Some remarks about the baffling Higgs physics
and the particle mass problem

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Abstract - A statistical model is proposed which ascribes the $Z^0$ mass to
the screening properties of the neutrino Fermi sea (neutrino vacuum). Concern-
ing the fermion masses, some puzzling features of the Higgs mechanism are
examined. Arguments are advanced, based on the Zitterbewegung theory of
the electron substructure, showing that in low energy experiments electron behaves
as an extended distribution of charge, though its size comes out less than $10^{-16}$
cm in high energy experiments. This might be a clue for explaining the origin
of fermion masses without resorting to Higgs field.

Keywords: Higgs physics, weak boson masses, fermion masses, Zitterbewe-
gung theory.

1 - Introduction. - This year, the reconditioned LHC at CERN in Geneva
will start to operate so that experiments either confirming or refuting the Higgs
boson existence will be at last carried out. This renews the interest for this
elusive particle which was hypothized to explain masses of bosons which mediate
the weak force and which should give mass to all massive fermions, leptons and
quarks. For this reasons, it seems timely reconsider the argument of masses
even to highlight possible explanations different from the Higgs hypothesis.

2 - The weak boson masses. - The main feature of weak interaction is
its very short range. In the 1933 Fermi’s theory, it was regarded as a "contact"
interaction acting at zero spatial separation in contrast with electric force me-
diated by photons which acts at large distances. Higgs mechanism ascribes to
weak bosons a "true" inertial mass originated by interaction with a doublet of
scalar fields in $SU(2)$ space, that is, the Higgs field [1]. Owing to energy-time
uncertainty principle, weak bosons of mass $M$ last a time $\delta t \leq \hbar / Mc^2$ so that
range of weak force is $\hbar / Mc$, the boson Compton wavelength. This is like what
occurs with Yukawa force mediated by massive pions. An alternative, more con-
servative, explanation is based on the effect of the neutrino Fermi sea (neutrino
vacuum) on the weak boson propagation. Indeed, owing to the vanishingly small
neutrino mass, neutrino sea is not above-bordered by a forbidden energy gap as
the electron Fermi sea. Consequently, it screens the weak force quite as elec-
trons in the conduction band of metals screen the electric force [2]. In this way,
range of weak force is curtailed, which is equivalent to have massive bosons.

To work out a rigorous treatment of the above screening effect is a rather
exacting task so that it appears suitable use a simplified approach. It is based
on a special application of the Thomas-Fermi method (TF), already examined
ten years ago [3]. A short account of this is given here.

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3 - Screening of weak force in neutrino vacuum. - According to $SU(2) \otimes U(1)$ symmetry, weak interactions are mediated by $W^\pm$ and $Z^0$ bosons. If mass is assigned to these particles, their masses turn out to be related by

$$M_{W^\pm} = M_{Z^0} \cos \vartheta_w,$$

(1)

$\vartheta_w \simeq 28.6^\circ$ standing for the electroweak angle [1]. Equation (1) holds independently of the mass-generating mechanism. Therefore, it is sufficient consider only the neutral $Z^0$ boson. In the unperturbed sea, $\nu_L$ neutrinos of kinetic energies $w = cp$ ranging from 0 to $-\infty$ are present. Taking into account only one spin component, neutrino density of states is related to kinetic energy by [4]

$$\rho(cp) = \frac{(cp)^2}{\pi^2 (\hbar c)^3}.$$

(2)

In presence at the point $r = 0$ of a steady fermion [2], only the time-component $Z^0$ of its potential is different from zero. So, the perturbed neutrino energy is

$$w = cp + U(r),$$

(3)

$U(r) = -Q_{\nu_L} e Z_0(r)$ standing for the neutrino potential energy and $Q_{\nu_L} = 1/\sin 2\vartheta_w$ for the electroweak neutrino charge in units of $e$. Let us examine first the case in which $U$ is negative, that is, $U = -|U|$. On line $w = U(r)$, it follows from equation (3) $cp = 0$. Consequently, the neutrino sea is divided in two regions: one above line $U(r)$ where $cp$ is positive, the other below this line where $cp$ is negative. By denying that neutrino sea is perturbed up to infinite depth [3], we assume that an energy $w_F$ ($w_F = |w_F|$) can be found great enough with respect to $|U(r)|$ ($w_F \gg |U(r)|$ for any $r$) that neutrinos with energy below $-|w_F - |U(r)||$ remain unperturbed. This amounts to say that energy $-w_F$ sets a cut-off in neutrino sea depth. So, neutrino density in the unperturbed Fermi sea is

$$n_F = \int_0^{w_F} \rho (cp) \, d(cp) = \frac{w_F^3}{3\pi^2 (\hbar c)^3}.$$

(4)

In the perturbed sea, integration over $cp$ covers the range $-|w_F - |U||$ to 0 in the negative region and from 0 to $|U|$ in the positive one. Accordingly, the perturbed neutrino density is

$$n = \int_{-(w_F - |U|)}^{0} \rho (cp) \, d(cp) + \int_{0}^{|U|} \rho (cp) \, d(cp) =$$

$$= \frac{1}{3\pi^2 (\hbar c)^3} \left( w_F^3 - 3 |U| w_F^2 + 3 |U|^2 w_F \right).$$

(5)

2) We consider here a point-like fermion since we are concerned with perturbation of Fermi sea in space around the fermion centre. In Section 5 this issue is re-examined.

3) Perturbation in neutrino sea is defined as: $P(cp) = |U| / (|cp| + |U|)$. On line $U(r)$, we have $P(0) = 1$, on the sea border $P(|U|) = 1/2$, but $P(cp) \to 0$ when $cp \to -\infty$. 
By subtracting the unperturbed density \( n_F \), we obtain
\[
 n - n_F = \frac{|U| w_F^2}{\pi^2 (\hbar c)^3} \left( 1 - \frac{|U|}{w_F} \right).
\] (6)

When potential energy \( U \) is positive, that is, \( U = |U| \), equation (6) is found again but with term \( |U| / w_F \) reversed in sign [3]. In reality, taking into account that \( w_F \) is large with respect to \( |U(r)| \), term \( |U| / w_F \) is small and can be disregarded.

By letting \( Q_{fe} \) be the fermion electroweak charge, potential \( Z_0 \) turns out to be ruled by the Poisson-like equation
\[
 \Delta_2 Z_0 = -4\pi Q_{eL} (n - n_F) - 4\pi Q_{f} e \delta (r),
\] (7)
that is, utilizing equation (6),
\[
 \Delta_2 Z_0 - \frac{1}{\lambda_S^2} Z_0 = -4\pi Q_{f} e \delta (r),
\] (8)
where
\[
 \lambda_S = \sqrt{\frac{\pi}{2}} \frac{(\hbar c)^{3/2}}{Q_{eL} \epsilon w_F}.
\] (9)

It follows that the screened potential is
\[
 Z_0 = Q_{fe} \frac{\exp (-r/\lambda_S)}{r}.
\] (10)
This result entitles us to define a “screening mass” \( M_{Z^0} \) by means of a formal Compton wavelength
\[
 \frac{\hbar}{M_{Z^0} c} = \lambda_S.
\] (11)

We consider now the high-energy collisions. Let \( w^{(-)} \), \( w^{(+)} \) be the energies and \( \vec{p}^{(-)} \), \( \vec{p}^{(+)} \) the momenta of the colliding electron-positron pairs. Assuming momentum \( \vec{p}^{(-)} \) opposite to momentum \( \vec{p}^{(+)} \), we have
\[
 w_{c.m.}^2 = (2m_e c^2 + (\vec{p}_{c.m.})^2, 
\] (12)
\[
 w_{c.m.} = w^{(-)} + w^{(+)} \quad \text{and} \quad \vec{p}_{c.m.} = \vec{p}^{(-)} - \vec{p}^{(+)} \text{ standing for energy and momentum in the centre of mass. Disregarding the rest energy } 2m_e c^2, \text{ wavelength corresponding to } p_{c.m.} \text{ is}
\[
 \lambda_{c.m.} = \frac{h}{p_{c.m.}} = \frac{\hbar c}{w_{c.m.}}.
\] (13)
A resonant collision is originated when wavelength \( \lambda_{c.m.} \) becomes equal to the width of the potential well which allows for the electron-positron interaction, that is,
\[
 \lambda_{c.m.} = \lambda_S.
\] (14)
So, taking into account equations (14), (11) and (13), we get

\[ w_{c.m.} = M_{Z^0} c^2, \]

which relates collision energy to \( Z^0 \) mass. When resonance occurs, energy \( w_{c.m.} \) is released through lepton and quark emissions mediated by flavor-diagonal interactions [1]. Utilizing equations (11) and (9), \( Z^0 \) mass, in energy units, turns out to be

\[ M_{Z^0} c^2 = \sqrt{2\alpha / \pi} Q_{\nu_L} w_F = 8.11 \cdot 10^{-2} w_F, \]

\( \alpha \) standing for the fine structure constant. Apart from the neutrino electroweak charge, it is related only to \( w_F \), the energy cut-off in neutrino sea depth. Without this cut-off, that is, for \( w_F \to \infty \), the \( Z^0 \) mass diverges so reducing to zero the range \( \lambda_S \) of weak forces and recovering the old Fermi’s “contact” theory.

By comparing \( M_{Z^0} c^2 \) with its experimental value of about 91 GeV, we obtain \( w_F \) as large as 1120 GeV, while the screening length \( \lambda_S \) turns out to be 2.2 \( \cdot 10^{-16} \) cm. As for the meaning of these figures, it is to be pointed out that on a temperature scale \( w_F \) corresponds to 1.3 \( \cdot 10^{16} \) K. Consequently, for \( T > 10^{16} \) K negative kinetic energy neutrinos are excited to positive energies and the neutrino sea becomes partially empty. This hinders the sea screening properties. Higgs mechanism also gives 10\(^{16} \) K as the temperature which restores the symmetry broken at low temperature [5]. This fact is not surprising because Higgs and screening mechanisms are calibrated on equivalent experimental data. The range 2.2 \( \cdot 10^{-16} \) cm of the weak force entails that the electron size is at least three orders of magnitude smaller than the classic electron radius 2.8 \( \cdot 10^{-13} \) cm.

Conclusion: the above statistical treatment, though not rigorous, readily explains how massless weak bosons may originate short range interactions without resorting to Higgs physics.

4 - The fermion masses. - The Higgs field, which has been assumed to be at the origin of weak boson masses, is also considered in connection with the fermion masses. The theory accommodates the masses of electrons and quarks of the three flavors and sets to zero neutrino masses as a consequence of the non-existence of right neutrinos. Masses are assumed proportional to the Higgs vacuum-expectation-value. Accordingly, three arbitrary coupling factors \( g_e, g_u, g_d \) are considered for the first flavor. So, taking into account the second and third flavor, the theory contains nine undetermined parameters [1].

To detect Higgs boson, various kinds of experiments have been devised based on its decay. But, the mere existence of a decay showing the features expected for the Higgs is not sufficient to conclude that it really concerns the "true" Higgs. It is also necessary that the found results allow an independent determination of the above mentioned nine parameters. Obviously, in lack of this, the Higgs remains nothing more than a conjecture.

But, another point challenges the correctness of the Higgs hypothesis. It is a well-established result of classic electrodynamics that the electromagnetic (e.m.) field shows inertial properties. This field, indeed, shows a momentum density
\[ (\mathbf{E} \times \mathbf{H}) / 4\pi c \] parallel to the Pointing’s propagation vector. On this ground, at the beginning of the past century considerable endeavour was devoted to explain the electron mass as the e.m. mass of a charge distribution of definite size. The advent of quantum mechanics set an end to these attempts, because the supposed point-like nature of elementary particles entails a divergent electron mass. For this reason, the unsolved problem of electron mass was merely put aside by applying renormalization procedures purposely devised. It follows that a viable Higgs mechanism, besides the mass, should likewise explain how the divergent electron e.m. mass is turned off.

In our opinion, in order to manage the tough problem of particle masses, two basic arguments should be considered. The first is that, according to the Copenhagen interpretation of quantum physics, electron is an observable object, not an absolute entity. This means that its features, as expected from theory, depend on the special experiment considered. The second concerns a peculiar property of Dirac equation: the so-called electron Zitterbewegung (Zb) [6,7,8].

It has been shown, when dealing with the expected electron velocity, that in the Fourier expansion of the spinor components, each element of momentum space is associated with oscillations on \( x, y, z \) axes of amplitudes and phases depending on \( \mathbf{p} \). These oscillations are caused by interference beats between positive and negative energy states. By integrating over momenta, we have for \( x \) axis,

\[
\langle x(t) \rangle_{Zb} = \frac{\lambda_C}{4\pi} \int \int \int A_1 \sin \left( 4\pi \frac{t}{T_{Zb}} + \varphi_1 \right) d^3 \mathbf{p},
\]

\( \lambda_C \) standing for the Compton wavelength and

\[
T_{Zb} = \frac{\hbar}{m_e c^2} = 8.1 \cdot 10^{-21} \text{s}
\]

for twice the oscillation period. It follows that a dynamic substructure is originated in which electron move in space around its centre of mass along a complex tissue of closed trajectories [9].

Keeping the above arguments in mind, we point out that high-energy collisions are very fast processes. The collision time \( \tau \) can be roughly identified with the ratio between impact parameter \( b \) and the electron-positron velocities \( c \), that is, \( \tau \simeq b/c \) [10]. But, while impact parameter in direction orthogonal to velocities can be assumed equal to the electron size, that is, \( b_\perp \simeq 10^{-16} \text{cm} \), in parallel direction it is reduced by Lorentz contraction, that is, \( b_\parallel \simeq \sqrt{1 - \beta^2} \cdot 10^{-16} \text{cm} \). Since for 45 GeV electrons we have: \( \sqrt{1 - \beta^2} = 1.1 \cdot 10^{-5} \), we obtain \( \tau \simeq b_\parallel / c = 3.7 \cdot 10^{-32} \text{s} \). This time is short in comparison with the Zb period, in fact: \( \tau / T_{Zb} = 4.6 \cdot 10^{-12} \). It follows that when electrons collide oscillations are stopped and the Zb substructure is not observable. This is like what occurs with a high-speed camera which allows to take steady pictures of a propeller even if it spins very fast.

The situation is opposite in low energy experiments where, in general, energies are determined with high accuracies. For instance, in atomic spectroscopy
indetermination $\delta w$ is less than about $10^{-7}$ eV, that is: $\delta w/m_e c^2 \lesssim 2 \cdot 10^{-13}$. This follows from the fact that Rydberg constant ($R_H = 13.6056981$ eV) is known with seven decimal digits. Considering that equation (18) allows us to write the energy-time uncertainty principle $\delta w \delta t \simeq \hbar$ as a reciprocity relation

$$\frac{\delta w}{m_e c^2} \frac{\delta t}{T_{zb}} \simeq 1,$$

we obtain $\delta t/T_{zb} \gtrsim 5 \cdot 10^{12}$. This large indetermination in time compels us to eliminate time in equation (17) so that oscillations are changed into distributions of probability lying along the electron trajectories \(^4\). Consequently, the observable electron turns out to be a static distribution of charge of definite size. In this way, it might allow for a finite e.m. mass evading divergent results \([11]\). Opposite to the previous propeller example, this is like what occurs with a low-speed camera which takes the picture of the spinning propeller in form of an uniform disk.

5 - Final remarks. - The opinion that the Zb electron substructure should be considered in connection with the mass problem is not new \([9]\). So far, however, it has found scarce attention because most physicists consider the Zb oscillations as a meaningless feature of Dirac’s equation and assume that electron behaves always as a point-like object which, in absence of external forces, cannot change its own velocity. Recently, an experiment has been performed showing clear evidence adverse to this belief \([12]\). In this experiment, a single \(^{40}\text{Ca}^+\) ion trapped in an electromagnetic cage simulates a free electron in an extremely fast quivering motion superimposed on a slow drift, that is, just the Zb motion.

It is to be pointed out , on the other hand, that the TF statistical model, just in order to allow for the finite mass of weak bosons, rules out the assumption of a point-like electron devoid of a dynamic substructure. In fact, this electrons would originate, a divergent neutrino potential energy: $|U(r)| \to \infty$ for $r \to 0$, which according to equation (3) would prevent us from considering a finite cut-off energy $w_F$ and, consequently, a finite boson mass.

References

[1] G. Kane, Modern Elementary Particle Physics (Addison-Wesley, 1987).
[2] N. H. March, Philos. Mag. Suppl. 6 (1957) and references therein.
[3] P. Brovetto, V. Maxia and M. Salis, Il Nuovo Cimento A 112, 531 (1999).
[4] L. D. Landau and E. M. Lifshitz, Statistical Physics (Pergamon Press, Oxford, 1969).
[5] S. Weinberg, Phys. Rev. D 9, 3357 (1974).
[6] P. Dirac, The Principles of Quantum Mechanics (Oxford, 1958) § 69.
[7] B. R. Holstein, Topics in Advanced Quantum Mechanics (Addison-Wesley, 1992).

\(^4\) This is like what occurs with a classic oscillator: $x = A \sin (2\pi t/T)$. Probability that the oscillating particle is found between $x$ and $x + dx$ is: $dP = 2 dt/T$, $dt$ standing for the time required to cross $dx$. We get thus: $dP/dx = 2/\sqrt{T^2 - x^2}$, that is, $dP/dx = 1/\sqrt{x^2 (A^2 - x^2)}$. 


[8] P. Brovetto, V. Maxia and M. Salis, arXiv:quant-ph/0702112v1 12 Feb 2007.
[9] A. O. Barut and A. J. Bracken, Phys. Rev. D 23, 2454 (1981).
[10] E. Fermi, Nuclear Physics (University of Chicago Press, 1955) Ch. II.
[11] P. Brovetto, V. Maxia and M. Salis, arXiv:quant-ph/0512047v1 6 Dec 2005.
[12] R. Gerritsma, G. Kirchmair, F. Zähringer, E. Solano, R. Blatt and C. F. Ross, Nature 463, 68 (2010).