of the W^+, W^- and Z_0 intermediate vector bosons (1983) with masses that were generated by the Higgs mechanism, and then finally the Higgs boson itself (2012). What gave the Higgs boson the prominence that it ultimately came to have was the realization that in the electroweak $SU(2) \times U(1)$ theory the Higgs boson not only gives masses to the gauge bosons while maintaining renormalizability, but through its Yukawa couplings to the quarks and leptons of the theory it gives masses to the fermions as well. The Higgs boson is thus responsible not just for the masses of the gauge bosons then but for the masses of all the other fundamental particles as well, causing it to be dubbed the "god particle".

from one or more compound states, probably in the ^3P and ^1S configurations.[1,2]

The position of the hydrogen resonance on the energy scale is in very good agreement with theoretical predictions, which range from 9.6 to 9.8 eV.

Because of the difficulty of the present experiment the author had to seek advice on many aspects of the experiment. He is indebted to A. O. McCoubrey, R. F. C. Vessot, and F. Kaufman for advice on handling of atomic hydrogen; to B. R. McAvoy, J. L. Pack, and J. L. Moruzzi for advice on and loan of high-power microwave equipment; to A. V. Phelps and P. J. Chantry for frequent discussions; and to W. J. Uhlig, J. Kearney, and H. T. Garstka for technical assistance.

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[12]The elastic cross section in both molecular and atomic hydrogen decreases with electron energy; thus the transmitted current vs electron energy under our operating conditions is a steeply rising function. On such a curve it would be very difficult to observe a resonance. Fortunately, the elastic cross section of H_2O increases with energy in the 9- to 10-eV range and thus it is possible to alter the slope of the transmitted current vs electron energy by admixing various amounts of H_2O to H_2.

[13]In a mixture of H_2 and H_2O it is difficult to establish the proper energy scale. In a mixture of H_2 and Ne, the rare gas serves both as a buffer gas for enhanced dissociation and as a calibrating gas.

GLOBAL CONSERVATION LAWS AND MASSLESS PARTICLES*

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In all of the fairly numerous attempts to date to formulate a consistent field theory possessing a broken symmetry, Goldstone's remarkable theorem[1] has played an important role. This theorem, briefly stated, asserts that if there exists a conserved operator Q_i such that

$$[Q_i, A_j(x)] = \sum_k t_{ijk} A_k(x),$$

and if it is possible consistently to take $\sum_k t_{ijk} \times \langle 0|A_k|0\rangle \neq 0$, then $A_j(x)$ has a zero-mass particle in its spectrum. It has more recently been observed that the assumed Lorentz invariance essential to the proof[2] may allow one the hope of avoiding such massless particles through the introduction of vector gauge fields and the consequent breakdown of manifest covariance.[3] This, of course, represents a departure from the assumptions of the theorem, and a limitation on its applicability which in no way reflects on the general validity of the proof.

In this note we shall show, within the framework of a simple soluble field theory, that it is possible consistently to break a symmetry (in the sense that $\sum_k t_{ijk}\langle 0|A_k|0\rangle \neq 0$) without requiring that $A(x)$ excite a zero-mass particle. While this result might suggest a general procedure for the elimination of unwanted massless bosons, it will be seen that this has been accomplished by giving up the global conservation law usually implied by invariance under a local gauge group. The consequent time dependence of the generators Q_i destroys the usual global operator rules of quantum field theory (while leaving the local algebra unchanged), in such a way as to preclude the possibility of applying the Goldstone theorem. It is clear that such a modification of the basic operator relations is a far more drastic step than that taken in the usual broken-symmetry theories in which a degenerate vacuum is the sole symmetry-breaking agent, and the operator algebra possesses the full symmetry. However, since superconductivity appears to display a similar behavior, the possibility of breaking such global conservation laws must not be lightly discarded.

Normally, the time independence of

$$Q_i = \int d^3x\, j_i^{\,0}(\vec{x},t)$$

is asserted to be a consequence of the local conservation law $\partial_\mu j^\mu = 0$. However, the relation

$$\partial_\mu \langle 0 | [j_i^{\,\mu}(x), A_j(x')] | 0\rangle = 0$$

implies that

$$\int d^3x \langle 0 | [j_i^{\,0}(x), A_j(x')] | 0\rangle = \mathrm{const}$$

only if the contributions from spatial infinity vanish. This, of course, is always the case in a fully causal theory whose commutators vanish outside the light cone. If, however, the theory is not manifestly covariant (e.g., radiation-gauge electrodynamics), causality is a requirement which must be imposed with caution. Since Q_i consequently may not be time independent, it will not necessarily generate local gauge transformations upon $A_j(x')$ for $x^0 \neq x'^0$ despite the existence of the differential conservation laws $\partial_\mu j^\mu = 0$.

The phenomenon described here has previously been observed by Zumino[4] in the radiation-gauge formulation of two-dimensional electrodynamics where the usual electric charge cannot be conserved. The same effect is not present in the Lorentz gauge where zero-mass excitations which preserve charge conservation are found to occur. (These correspond to gauge parts rather than physical particles.) We shall, however, allow the possibility of the breakdown of such global conservation laws, and seek solutions of our model consistent only with the differential conservation laws.

We consider, as our example, a theory which was partially solved by Englert and Brout,[5] and bears some resemblance to the classical theory of Higgs.[6] Our starting point is the ordinary electrodynamics of massless spin-zero particles, characterized by the Lagrangian

$$\mathcal{L} = -\tfrac{1}{2}F^{\mu\nu}(\partial_\mu A_\nu - \partial_\nu A_\mu) + \tfrac{1}{4}F^{\mu\nu}F_{\mu\nu}$$
$$+\varphi^\mu \partial_\mu \varphi + \tfrac{1}{2}\varphi^\mu \varphi_\mu + ie_0 \varphi^\mu q \varphi A_\mu,$$

where φ is a two-component Hermitian field, and q is the Pauli matrix σ_2. The broken-symmetry condition

$$ie_0 q \langle 0 | \varphi | 0 \rangle = \eta \equiv \begin{pmatrix}\eta_1\\ \eta_2\end{pmatrix}$$

will be imposed by approximating $ie_0\varphi^\mu q\varphi A_\mu$ in the Lagrangian by $\varphi^\mu \eta A_\mu$. The resulting equations of motion,

$$F^{\mu\nu} = \partial^\mu A^\nu - \partial^\nu A^\mu,$$
$$\partial_\nu F^{\mu\nu} = \varphi^\mu \eta,$$
$$\varphi^\mu = -\partial^\mu \varphi - \eta A^\mu,$$
$$\partial_\mu \varphi^\mu = 0,$$

are essentially those of the Brout-Englert model, and can be solved in either the radiation[7] or Lorentz gauge. The Lorentz-gauge formulation, however, suffers from the fact that the usual canonical quantization is inconsistent with the field equations. (The quantization of A_μ leads to an indefinite metric for one component of φ.) Since we choose to view the theory as being imbedded as a linear approximation in the full theory of electrodynamics, these equations will have significance only in the radiation gauge.

With no loss of generality, we can take $\eta_2 = 0$, and find

$$(-\partial^2 + \eta_1^{\,2})\varphi_1 = 0,$$
$$-\partial^2 \varphi_2 = 0,$$
$$(-\partial^2 + \eta_1^{\,2})A_k^{\,T} = 0,$$

where the superscript T denotes the transverse part. The two degrees of freedom of $A_k^{\,T}$ combine with φ_1 to form the three components of a massive vector field. While one sees by inspection that there is a massless particle in the theory, it is easily seen that it is completely decoupled from the other (massive) excitations, and has nothing to do with the Goldstone theorem.

It is now straightforward to demonstrate the failure of the conservation law of electric charge. If there exists a conserved charge Q, then the relation expressing Q as the generator of rotations in charge space is

$$[Q, \varphi(x)] = e_0 q \varphi(x).$$

Our broken symmetry requirement is then

$$\langle 0 | [Q, \varphi_1(x)] | 0 \rangle = -i\eta$$

or, in terms of the soluble model considered here,

$$\int d^3x' \, \eta_1 \langle 0 | [\varphi_1^0(x'), \varphi_1(x)] | 0 \rangle = -i\eta_1.$$

From the result

$$\langle 0 | \varphi_1^0(x') \varphi_1(x) | 0 \rangle = \partial_0 \Delta^{(+)}(x'-x; \eta_1^2),$$

one is led to the consistency condition

$$\eta_1 \exp[-i\eta_1(x_0' - x_0)] = \eta_1,$$

which is clearly incompatible with a nontrivial η_1. Thus we have a direct demonstration of the failure of Q to perform its usual function as a conserved generator of rotations in charge space. It is well to mention here that this result not only does not contradict, but is actually required by, the field equations, which imply

$$(\partial_0^2 + \eta_1^2) Q = 0.$$

It is also remarkable that if A_μ is given any bare mass, the entire theory becomes manifestly covariant, and Q is consequently conserved. Goldstone's theorem can therefore assert the existence of a massless particle. One indeed finds that in that case φ_1 has only zero-mass excitations.

In summary then, we have established that it may be possible consistently to break a symmetry by requiring that the vacuum expectation value of a field operator be nonvanishing without generating zero-mass particles. If the theory lacks manifest covariance it may happen that what should be the generators of the theory fail to be time-independent, despite the existence of a local conservation law. Thus the absence of massless bosons is a consequence of the inapplicability of Goldstone's theorem rather than a contradiction of it. Preliminary investigations indicate that superconductivity displays an analogous behavior.

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[7]This is an extension of a model considered in more detail in another context by D. G. Boulware and W. Gilbert, Phys. Rev. 126, 1563 (1962).

5 What exactly is the Higgs field?

Given that the existence of the Higgs boson has now been confirmed, we need to ask what exactly the field ϕ represents. There are two possibilities. It is either a q-number elementary field ϕ that appears in the fundamental $SU(3)\times SU(2)\times U(1)$ Lagrangian (to thereby be on an equal footing with the fundamental quarks, leptons and gauge bosons), or it is generated as a dynamical bound state, with the field in a dynamically induced Higgs potential then being the c-number matrix element $\langle\Omega|\bar{\psi}\psi|\Omega\rangle$, a dynamical bilinear fermion condensate. The Mexican Hat potential is thus either part of the fundamental Lagrangian or it is generated by dynamics. If the Higgs field is elementary, then while the potential $V(\phi)=\lambda^2\phi^4/4-\mu^2\phi^2/2$ would be its full quantum-mechanical potential, the discussion given earlier of the minima of the potential would correspond to a c-number tree approximation analysis with the ϕ that appeared there being the c-number $\langle\Omega|\phi|\Omega\rangle$. However, in the fermion condensate case there is no tree approximation, with the theory being given by radiative loop diagrams alone. To see how to generate the Mexican Hat potential in this case we consider the Nambu-Jona-Lasinio (NJL) 4-Fermi model.

5.1 Nambu-Jona-Lasinio Chiral Model as a Mean Field Theory

The NJL model is a chirally-symmetric 4-Fermi model of interacting massless fermions with action

$$I_{\rm NJL}=\int d^4x\left[i\bar{\psi}\gamma^{\mu}\partial_{\mu}\psi-\frac{g}{2}[\bar{\psi}\psi]^2-\frac{g}{2}[\bar{\psi}i\gamma_5\psi]^2\right]. \qquad (10)$$

As such it is a relativistic generalization of the BCS model. In the mean field, Hartree-Fock approximation one introduces a trial wave function parameter m that is not in the original action, and then decomposes the action into two pieces, a mean field piece and a residual interaction according to:

$$\begin{aligned}I_{\rm NJL}&=\int d^4x\left[i\bar{\psi}\gamma^{\mu}\partial_{\mu}\psi-m\bar{\psi}\psi+\frac{m^2}{2g}\right]+\int d^4x\left[-\frac{g}{2}[\bar{\psi}\psi-\frac{m}{g}]^2-\frac{g}{2}[\bar{\psi}i\gamma_5\psi]^2\right]\\&=I_{\rm Mean\ Field}+I_{\rm Residual\ Interaction}\end{aligned} \qquad (11)$$

where $I_{\rm Mean\ Field}$ contains the kinetic energy of a now massive fermion and a self-consistent $m^2/2g$ term. This $m^2/2g$ term acts like a cosmological constant and contributes to the mean field vacuum energy. In the mean field approximation one sets

$$\langle S|[\bar{\psi}\psi-m/g]^2|S\rangle=\langle S|[\bar{\psi}\psi-m/g]|S\rangle^2=0,\quad \langle S|\bar{\psi}\psi|S\rangle=m/g,\quad \langle S|\bar{\psi}i\gamma_5\psi|S\rangle=0. \qquad (12)$$

In this approximation the physical mass M is the value of m that satisfies that satisfies $\langle S|\bar{\psi}\psi|S\rangle = m/g$, and the one loop contribution of the fermionic negative energy Dirac sea to $\langle S|\bar{\psi}\psi|S\rangle$ yields the gap equation

$$-\frac{M\Lambda^2}{4\pi^2} + \frac{M^3}{4\pi^2}\ln\left(\frac{\Lambda^2}{M^2}\right) = \frac{M}{g}, \tag{13}$$

where Λ is an ultraviolet cut-off, as needed since the NJL model is not renormalizable.

Given this gap equation we can calculate the one loop mean field vacuum energy $\epsilon(m)$ as a function of m to obtain

$$\epsilon(m) = i\int \frac{d^4p}{(2\pi)^4}\mathrm{Tr}\ \mathrm{Ln}\left[\frac{\gamma^\mu p_\mu - m + i\epsilon}{\gamma^\nu p_\nu + i\epsilon}\right] - \frac{m^2}{2g} = \frac{m^4}{16\pi^2}\ln\left(\frac{\Lambda^2}{m^2}\right) - \frac{m^2M^2}{8\pi^2}\ln\left(\frac{\Lambda^2}{M^2}\right) + \frac{m^4}{32\pi^2}. \tag{14}$$

We thus see that while the energy $i\int d^4p/(2\pi)^4\mathrm{Tr}\ \mathrm{Ln}[\gamma^\mu p_\mu - m]$ has quartic, quadratic and logarithmically divergent pieces, the subtraction of the massless vacuum energy $i\int d^4p/(2\pi)^4\mathrm{Tr}\ \mathrm{Ln}[\gamma^\mu p_\mu]$ removes the quartic divergence, with the subtraction of the self-consistent induced mean field term $m^2/2g$ then leaving $\epsilon(m)$ only logarithmically divergent. We recognize the resulting logarithmically divergent $\epsilon(m)$ as having a local maximum at $m = 0$, and a global minimum at $m = M$ where M itself is finite. We thus induce none other than a dynamical double-well Mexican Hat potential, and identify M as the matrix element of a fermion bilinear according to $M/g = \langle S|\bar{\psi}\psi|S\rangle$.

If instead of looking at matrix elements in the translationally-invariant vacuum $|S\rangle$ we instead look at matrix elements in coherent states $|C\rangle$ where $m(x) = \langle C|\bar{\psi}(x)\psi(x)|C\rangle$ is now spacetime dependent, we then find [T. Eguchi and H. Sugawara, Phys. Rev. D 10, 4257 (1974); P. D. Mannheim, Phys. Rev. D. 14, 2072 (1976)] that the resulting mean field effective action has a logarithmically divergent part of the form

$$I_{\mathrm{EFF}} = \int d^4x \frac{1}{8\pi^2}\ln\left(\frac{\Lambda^2}{M^2}\right)\left[\frac{1}{2}\partial_\mu m(x)\partial^\mu m(x) + m^2(x)M^2 - \frac{1}{2}m^4(x)\right]. \tag{15}$$

If we introduce a coupling $g_{\mathrm{A}}\bar{\psi}\gamma_\mu\gamma_5 A^\mu_5\psi$ to an axial gauge field $A^\mu_5(x)$, on setting $\phi = \bar{\psi}(1+\gamma_5)\psi$ the effective action becomes

$$I_{\mathrm{EFF}} = \int d^4x \frac{1}{8\pi^2}\ln\left(\frac{\Lambda^2}{M^2}\right)\left[\frac{1}{2}|(\partial_\mu - 2ig_{\mathrm{A}}A_{\mu 5})\phi(x)|^2 + |\phi(x)|^2M^2 - \frac{1}{2}|\phi(x)|^4 - \frac{g_{\mathrm{A}}^2}{6}F_{\mu\nu 5}F^{\mu\nu 5}\right]. \tag{16}$$

We recognize this action as a double-well Ginzburg-Landau type Higgs Lagrangian, only now generated dynamically. We thus generalize to the relativistic chiral case Gorkov's derivation of the Ginzburg-Landau order parameter action starting from the BCS 4-Fermi theory. In the I_{EFF} effective action associated with the NJL model there is a double-well Higgs potential, but since $m(x) = \langle C|\bar{\psi}(x)\psi(x)|C\rangle$ is a c-number, $m(x)$ does not itself represent a q-number scalar field. Rather, as we now show, the q-number fields are to be found as collective modes generated by the residual interaction.

5.2 The Collective Goldstone and Higgs Modes

To find the collective modes we calculate $\Pi_{\rm S}(x) = \langle\Omega|T(\bar{\psi}(x)\psi(x)\bar{\psi}(0)\psi(0))|\Omega\rangle$, $\Pi_{\rm P}(x) = \langle\Omega|T(\bar{\psi}(x)i\gamma_5\psi(x)\bar{\psi}(0)i\gamma_5\psi(0))|\Omega\rangle$ in the scalar and pseudoscalar sectors, as is appropriate to a chiral-invariant theory. If first we take the fermion to be massless (i.e. setting $|\Omega\rangle = |N\rangle$ where $\langle N|\bar{\psi}\psi|N\rangle = 0$) to one loop order in the 4-Fermi residual interaction we obtain

$$\Pi_{\rm S}(q^2) = \Pi_{\rm P}(q^2) = -\frac{1}{8\pi^2}\left(2\Lambda^2 + q^2 \ln\left(\frac{\Lambda^2}{-q^2}\right) + q^2\right). \tag{17}$$

The scattering matrices in the two channels are given by

$$T_{\rm S}(q^2) = \frac{1}{g^{-1} - \Pi_{\rm S}(q^2)}, \quad T_{\rm P}(q^2) = \frac{1}{g^{-1} - \Pi_{\rm P}(q^2)}, \tag{18}$$

and with g^{-1} given by the gap equation, near $q^2 = -2M^2$ both scattering matrices behave as

$$T_{\rm S}(q^2) = T_{\rm P}(q^2) = \frac{Z^{-1}}{(q^2 + 2M^2)}, \qquad Z = \frac{1}{8\pi^2}\ln\left(\frac{\Lambda^2}{M^2}\right), \tag{19}$$

to give degenerate (i.e. chirally symmetric) scalar and pseudoscalar tachyons at $q^2 = -2M^2$ (just like fluctuating around the local maximum in a double-well potential), with $|N\rangle$ thus being unstable.

However, suppose we now take the fermion to have non-zero mass M (i.e. we set $|\Omega\rangle = |S\rangle$). Now we obtain

$$\begin{aligned}
\Pi_{\rm P}(q^2) &= -\frac{\Lambda^2}{4\pi^2} + \frac{M^2}{4\pi^2}\ln\left(\frac{\Lambda^2}{M^2}\right) - \frac{q^2}{8\pi^2}\ln\left(\frac{\Lambda^2}{M^2}\right) - \frac{(q^2 - 4M^2)}{8\pi^2} \\
&- \frac{(8M^4 - 8M^2q^2 + q^4)}{8\pi^2 q^2}\left(\frac{-q^2}{4M^2 - q^2}\right)^{1/2} \ln\left(\frac{(4M^2 - q^2)^{1/2} + (-q^2)^{1/2}}{(4M^2 - q^2)^{1/2} - (-q^2)^{1/2}}\right), \\
\Pi_{\rm S}(q^2) &= -\frac{\Lambda^2}{4\pi^2} + \frac{M^2}{4\pi^2}\ln\left(\frac{\Lambda^2}{M^2}\right) + \frac{(4M^2 - q^2)}{8\pi^2}\ln\left(\frac{\Lambda^2}{M^2}\right) + \frac{(4M^2 - q^2)}{8\pi^2} \\
&- \frac{(4M^2 - q^2)}{8\pi^2}\left(\frac{4M^2 - q^2}{-q^2}\right)^{1/2} \ln\left(\frac{(4M^2 - q^2)^{1/2} + (-q^2)^{1/2}}{(4M^2 - q^2)^{1/2} - (-q^2)^{1/2}}\right).
\end{aligned} \tag{20}$$

Given the form for g^{-1}, we find a dynamical pseudoscalar Goldstone boson bound state at $q^2 = 0$ and a dynamical scalar Higgs boson bound state at $q^2 = 4M^2$ $(= -2 \times M^2(\text{tachyon}))$. The two dynamical bound states are not degenerate in mass (spontaneously broken chiral symmetry), and the dynamical Higgs scalar mass $2M$ is twice the induced mass of the fermion.

6 What Comes Next?

The most crucial question for the Higgs boson is determining its underlying nature. Thus we need to find some experimental way to determine whether the Higgs boson is elementary or dynamical. If the Higgs field is elementary, one would have to treat the Higgs Lagrangian as a bona fide quantum field theory with an a priori double-well potential being present in the fundamental Lagrangian itself. The great appeal of a Higgs potential of the specific form $V(\phi) = \lambda^2\phi^4/4 - \mu^2\phi^2/2$ is that the radiative corrections that it generates are renormalizable. Thus not only is the massive gauge boson sector of the theory renormalizable (because of the Higgs mechanism), the Higgs sector itself would be too. Moreover, with an elementary Higgs field electroweak radiative corrections are not only finite, they are straightforwardly calculable.

However, having an elementary Higgs field also has some disturbing consequences. First we note that the radiative corrections in the Higgs sector itself lead to a quadratically divergent Higgs self energy. While this divergence can be made finite by regularization, there is no control on the ensuing magnitude that the Higgs mass would take, as it would be as large as the regulator masses. These masses could be at the 10^{16} GeV or so grandunified scale (a very large scale if the proton lifetime is to be greater than its current experimental lower bound) or even at the 10^{19} GeV quantum gravitational Planck scale. Such mass values are orders of magnitude larger than the Higgs boson's observed 126 GeV mass. In order to resolve such a disparity (known as the hierarchy problem) it had long been conjectured that there would be a supersymmetry between bosons and fermions, and with fermion and boson loops having opposite signs, a fermionic superparticle could then cancel the quadratic divergence in the Higgs self-energy. For it to do so to the degree needed, the fermionic superparticle would need to be close in mass to the Higgs boson itself. However, data from the very same LHC that was used to find the Higgs boson have sharply curtailed the possibility that there might be any fermionic superparticle in the requisite mass region. Thus the Higgs self-energy problem is open at the present time.

A second concern for an elementary Higgs field is that the group theoretic structure of the $SU(2) \times U(1)$ theory does not at all constrain the Yukawa couplings of the Higgs field to the fundamental fermions. Thus while the Higgs field can give masses to the quarks and leptons, those masses are totally unconstrained, with all the Yukawa coupling constants being free parameters that one has to introduce by hand. While embedding $SU(2) \times U(1)$ in some grandunified theory of the strong, electromagnetic and weak interactions, say, might solve this problem, the issue is open at the present time.

A third concern for an elementary Higgs field is that, as had been noted earlier, the very minimization of the Higgs potential generates an enormous contribution to the cosmological constant. For a one TeV or so Higgs mass breaking scale one would get a cosmological constant that would be of order 10^{60} times larger than the standard Einstein gravitational theory could possibly permit. This problem is very severe, and it also is open at the present time.

A fourth concern for an elementary Higgs field is an aesthetic, in principle one. At the level of the $SU(3)\times SU(2)\times U(1)$ Lagrangian all fermions and gauge bosons are massless and all coupling constants are dimensionless. The only place where there is a fundamental mass scale is in the $-\mu^2\phi^2/2$ term in the Higgs potential. It would be much more natural and elegant if this μ^2 scale were to be generated dynamically in a then scale invariant theory that possesses no fundamental scale at all.

It is thus of interest to note that all of these concerns can be addressed if the Higgs boson is dynamical. Moreover, there already are two working models that we know of in which the symmetry breaking actually is done dynamically, the BCS theory, and the generation of a Goldstone boson pion in QCD. In the BCS theory the symmetry breaking is done by a dynamically generated non-zero fermion bilinear vacuum expectation value $\langle S|\psi\psi|S\rangle$ in a theory that only contains electrons and phonons while not containing any fundamental scalar fields whatsoever. For the pion, we note that the QCD local color $SU(3)$ Lagrangian of quarks and gluons possesses a global chiral flavor symmetry. This chiral symmetry is broken dynamically by QCD dynamics to yield a Goldstone pion. This pion then acquires a mass via the weak interaction since the electroweak $SU(2)\times U(1)$ action breaks the flavor chiral symmetry intrinsically at the level of weak interaction Lagrangian itself, to thereby make the chiral favor symmetry be only an approximate one.

The generation of a Goldstone pion in QCD is particularly of interest since unlike the NJL model, in the QCD case it occurs in a theory that is renormalizable, and in which all coupling constants are dimensionless and all quark and gluon masses are zero at the level of the Lagrangian. In such a theory not only is there no fundamental $-\mu^2\phi^2/2$ term, there is not even any fundamental ϕ at all, and the theory is scale invariant at the level of the Lagrangian. So let us suppose that dynamical symmetry breaking occurs in some renormalizable Yang-Mills theory of interacting massless fermions and gauge bosons. In such a case, if non-zero gauge boson masses are produced by a dynamical Higgs mechanism, renormalizability would not be impaired since bound state production in a renormalizable theory does not violate renormalizability. Without needing to specify any particular such Yang-Mills theory (i.e. without needing to specify the group structure or the specific matter content of the theory), we note that if the symmetry breaking occurs at all, then since all such theories have no divergences higher than logarithmic, there will be no quadratic divergence associated with any dynamical scalar bound states that might be produced. Thus with a dynamical Higgs boson there is no Higgs self-energy hierarchy problem at all.

As regards the couplings of fermions to the Higgs boson, these couplings are given as the residues of the dynamical poles in the requisite channels. Thus they are determined by the theory itself and are not free parameters at all. To see how things work, let us consider the NJL model as a stand-in for a renormalizable field theory. In its $T_{\rm S}$ and $T_{\rm P}$ channels there are scalar and pseudoscalar bound states, and near the respective poles the scattering amplitudes behave as:

$$T_{\rm S}(q^2)=\frac{Z_{\rm S}^{-1}}{(q^2-4M^2)},\qquad T_{\rm P}(q^2)=\frac{Z_{\rm P}^{-1}}{q^2},\qquad Z_{\rm S}=Z_{\rm P}=\frac{1}{8\pi^2}\ln\left(\frac{\Lambda^2}{M^2}\right). \tag{21}$$

We thus identify the Yukawa coupling of the scalar and psuedoscalar bound states to a fermion anti-fermion pair to be $Z_{\rm S}^{-1/2}$ and $Z_{\rm P}^{-1/2}$, with the equality of $Z_{\rm S}$ and $Z_{\rm P}$ that is found reflecting the underlying chiral symmetry of the NJL theory.

As regards the cosmological constant problem, we had noted above that the self-consistent mean-field treatment of the NJL model automatically generated an $m^2/2g$ term in the mean-field action, with this term serving to cancel the quadratic divergence in the vacuum energy. Unlike the situation with a fundamental Higgs field where there is no control of the vacuum energy, with dynamical symmetry breaking we see that contributions to the vacuum energy are under control. Breaking symmetries dynamically thus provides a good starting point to address the cosmological constant problem. A discussion of how things works in the renormalizable situation may be found in P. D. Mannheim, Found. Phys. 42, 388 (2012).

However, if the Higgs boson is to be dynamical, we would then have to be able to reproduce those aspects of the radiative correction structure associated with a fundamental Higgs field that have been tested. To this end we note that following Gaussian integration on a dummy scalar field variable σ the path integral representation of the generator $Z(\bar{\eta},\eta)$ of fermion Green's functions in the NJL theory can be written as:

$$\begin{aligned} Z(\bar{\eta},\eta) &= \int[d\bar{\eta}d\eta]\exp\left[i\int d^4x\left(i\bar{\psi}\gamma^\mu\partial_\mu\psi-\frac{g}{2}(\bar{\psi}\psi)^2+\bar{\eta}\psi+\bar{\psi}\eta\right)\right] \\ &= \int[d\bar{\eta}d\eta d\sigma]\exp\left[i\int d^4x\left(i\bar{\psi}\gamma^\mu\partial_\mu\psi-\frac{g}{2}(\bar{\psi}\psi)^2+\frac{g}{2}\left(\frac{\sigma}{g}-\bar{\psi}\psi\right)^2+\bar{\eta}\psi+\bar{\psi}\eta\right)\right] \\ &= \int[d\bar{\eta}d\eta d\sigma]\exp\left[i\int d^4x\left(i\bar{\psi}\gamma^\mu\partial_\mu\psi-\sigma\bar{\psi}\psi+\frac{\sigma^2}{2g}+\bar{\eta}\psi+\bar{\psi}\eta\right)\right]. \end{aligned} \tag{22}$$

As we see, the fermion Green's functions of the NJL theory are given as the fermion Green's functions of a scalar field theory whose action is precisely the NJL mean field action. In consequence, the perturbative expansions in the two theories are in one to one correspondence. However, in this scalar field theory there is no source term $J\sigma$ for the scalar field (in a true fundamental Higgs Lagrangian there would be such a source term), and thus the scalar field theory only generates Green's functions with external fermion legs and does not generate any Green's functions with external scalar field legs. Thus in the dynamical Higgs case one can generate the fermion Green's functions using a fundamental Higgs field theory in which the fundamental Higgs field only role is to contribute internally in Feynman diagrams and never to appear in any external legs. In such a case, the all-order iteration of internal σ exchange diagrams then generates the dynamical Higgs and Goldstone poles in $T_{\rm S}(q^2)$ and $T_{\rm P}(q^2)$.

7 The Moral of the Story

With the vacuum of quantum field theory being a dynamical one, in a sense Einstein's ether has reemerged. Only it has reemerged not as the mechanical ether of classical physics that was excluded by the Michaelson-Morley experiment, but as a dynamical, quantum-field-theoretic one full of Dirac's negative energy particles, whose dynamics can spontaneously break symmetries. The type of physics that would be taking place in this vacuum depends on how symmetries are broken, i.e. on whether the breaking is by elementary Higgs fields or by dynamical ones. If the symmetry is broken by a fundamental Higgs field, then the Higgs boson gives mass to fundamental gauge bosons and fermions alike. However, if the breaking is done dynamically, then it is the structure of an ordered vacuum itself that generates masses, with the mass generation mechanism in turn then producing the Higgs boson. In the dynamical case then mass produces Higgs rather than Higgs produces mass. In the dynamical case we should not be thinking of the Higgs boson as being the "god particle". Rather, if anything, we should be thinking of the vacuum as being the "god vacuum".