
FIZIKAI NOBEL-DÍJ, 2013


A tömeg
eredete
és a
Higgs-
mező
avagy a 2013. évi
fizikai Nobel-díj
ELTE TTK
Fizikai Intézet
Az atomoktól a csillagokig
Dávid Gyula
2012. 09. 13.


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2012. július 4.

Világszenzáció!

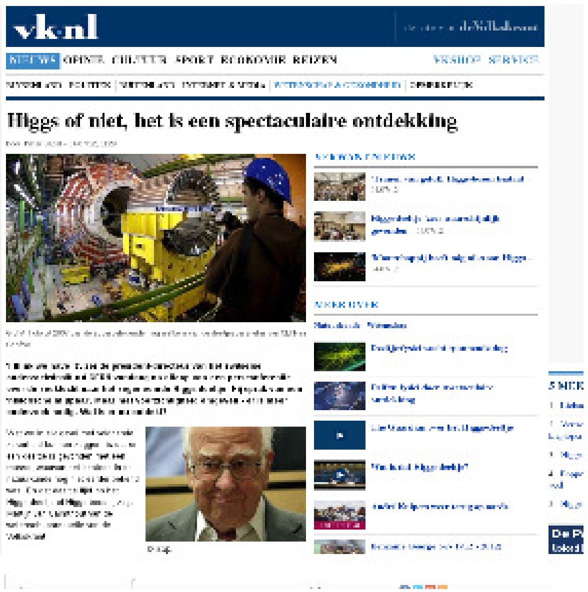
2012. július 4.

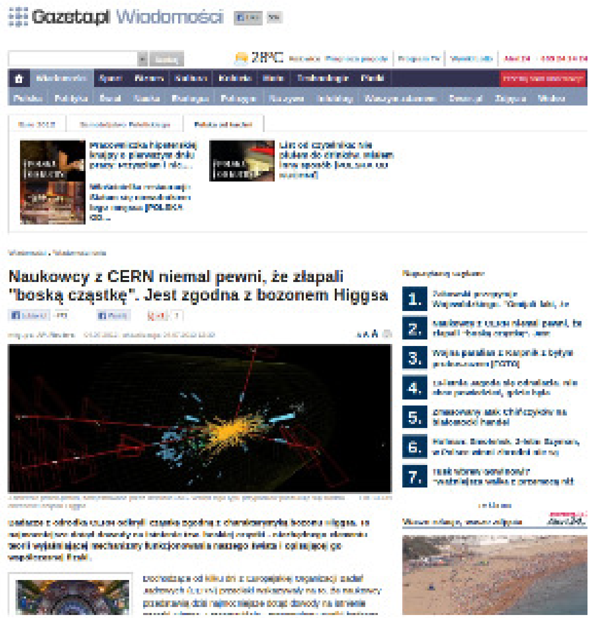
Felfedezték a Higgs-bozont!

Világszenzáció!


Tribune de Genève
Une nouvelle particule a été découverte


NÉPSZABADSÁG ONLINE
Úgy tűnik, megvan az 'isteni' részecske


vk.nl
Higgs of niet, het is een spectaculaire ontdekking


Gazeta.pl Wiadomości
Naukowcy z CERN niemal pewni, że złapali "boską cząstkę". Jest zgodna z bozonem Higgsa

2012. július 4.

Felfedezték a Higgs-bozont!

Világszenzáció!

Az eredeti bejelentés legrövidebb formája

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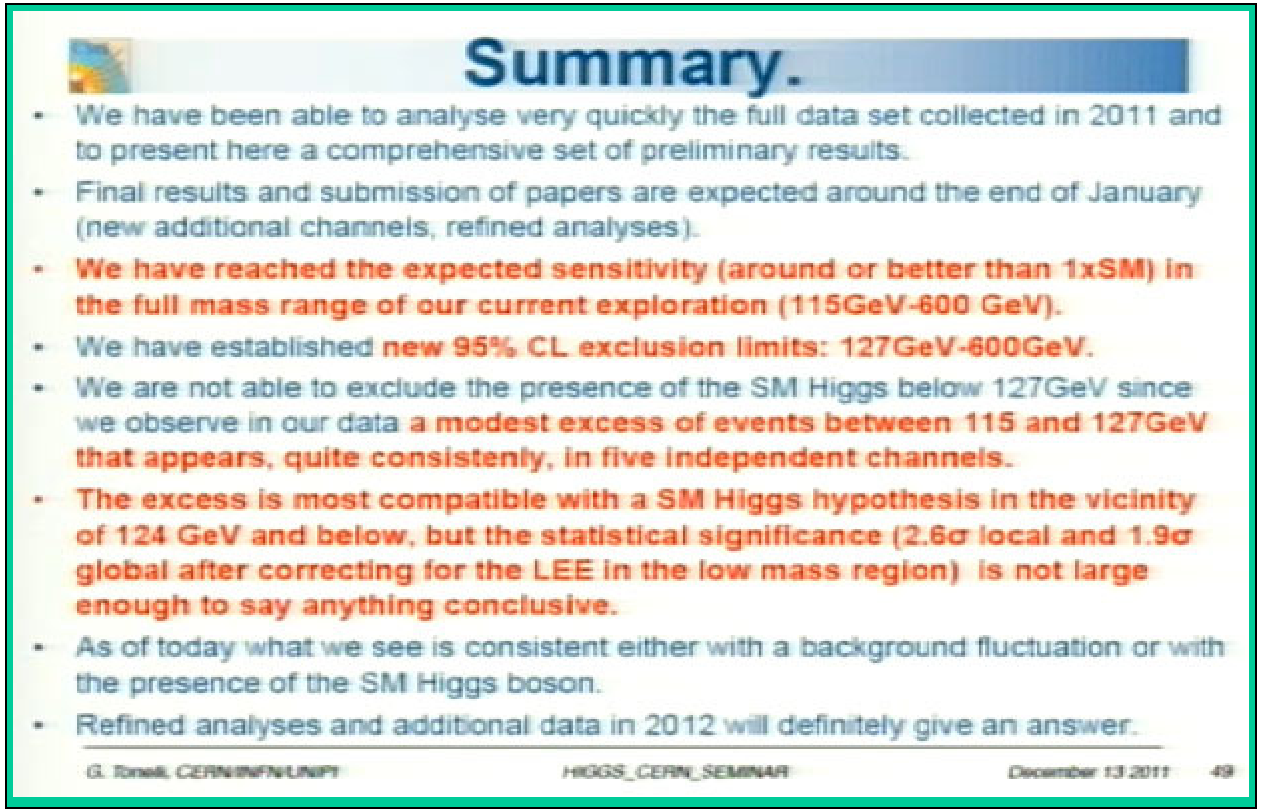

Az eredeti bejelentés kissé hosszabb formája

2012. szeptember 12.

megjelent a tudományos cikk a felfedezésről

Physics Letters B

Volume 716, Issue 1, 17 September 2012, Pages 30–61

Observation of a new boson at a mass of 125 GeV with the CMS experiment at the LHC ☆

Universally Available

This paper is dedicated to the memory of our colleagues who worked on CMS but have since passed away. In recognition of their many contributions to the achievement of this observation.

CMS Collaboration*

CERN, Switzerland

fejléc

(utána jön a több mint 2000 szerző felsorolása)

2012. szeptember 12.

megjelent a tudományos cikk a felfedezésről

http://dx.doi.org/10.1016/j.physletb.2012.08.021, How to Cite or Link Using DOI
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Abstract

Results are presented from searches for the standard model Higgs boson in proton–proton collisions at $\sqrt{s} = 7$ and 8 TeV in the Compact Muon Solenoid experiment at the LHC, using data samples corresponding to integrated luminosities of up to 5.1 fb^{-1} at 7 TeV and 5.3 fb^{-1} at 8 TeV. The search is performed in five decay modes: $\gamma\gamma$, ZZ, W^+W^-, $\tau^+\tau^-$, and $b\bar{b}$. An excess of events is observed above the expected background, with a local significance of 5.0 standard deviations, at a mass near 125 GeV, signalling the production of a new particle. The expected significance for a standard model Higgs boson of that mass is 5.8 standard deviations. The excess is most significant in the two decay modes with the best mass resolution, $\gamma\gamma$ and ZZ; a fit to these signals gives a mass of 125.3±0.4(stat.)±0.5(syst.) GeV. The decay to two photons indicates that the new particle is a boson with spin different from one.

Keywords

CMS; Physics; Higgs

rövid kivonat

2012. szeptember 12.

megjelent a tudományos cikk a felfedezésről

8. Conclusions

Results are presented from searches for the standard model Higgs boson in proton–proton collisions at $\sqrt{s}=7$ and 8 TeV in the CMS experiment at the LHC, using data samples corresponding to integrated luminosities of up to 5.1 fb^{-1} at 7 TeV and 5.3 fb^{-1} at 8 TeV. The search is performed in five decay modes: $\gamma\gamma$, ZZ, W^+W^-, $\tau^+\tau^-$, and $b\bar{b}$. An excess of events is observed above the expected background, with a local significance of 5.0σ, at a mass near 125 GeV, signalling the production of a new particle. The expected local significance for a standard model Higgs boson of that mass is 5.8σ. The global p-value in the search range of 115–130 (110–145) GeV corresponds to 4.6σ (4.5σ). The excess is most significant in the two decay modes with the best mass resolution, $\gamma\gamma$ and ZZ, and a fit to these signals gives a mass of 125.3±0.4(stat.)±0.5(syst.) GeV. The decay to two photons indicates that the new particle is a boson with spin different from one. The results presented here are consistent, within uncertainties, with expectations for the standard model Higgs boson. The collection of further data will enable a more rigorous test of this conclusion and an investigation of whether the properties of the new particle imply physics beyond the standard model.

következtetések

2012. július 4. szenzációja:

(majdnem biztos, hogy) felfedezték a Higgs-bozont!
felfedezték a Higgs-részecskét!
felfedezték a Higgs-mezőt!

Leggyakoribb szövegek:

- ez a részecskefizika Standard modelljének utolsó hiányzó építőkockája
- a Higgs-részecske speciális szerepet játszik a részecskék világában, mert ő adja a tömeget a többi részecskének

Egyéb gyakori szövegek:

- a Higgs-részecske tömege 125 GeV (a proton tömegének 125-szerese)
- a Higgs-részecske nagyon rövid ideig létezik, nem is látták, csak a közvetett nyomokból következtettek a létezésére
- még nem biztos a dolog, további mérésekre van szükség
- lehet, hogy nem is egy, hanem több részecskét találtak, vajon csak az egyik a Higgs, vagy mindegyik az?
- Peter Higgs már több mint 50 éve megjósolta a részecske létezését

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Felmerülő kérdések:

- mi az a Standard modell, és miért játszik benne olyan fontos szerepet a Higgs-bozon?
- mi a viszony a Higgs-mező, a Higgs-részecske és a Higgs-bozon között?
- ha ez a részecske mindenütt ott van (hiszen ő adja a tömeget minden más részecskének), miért kellett 50 évig várni a felfedezésére?
- ha csak nagyon rövid ideig létezik, hogy lehet ott mindig és mindenütt (hogy ő adhassa a tömeget minden más részecskének)?
- ha ő adja a tömeget minden más részecskének, akkor neki honnan van a 125 GeV-nyi tömege? Adott magának is?
- miért kell tömeget ADNI a részecskének? Nincs nekik?
- egyáltalán: hogyan lehet tömeget ADNI a részecskének?
- tulajdonképpen mi a csuda az a TÖMEG?

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- mi az a Standard modell, és miért játszik benne olyan font… szerepet a Higgs-bozon?
- mi a viszony a Higgs-mező, a Higgs-része… között?
- ha ez a részecske mindenütt ott v… …get minden más részecskének), mi… …fedezésére?
- ha csak nagyon r… …et ott mindig és mindenütt (hogy ő a… …más részecskének)?
- ha ő… …más részecskének, akkor neki honnan …tömege? Adott magának is?
- miér… …get ADNI a részecskének? Nincs nekik?
- egyált…an: hogyan lehet tömeget ADNI a részecskének?
- tulajdonképpen mi a csuda az a TÖMEG?

Ezekről az alapkérdésekről szól ez az előadás

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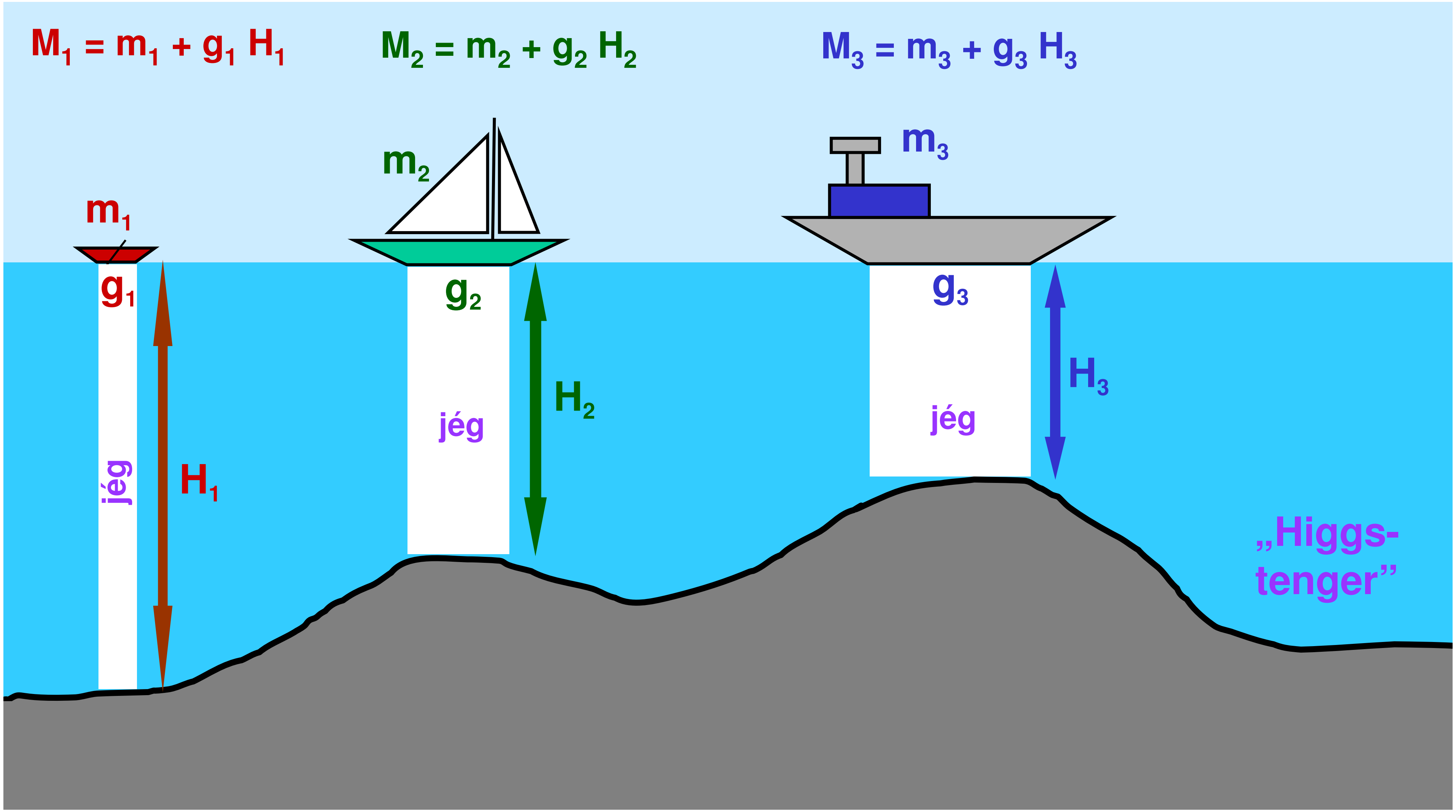
- mi a viszony a Higgs-mező, a Higgs-része... között?
- ha ez a részecske mindenütt ott v... den más részecskének), mi...
- ha csak nagyon r... és mindenütt (hogy ő a... csékének)?
- ha ő... énnek, akkor neki honnan ... magának is?
- miér... szecskének? Nincs nekik?
- egyált... ...enet tömeget ADNI a részecskének?
- tulajdo... ...ppen mi a csuda az a TÖMEG?

Ezeknek az alapkérdéseknek
egy parányi részéről
szólt a tavaly szeptemberi előadás

$M = m + g\,H$

Értsük meg a „tömegadást” matek nélkül!



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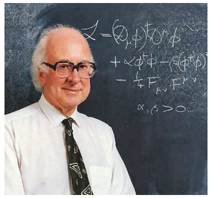
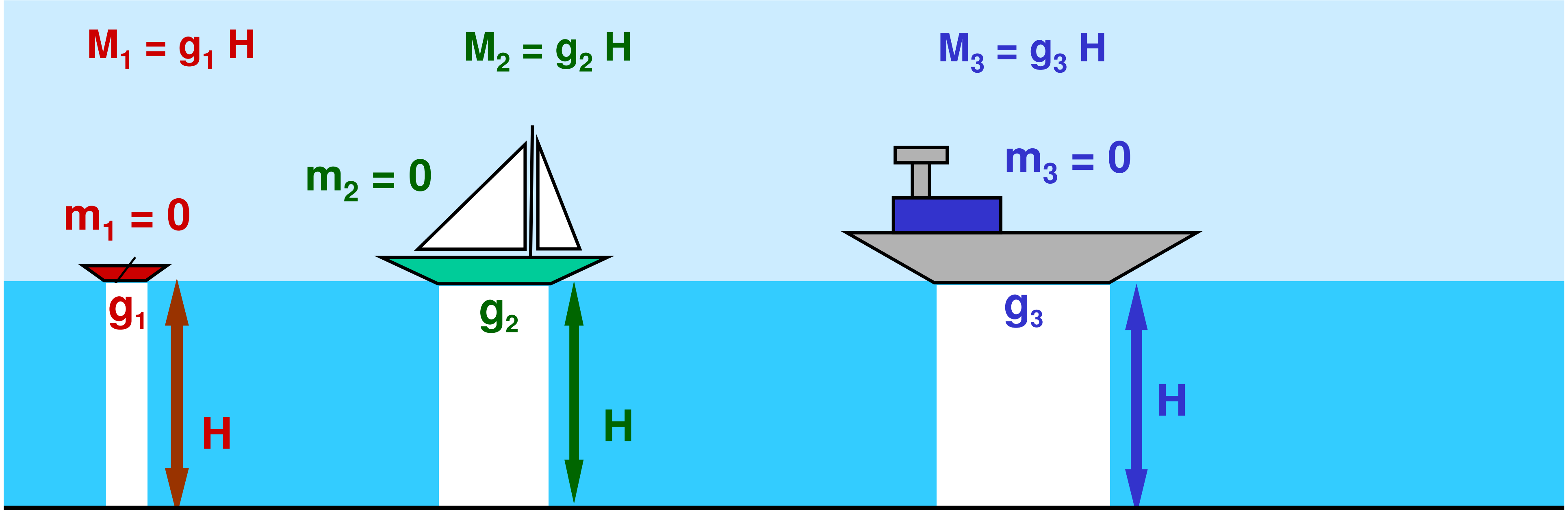


Peter Higgs módosításai



1/ a víz mindenhol egyforma mélységű → mozgás közben nem változik a tömeg

2/ a hajók eredeti tömege 0

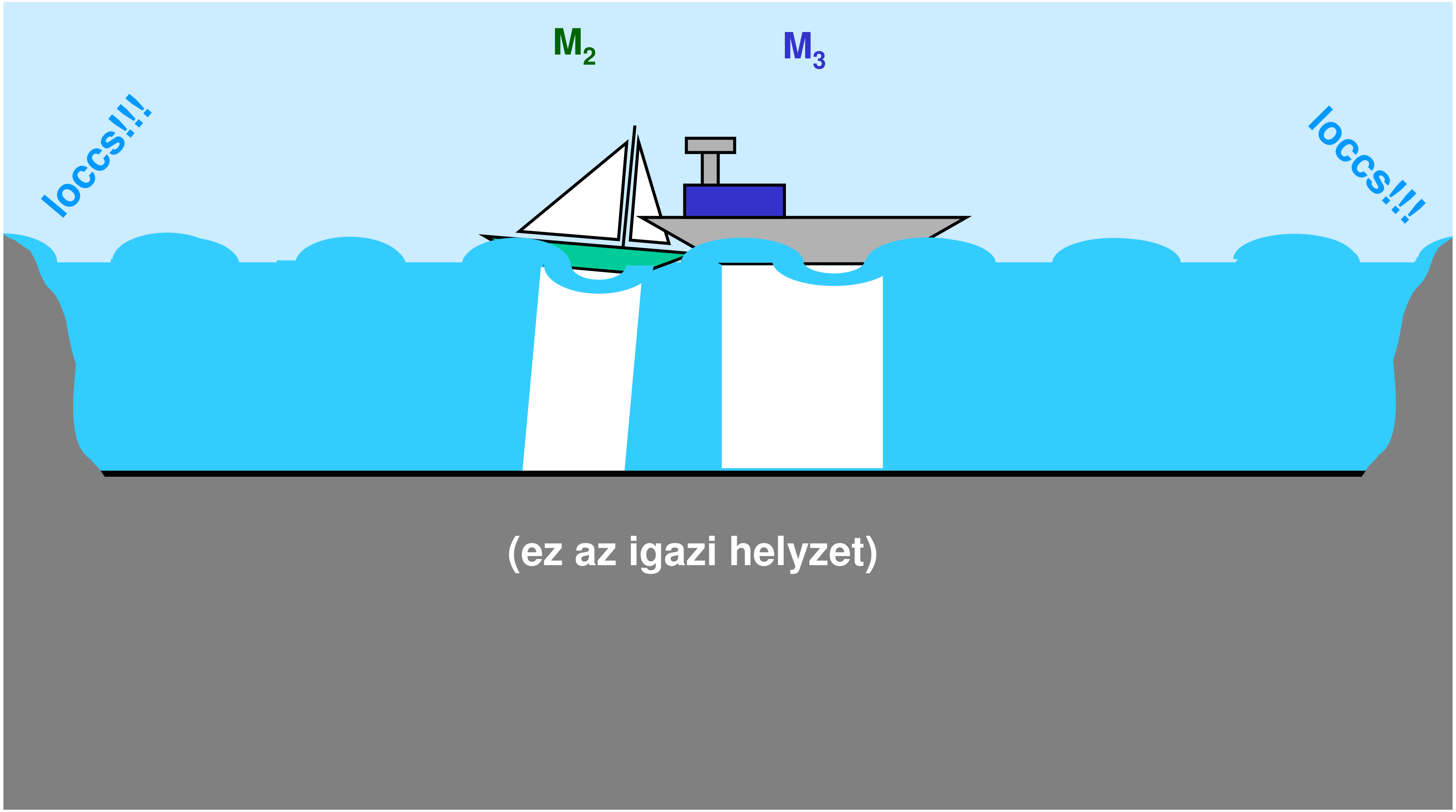
működik a QFT matekja!

Emlékeztető: A kvantum-mezőelmélet (QFT) matematikája a nulla nyugalmi tömegű részecskéket (és csak azokat) tudja jól leírni.

☺

Hogy lehetne kimutatni a láthatatlan „Higgs-tenger” létezését?

Ütköztessünk két hajót nagy energiával!



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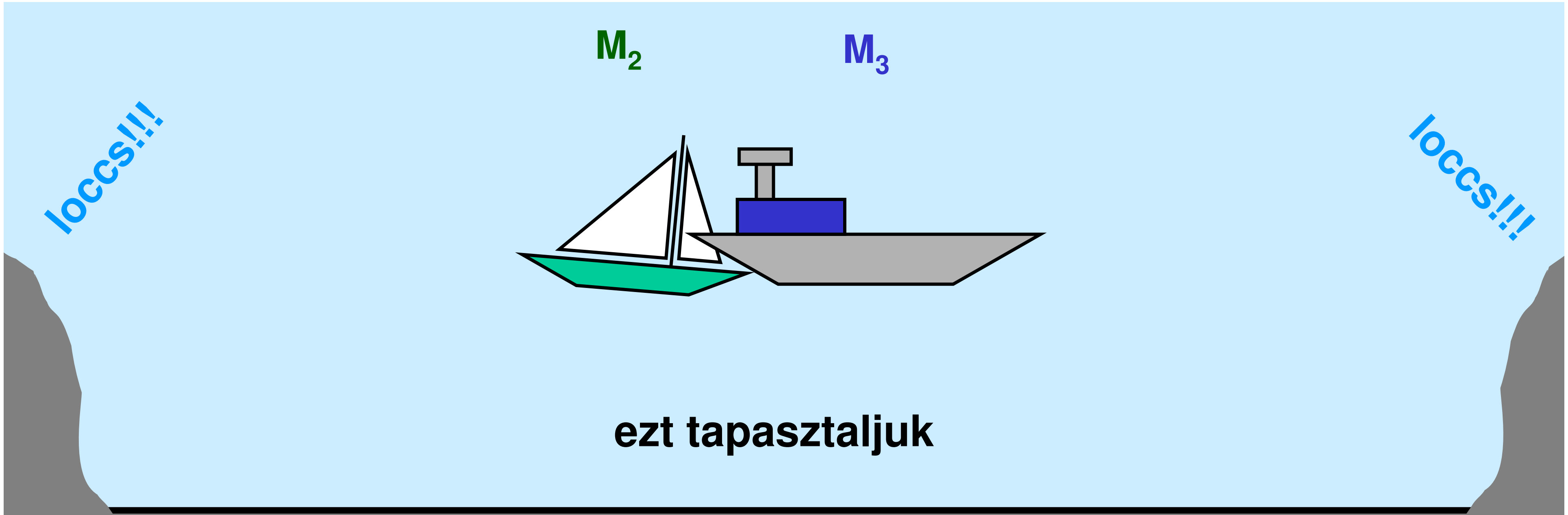


Szótár: a tenger vize - a Higgs mező

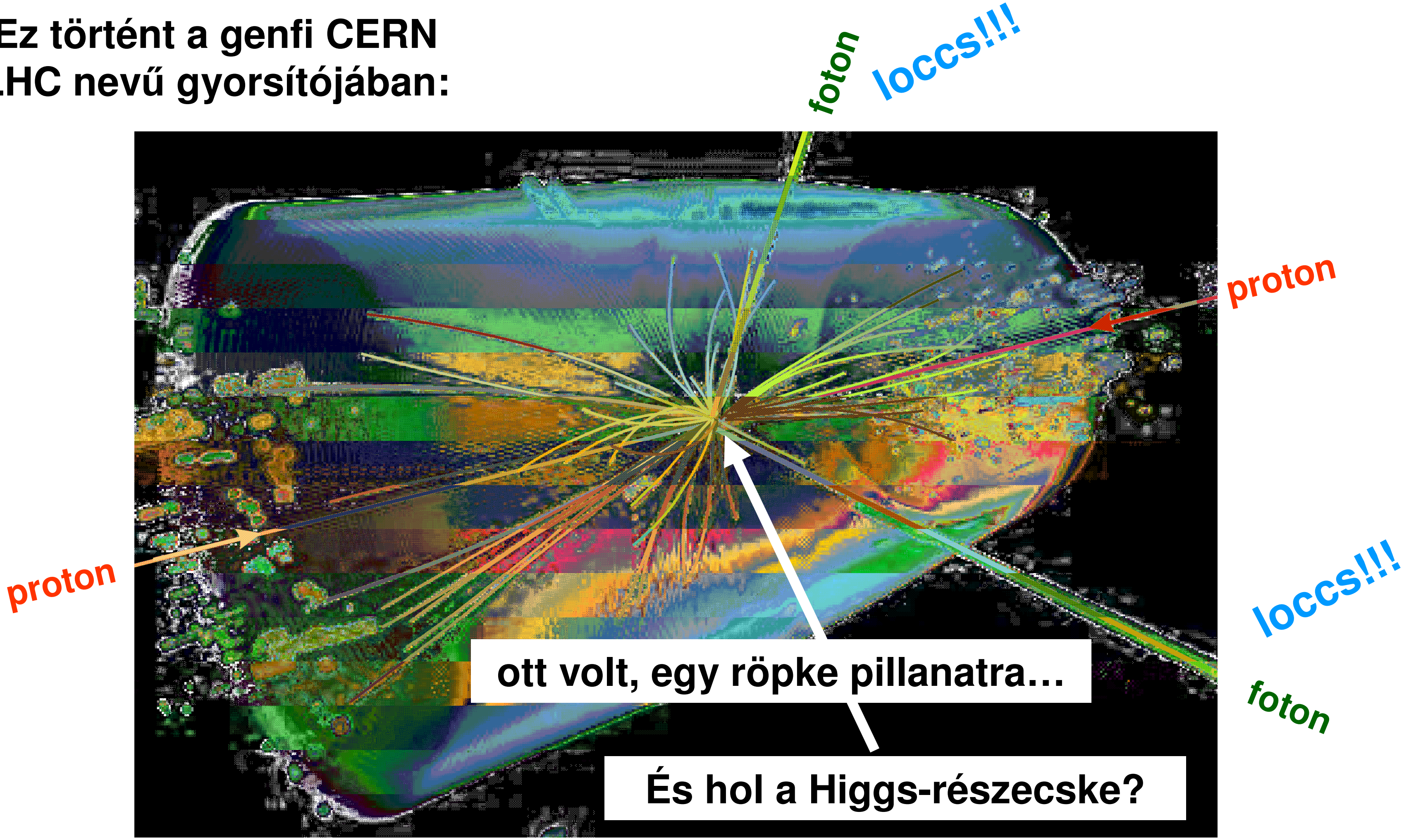
a tenger fodrozódása - a Higgs részecske

a partra csapó hullámok - a detektált másodlagos részecskék

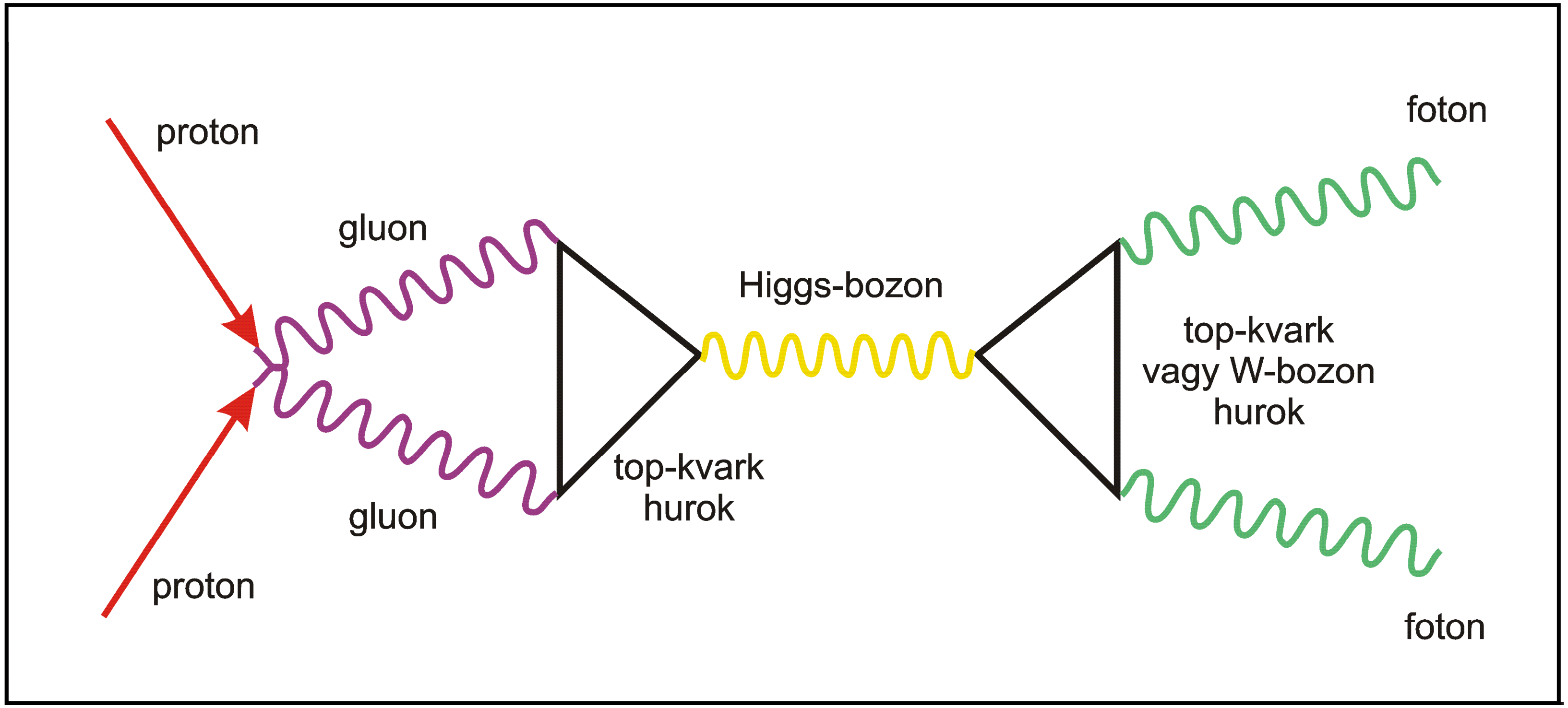
Hogy lehetne kimutatni a láthatatlan „Higgs-mező” létezését?

Ütköztessünk két protont nagy energiával!

Ez történt a genfi CERN LHC nevű gyorsítójában:

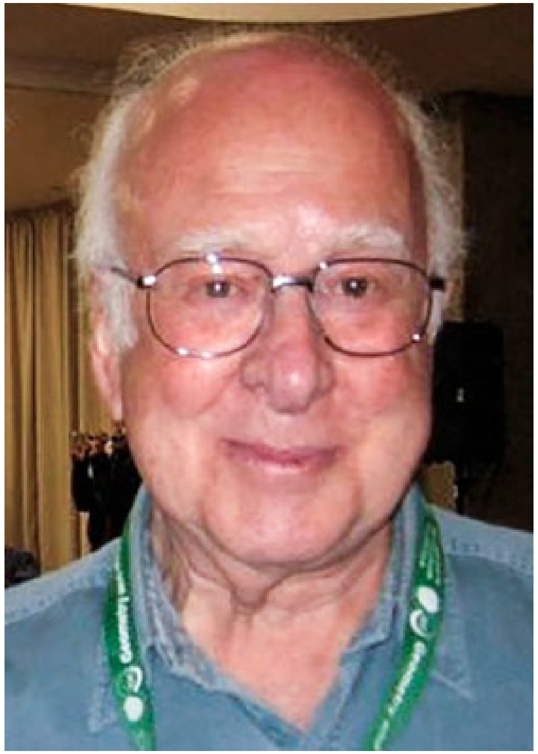


ugyanez, részecskefizikus titkosírással, a Feynman-diagramok nyelvén

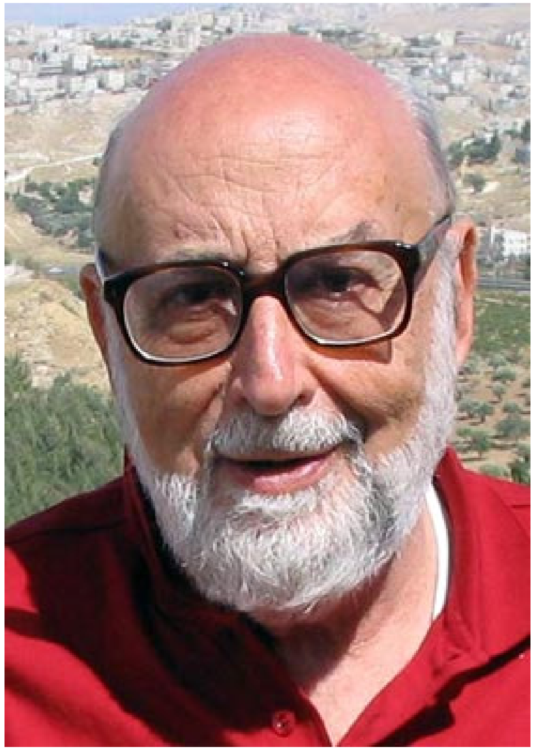


2013. október 8.

Fizikai Nobel-díj:



Peter Higgs (1929 -)



François Englert (1932 -)

„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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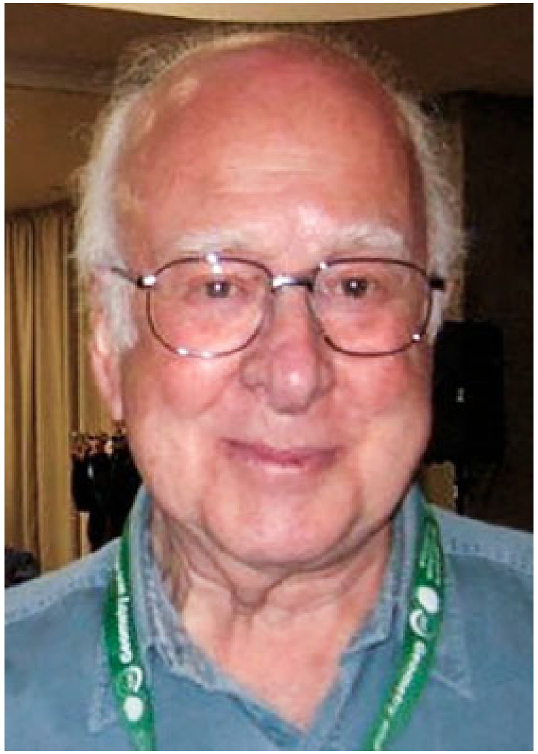

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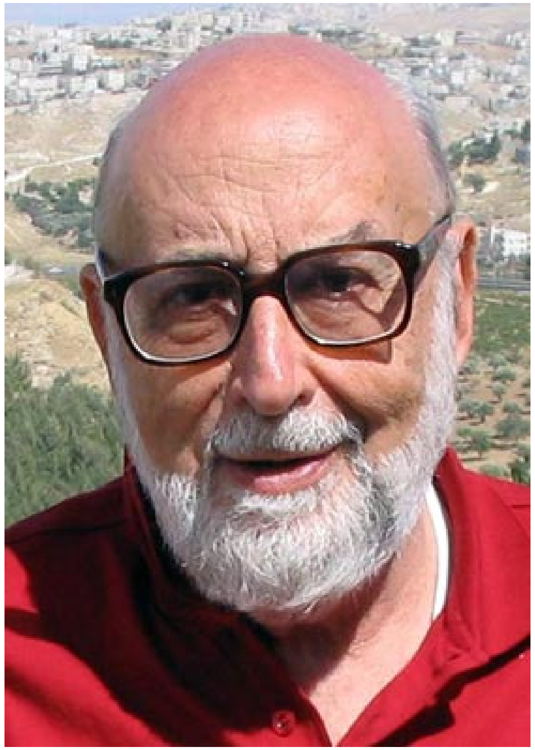

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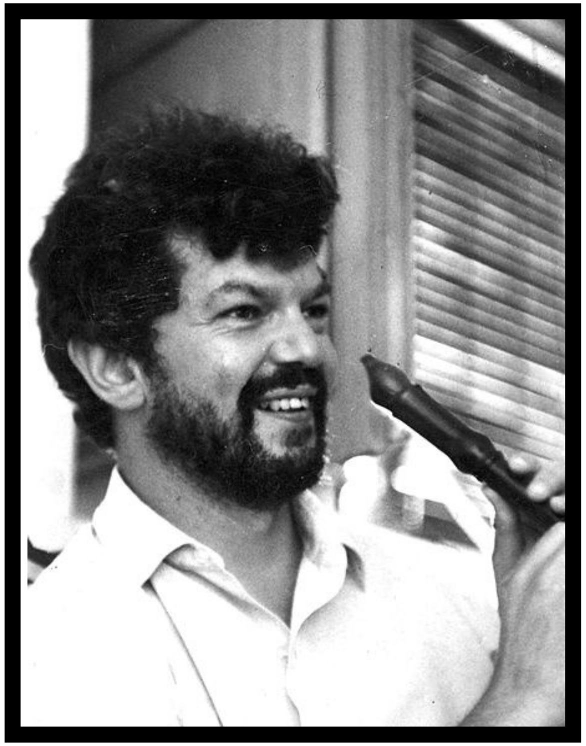

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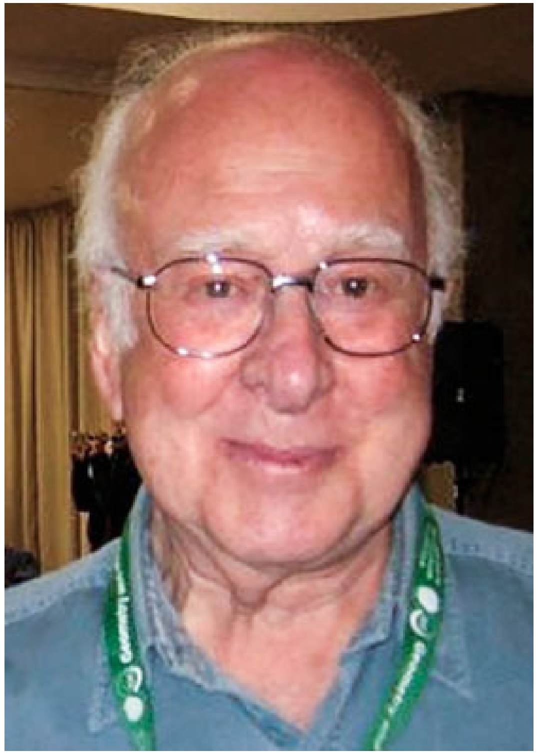

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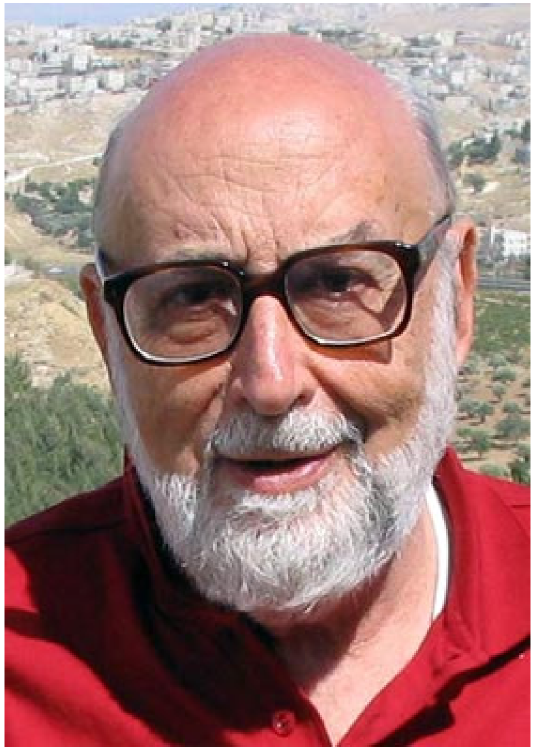

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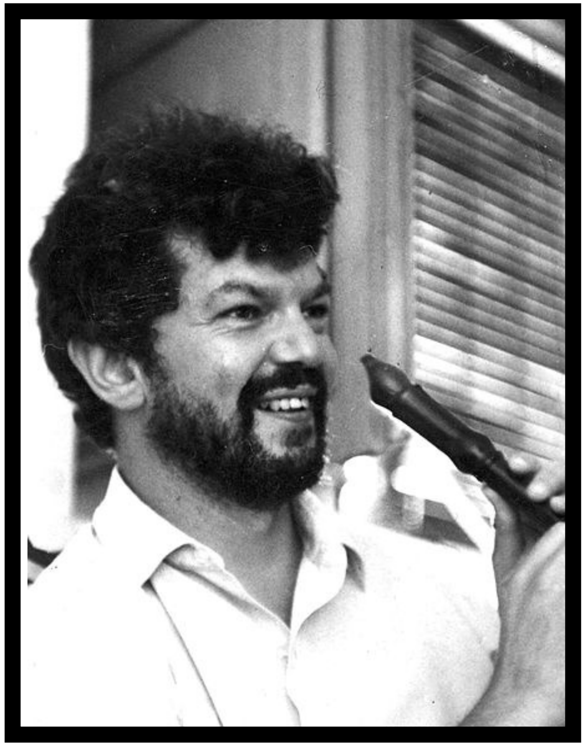

Robert Brout (1928 - 2011)

„egy olyan mechanizmus elméleti felfedezéséért, amely hozzájárult a szubatomi részecskék tömege eredetének megértéséhez, és amelyet igazolt az elmélet által megjósolt alapvető részecskének az ATLAS és a CMS kísérletek általi felfedezése a CERN Nagy Hadronütköztetőjében"

Négy alapvető cikk 1964-ből

BROKEN SYMMETRIES, MASSLESS PARTICLES AND GAUGE FIELDS

P. W. HIGGS
Tait Institute of Mathematical Physics, University of Edinburgh, Scotland

Received 27 July 1964

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

Volume 12, number 2 PHYSICS LETTERS 15 September 1964

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

$$\varphi_1 \to \varphi_1 \cos\alpha + \varphi_2 \sin\alpha\,,$$
$$\varphi_2 \to -\varphi_1 \sin\alpha + \varphi_2 \cos\alpha\,. \qquad (1)$$

Then there is a conserved current j_μ such that

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

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 $i\langle[j_\mu(x), \varphi_1(y)]\rangle$ 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 non-relativistic case by showing that the most general form of this Fourier transform is now, in Gilbert's notation,

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

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

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

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

$$j^\mu = \partial_\nu F'^{\mu\nu}\,,$$
$$F_{\mu\nu}' = \partial_\mu A_\nu' - \partial_\nu A_\mu'\,. \qquad (5)$$

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 $i\langle[A_\mu'(x), \varphi_1(y)]\rangle$ must be given by eq. (3). Applying eq. (5) to this commutator gives us as the Fourier transform of $i\langle[j_\mu(x), \varphi_1(y)]\rangle$ 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.

References

1) J.Goldstone, Nuovo Cimento 19 (1961) 154.
2) J.Goldstone, A.Salam and S.Weinberg, Phys.Rev. 127 (1962) 965.
3) A.Klein and B.W.Lee, Phys.Rev.Letters 12 (1964) 266.
4) W.Gilbert, Phys.Rev.Letters 12 (1964) 713.
5) J.Schwinger, Phys.Rev.127 (1962) 324.

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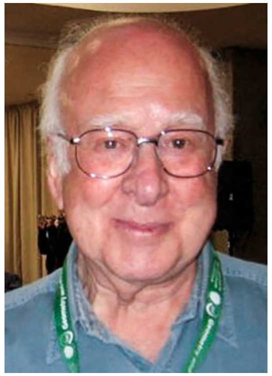

Peter Higgs
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Nem a méret számít… Négy alapvető cikk 1964-ből

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

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

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

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

$$j^\mu = \partial_\nu F'^{\mu\nu}\,,$$

$$F_{\mu\nu}' = \partial_\mu A_\nu' - \partial_\nu A_\mu'\,. \tag{5}$$

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 $i\langle[A_\mu'(x), \varphi_1(y)]\rangle$ must be given by eq. (3). Applying eq. (5) to this commutator gives us as the Fourier transform of $i\langle[j_\mu(x), \varphi_1(y)]\rangle$ 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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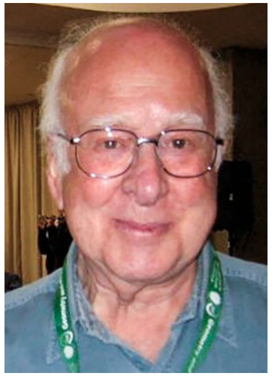
1) J. Goldstone, Nuovo Cimento 19 (1961) 154.
2) J. Goldstone, A. Salam and S. Weinberg, Phys. Rev. 127 (1962) 965.
3) A. Klein and B. W. Lee, Phys. Rev. Letters 12 (1964) 266.
4) W. Gilbert, Phys. Rev. Letters 12 (1964) 713.
5) J. Schwinger, Phys. Rev. 127 (1962) 324.

* * * * *

„…kiűztük a boszorkányokat az elméletből!”



Peter Higgs (1929 -)

Nem a méret számít… **Négy alapvető cikk 1964-ből**

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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

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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

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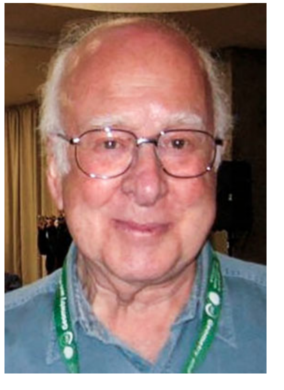

Peter Higgs
(1929 -)

Nem a méret számít… Négy alapvető cikk 1964-ből

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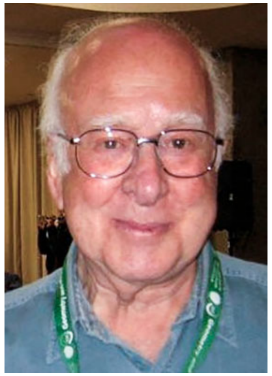

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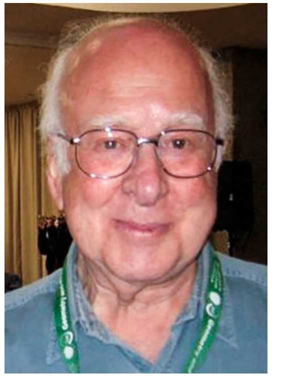

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Nem a méret számít… Négy alapvető cikk 1964-ből

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 loc phase and local γ_5-phase transformations. In this model the gauge fields themselves may br the γ_5 invariance leading to a mass for the orig inal 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 transformatio of a charged boson; we then present the genera ization to an arbitrary compact Lie group.

The interaction between the φ and the A_μ fields is

$$H_{\rm int} = ieA_\mu \varphi^* \overleftrightarrow{\partial}_\mu \varphi - e^2 \varphi^* \varphi A_\mu A_\mu,$$

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

3

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]

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

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 $\sum_{A,B'} T_{a,AB'}\varphi_A$

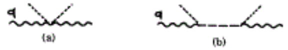

FIG. 1. Broken-symmetry diagram leading to a mass for the gauge field. Short-dashed line, $\langle\varphi_1\rangle$; long-dashed line, φ_2 propagator; wavy line, A_μ propagator. (a) $\to (2\pi)^4 ie^2 g_{\mu\nu}\langle\varphi_1\rangle^2$, (b) $\to -(2\pi)^4 ie^2 (q_\mu q_\nu/q^2) \times \langle\varphi_1\rangle^2$.

$\times\langle\varphi_{B'}\rangle$ is, in the lowest order,

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

With λ the coupling constant of the Yang-M field, the same calculation as before yield

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

giving a value for the mass

$$\mu_a{}^2 = -(\langle\varphi\rangle T_a T_a \langle\varphi\rangle).$$

(2) Consider the interaction Hamiltonian

$$H_{\rm int} = -\eta\bar\psi\gamma_\mu\gamma_5\psi B_\mu - \epsilon\bar\psi\gamma_\mu\psi A_\mu,$$

where A_μ and B_μ are vector and pseudovec gauge fields. The vector field causes attra whereas the pseudovector leads to repulsio tween particle and antiparticle. For a suit choice of ϵ and η there exists, as in Johnso model,[11] a broken-symmetry solution corre ing to an arbitrary mass m for the ψ field f the scale of the problem. Thus the fermion propagator $S(p)$ is

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

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 using Johnson's method[12]:

$$J_\mu{}^5 = -\eta \lim_{\xi\to 0} \bar\psi'(x+\xi)\gamma_\mu\gamma_5\psi'(x),$$
$$\psi'(x) = \exp\left[-i\int_{-\infty}^x \eta B_\mu(y)dy^\mu\gamma_5\right]\psi(x).$$

This gives for the polarization tensor of the

pseudovector field

$$\Pi_{\mu\nu}{}^5(q) = \eta^2\frac{i}{(2\pi)^4}\int {\rm Tr}\{S(p-\tfrac12 q)\Gamma_{\nu5}(p-\tfrac12 q; p+\tfrac12 q) \times S(p+\tfrac12 q)\gamma_\mu\gamma_5 - S(p)[\partial S^{-1}(p)/\partial p_\nu]S(p)\gamma_\mu\} d^4p, \quad (10)$$

where the vertex function $\Gamma_{\nu5} = \gamma_\nu\gamma_5 + \Lambda_{\nu5}$ satisfies the Ward identity[5]

$$q_\nu\Lambda_{\nu5}(p-\tfrac12 q; p+\tfrac12 q) = \Sigma(p-\tfrac12 q)\gamma_5 + \gamma_5\Sigma(p+\tfrac12 q), \quad (11)$$

which for low q reads

$$q_\nu\Gamma_{\nu5} = q_\nu\gamma_\nu\gamma_5[1-\Sigma_2] + 2\Sigma_1\gamma_5 - 2(q_\nu p_\nu)(\gamma_\lambda p_\lambda)(\partial\Sigma_2/\partial p^2)\gamma_5. \quad (12)$$

The singularity in the longitudinal $\Gamma_{\nu5}$ 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 $\Gamma_{\nu5}$ 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 $q\to 0$.

*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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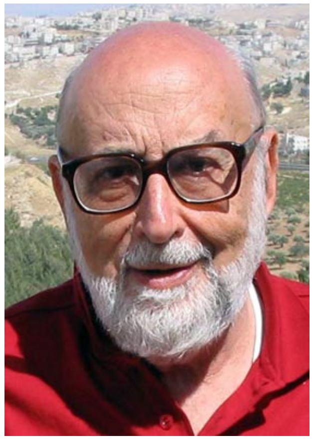
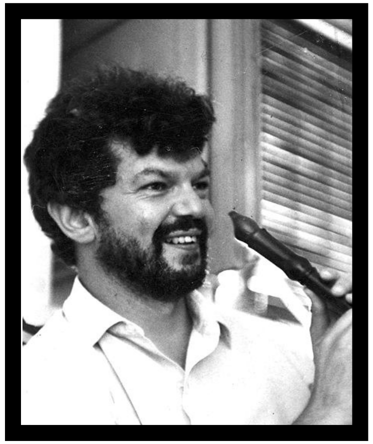
[12] K. A. Johnson, reference 6.



François Englert
(1932 -)

Robert Brout
(1928 - 2011)



Nem a méret számít… Négy alapvető cikk 1964-ből

GLOBAL CONSERVATION LAWS AND MASSLESS PARTICLES*

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(Received 12 October 1964)

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 in-

troduction of vector gauge fields and th
quent breakdown of manifest covarianc
of course, represents a departure fron
sumptions of the theorem, and a limita
its applicability which in no way reflect
general validity of the proof.

In this note we shall show, within the
work of a simple soluble field theory, t
possible consistently to break a symme
the sense that $\sum_k t_{ijk}\langle 0|A_k|0\rangle \neq 0$) witho
ing that $A(x)$ excite a zero-mass partic
this result might suggest a general pro
for the elimination of unwanted massle
it will be seen that this has been accom
by giving up the global conservation law

Gerald Stanford Guralnik (1936 -)

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 theor
of Higgs.[6] Our starting point is the ordinary
electrodynamics of massless spin-zero particl
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

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

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 mod
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
an indefinite metric for one component of φ.)
Since we choose to view the theory as being im
bedded as a linear approximation in the full
theory of electrodynamics, these equations wil
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 transvers
part. The two degrees of freedom of A_k^T com
bine with φ_1 to form the three components of a
massive vector field. While one sees by inspe
tion that there is a massless particle in the the
ory, it is easily seen that it is completely de-
coupled from the other (massive) excitations,

Carl Richard Hagen (1937 -)

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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†National Science Foundation Postdoctoral Fellow.

‡On leave of absence from the University of Rochester, Rochester, New York.

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Thomas Walter Bannerman Kibble (1932 -)

Történeti vázlat

kvantummechanika: 1923 - 25

az atomi elektronhéj elmélete

relativisztikus kvantummechanika:
1925 - 29

alkalmazások:

kvantumkémia
szilárdtestfizika
plazmafizika
magfizika
részecskefizika

kvantum-elektrodinamika (QED):
1928 - 39

háborús megszakítás (1939 – 46)

kvantumelektrodinamika (QED):
döntő áttörés 1949 – 51
(Nobel-díj: 1965)

atommagok + új elemi részecskék

az új elemi részecskék kísérleti vizsgálata + egyedi elméletek

IGÉNY egy általánosabb kvantummező-elmélet (QFT) iránt

a kvantum-elektrodinamika általánosítása:
Yang--Mills-elmélet 1952

a szimmetriák sikeres leírása

DE csak nulla nyugalmi tömegű részecskék lépnek fel!

egyes elemirészecske-kölcsönhatások sikeres leírása

DE fellépnek soha nem látott, nulla nyugalmi tömegű részecskék (Goldstone-bozonok) is

a kvantum-elektrodinamika általánosítása:
Yang--Mills-elmélet 1952

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DE csak nulla nyugalmi tömegű részecskék lépnek fel!

egyes elemirészecske-kölcsönhatások sikeres leírása

DE fellépnek soha nem látott, nulla nyugalmi tömegű részecskék (Goldstone-bozonok) is

Higgs, Englert és Brout, 1964:

Házasítsuk össze a kétféle, nehézségekkel küzdő elméletet!

Eltűnnek a nehézségek: a nulla nyugalmi tömegű részecskék helyett tömegesek lépnek fel --
és a szimmetria is megmarad!

Ez az ötlet az alapja az 1973-ra kikristályosodó részecskefizikai Standard Modellnek

Az összes ismert részecskefizikai szereplő, jelenség, folyamat sikeres (és igen konzervatív) leírása

új folyamatok, részecskék megjóslása

új folyamatok, részecskék kísérleti igazolása (sok Nobel-díj)

a matematikai konzisztencia (renormálhatóság) igazolása (Nobel-díj 1999)

Ez az ötlet az alapja az 1973-ra kikristályosodó részecskefizikai Standard Modellnek

Az összes ismert részecskefizikai szereplő, jelenség, folyamat sikeres (és igen konzervatív) leírása

Már csak az egyik alapvető szereplő, a Higgs által megjósolt skalár bozon hiányzott…

2012-ben az LHC-ben ezt is felfedezték…

összeállt a kirakós játék, a Standard Modell…

A Standard Modell:

késsz, komplett építőkocka a jövő még nagyobb rejtvényeinek kirakásához

(kvantumgravitáció, sötét anyag, stb…)

Englert és Higgs első találkozása 2012-ben a CERN-ben,
a Higgs-bozon felfedezése után

Köszönöm a figyelmet!