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INSTABILITY OF WEAK VACUUM

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1 Introduction

The Weinberg-Salam theory [1] presents both the unique and instructive method to understand the nature of breaking the weak gauge group by the Higgs particles in wide sector of their masses. The aim to obtain the relations between the set of quarks (in heavy sector), gauge boson and the Higgs boson masses is considered as one of the biggest problems which is provided by great experimental interest for last few years [2].

What are the physically clear bounds on the χ-boson (or the Higgs particle) and heavy quark masses that can be derived by the requiring that the vacuum be stable? For instance, the recent experimental estimation of the top quark mass $m_{top} = 174 \pm 16\ GeV$ [3] allows one, based on the one-loop approximation calculations, to expect the χ-boson mass around the value of $160\ GeV$, taking into account the summarized Standard Model (SM) data analysis on the (m_{top}, m_χ) mass restriction plane inside the allowed area of stable vacuum [2].

Since for certain values of the top quark and the Higgs boson masses there is the metastability of the physical vacuum [4,5] with the lifetime exceeding the present age of the Universe, we shall concentrate our attention on clearing the above-mentioned mass restriction taking into account the scale Λ of validity of the theory.

In the next section, we shall briefly give some semiintuitive estimations of the lower bound of the χ-boson mass. The problem of the vacuum instability will be discussed in section 3. We also review the effects which can be obtained for the χ-boson and the top quark mass constraints using the one-loop effective potential (sec.4). The last section contains the conclusion.

2 Semiintuitive Estimations

The renormalizable scalar field theory is described by the Lagrangian density

$$L = \frac{1}{2}(\partial_\mu \chi)^2 + \frac{1}{2} m_0^2 \chi^2 - \frac{1}{4! N} g_0^2 \chi^4 , \qquad (1)$$

where we consider the real scalar field $\chi(x)$ as a multiplet $\chi^j(x)$ with N components $\chi^j(x) = (\chi^1(x), ..., \chi^N(x))$ belonging to the regular representation of $O(N)$. The scalar field vacuum expectation value $< \chi >$ is of order m_0/g_0, while the χ-boson masses m_χ will be of order of the bare parameter m_0. As was noted by S. Weinberg [6], even for fixed $< \chi > \sim G_F^{-1/2}$ (G_F is the Fermi coupling constant), we can make the χ-boson mass $m_\chi \cong g_0 < \chi >$ as small as we like by taking both g_0^2 and m_0 to be sufficiently small. More instructive models contain the terms of interaction between the field of the χ-particles with others (e.g., the quark field $\Psi(x)$ with the mass m_q). The simplest model is gradient. The classical variant of this model is characterized by the Lagrangian density

$$L = \bar{\Psi}(i\hat{\partial} - m_q)\Psi - g\bar{\Psi}\gamma^\mu\Psi\partial_\mu\chi + \frac{1}{2}(\partial_\mu\chi)^2 + \frac{1}{2}m_0^2\chi^2 - \frac{1}{4!N}g_0^2\chi^4 \,. \qquad (2)$$

The solution of the Lagrange-Euler equation

$$(i\hat{\partial} - m_q)\Psi(x) = g\gamma^\mu\Psi(x)\partial_\mu\chi(x)$$

has the form

$$\Psi(x) = \exp[-ig\chi(x)]\Psi_{free}(x) \,,$$

where $\Psi_{free}(x)$ is the solution of the free Dirac equation. In the quantum case we define the quark field as

$$\Psi(x) =: \exp[-ig\chi(x)] : \Psi_{free}(x) \,.$$

Anticipating the symmetry breaking, the field $\chi(x)$ acquires the classical term χ_c

$$\chi(x) \to \chi'(x) + \chi_c \,.$$

In the tree approximation the mass of the χ-field is

$$m_\chi = \sqrt{\frac{1}{2}g_0^2\chi_c^2 - m_0^2} \,.$$

There is the scale Λ where the scalar field theory described by (2) ceases to be physical. This scale should be chosen as

$$\Lambda \cong m_0 \exp(const/g_0^2) \,,$$

where the interaction (we ignore the $\chi - \Psi$ interaction) became strong. If we need considering such a model at the lower bound, then a replacement of the model would be necessary below energies of order Λ.
Many physicists both in experimental and theoretical areas now believe that the electroweak theory correctly describes phenomena at/or below $O(100\ GeV)$. The characteristic feature of this theory is that particle masses are produced through spontaneous symmetry breaking (plus large Yukawa couplings). In paper [7] the authors claimed that the Born approximation based on the characteristic couplings $\alpha(M_Z) = 1/(128.87 \pm 0.12)$ [8] reproduces all electroweak precision measurements. Taking into account the approximate symmetry

$$\chi(x) = \bar{\chi}(x) + \chi_0 \qquad (3)$$

one can get the following clear representation

$$\exp(iC\chi_0)\exp[i\int\bar{\chi}(x)f(x)d^4x]\exp(-iC\chi_0)=\exp[i\int\chi(x)f(x)d^4x]\ ,$$

where we stipulate that $<\bar{\chi}(x)>=0$. Here

$$C=a(r)+a^*(r)\ ,$$

where the creation $a^*(r)$ and annihilation $a(r)$ bose-operators are defined by means of $[\tilde{\chi}(p)=A(p)+A^*(-p)]$

$$\int A^*(p)\tilde{f}(p)d_4p=a^*(r)\ ,$$

$$\int A(p)\bar{\tilde{f}}(p)d_4p=a(r)$$

for any vector r belonging to the pseudohilbert space $h_1, \tilde{f}(p)\in S(\Re_N)$, $f(x)\in S(\Re^N)$ in the space of S^∞-real functions on $\Re^N$, $\chi_0\in\Re^N$.
The symmetry (3) could be broken by all interactions that have to do with gauge transformation

$$\chi(x)\to\Omega(x)\chi(x)\ . \tag{4}$$

The transformations (3) and (4) form a closed group if there is also invariance under

$$\chi(x)\to\bar{\chi}(x)+\chi_0(x)\ .$$

Using the claim by Novikov et al. [7], we can estimate the strengh $g^2\sim O[4\pi\alpha(M_Z)]$ of the weak gauge field transformation symmetry (4). Then the model restriction scale $\Lambda^2\leq m_\chi^2 O[\alpha^{-1}(M_Z)]$ since the symmetry (3) is broken by $O[\alpha(M_Z)]$ effect. Breaking the above-mentioned symmetry (3) by the $g_0^2\chi^4$ term in (2) leads to the following restriction of m_χ^2:

$$(m_\chi^2/\Lambda^2)\geq O(g_0^2)\geq O[\alpha(M_Z)]\ .$$

On the scale Λ we can get $\Lambda\leq O(\sigma_0)$ while $m_\chi=O(\sigma_0)$ and $\sigma_0=(2\sqrt{2}G_F)^{-1/2}\cong 174\ GeV$. The scale constraint for $\Lambda\leq O(174\ GeV)$ indicates that the model should be unphysical if the energy scale is much higher than $O(\sigma_0)$. Finally, the lower bound constraint of m_χ is defined by

$$m_\chi\geq O[\sigma_0\sqrt{\alpha(M_Z)}]\sim 15\ GeV\ ,$$

that coincides with the ALEPH data of the excluded region for the SM Higgs-boson mass [9]: 32 MeV – 15 GeV at 95% C.L. and 40 MeV – 12 GeV at 99% C.L.

3 The Vacuum Instability

First af all, let us recall some facts about the χ-particle self-couplings and the heavy (top) quark Yukawa couplings to obtain the constraints on the physical χ-boson and the top quark masses. The renormalization group flow to the unknown low χ-boson masses, $m_\chi = \sqrt{2} g_0 v$, and a large enough top quark mass, $m_{top} = g_{top} v$, is given by the equation

$$\frac{dg_0^2}{dq} \cong \frac{3}{4\pi^2}[(g_{top}^2 - A)g_0^2 + B - g_{top}^4] , \tag{5}$$

where $q = \ln(\mu/\mu_0)$, μ is an arbitrary high scale and μ_0 is the χ-boson mass normalized point. The coupling constants A and B are connected with each other and can be read in the standard form

$$A = \frac{1}{4}(g_1^2 + 3g_2^2) \ , \ B = \frac{1}{16}(g_1^4 + 2g_1^2 g_2^2 + 3g_2^4) \ ,$$

$$g_1^2 = \frac{2}{v^2}(m_Z^2 - m_W^2) \ , \ g_2^2 = \frac{2}{v^2} m_W^2 \ ,$$

where v is the weak scale for the χ-boson in the framework of the scale $\Lambda = ve^q$ where g_0^2 is evaluated; m_W and m_Z are the masses of the W- and Z-bosons, respectively; Λ is the scale beyond which the SM is no longer valid. The solution [10]

$$g_0^2(\mu) = R[1 - (\mu_0/\mu)^a] + g_0^2(\mu_0)(\mu_0/\mu)^a \tag{6}$$

obeys equation (5) at $m_{top} \neq \sqrt{m_Z^2/2 + m_W^2}$. The ratio

$$R = \frac{1}{v^2}[\frac{m_W^2(m_W^2/2 + m_Z^2)}{m_{top}^2 - (m_Z^2/2 + m_W^2)} + m_{top}^2 + \frac{m_Z^2}{2} + m_W^2] , \tag{7}$$

$$a = \frac{3}{4\pi^2}(A - g_{top}^2) .$$

The leading positive (negative) contribution to R comes from the first term in (7), if m_{top} is slightly larger (less) as compared to the value of $\sim 103\ GeV$. For the scale, where

$$\Lambda < m_\chi [1 - m_\chi^2/(2v^2 R)]^{-b} \ ,$$

the instability appears, because the coupling constant $g_0^2(\mu) < 0$, if $m_{top} > \sqrt{m_Z^2/2 + m_W^2}$ and

$$b = \frac{4\pi^2 v^2}{3} \frac{1}{\mid m_Z^2/2 + m_W^2 - m_{top}^2 \mid} \ .$$

It is known that the bound for the small χ-boson mass does not depend on the Λ-scale at all. This could be explained by the observation that Λ represents the cutoff where g_0^2 becomes negative. This fact can be used in the instability of the effective potential

$$V_{eff}(\chi) = -\frac{1}{2} m_0^2(\chi)\chi^2 + \frac{1}{4!N} g_0^2(\chi)\chi^4$$

for $\chi >> v$. It means that the χ-boson sector becomes unstable at short distances. In the narrow region of the top quark masses $(m_Z^4/4 + m_W^4/2)^{1/4} < m_{top} < (m_Z^2/2 + m_W^2)^{1/2}$ the ratio (7) is negative in sign. The following restriction on m_χ takes place for the above-mentioned narrow window for m_{top}

$$m_\chi [1 + m_\chi^2/(2v^2 \mid R \mid)]^{1/a} < \Lambda \ ,$$

if $g_0^2(\mu) > 0$. Thus, for the first step, we can conclude that if there is the stable vacuum, there is only a narrow window of the top quark mass allowed. Starting at the point $g_0^2(\mu_0)$ (fixed by the physical χ-boson mass), the $g_0^2(\mu)$-function falls down with increasing μ. For large enough values of μ the $g_0^2(\mu)$-curve reaches the zero point at $\mu = \mu_0[1 - g_0^2(\mu_0)/R]^{1/a}$. The beta-function for the top quark couplings g_{top} [11]

$$\frac{\partial g_{top}(\mu)}{\partial q} = \frac{1}{(4\pi)^2}(\frac{9}{2}g_{top}^2 - \frac{17}{12}g_1^2 - \frac{9}{4}g_2^2 - 8g_3^2)g_{top}$$

allows us to approach the solution of the hierarchy problem only partly in the case of arbitrary m_{top} when the range Λ of validity of our model $\Lambda > \mu$. Since we are interested in the heavy m_{top}, the running coupling $g_{top}(\mu)$ develops the known Landau pole at $g^2_{top}(\mu) = 16\pi^2/[9\ln(\mu/\mu_0)]$. In the physical region the heavy enough top quark mass is restricted by the scale Λ:

$$m_{top} \exp[16\pi^2 v^2/(9m^2_{top})] > \Lambda \ . \tag{8}$$

As we can see from (8), the scale Λ is not so large for the physical true m_{top}. Supposing that our (physical) vacuum is stable, the formula

$$m_\chi [1 - m^2_\chi/(2v^2 R)]^{1/a} < \Lambda \ ,$$

obtained by solving (6) together with (8), translates into the Λ-independent bound on the χ-boson mass, m_χ:

$$\ln \frac{m_\chi}{m_{top}} < (\frac{4\pi v}{3m_{top}})^2 (1 + \frac{3}{2} \frac{m^2_\chi m^2_{top}}{\sum_{V:W\pm,Z} m^4_V - 4m^4_{top}}) \ . \tag{9}$$

Let us very briefly consider the uncertainties which are a direct consequence of the one-loop effect calculations. The one-loop effective potential with keeping only the gauge boson mass and heavy quark mass terms is given by [6]

$$V_{eff}(\chi) = -\frac{1}{2} m^2_0 \chi^2 + [\frac{g^2_0}{4!N} + P(\chi, m)]\chi^4 \ , \tag{10}$$

where

$$P(\chi, m) = \frac{1}{(8\pi\chi^2)^2} Tr[3 \sum_{V:W\pm,Z} m^4_V \ln m^2_V - 4 \sum_{q:c,b,top} m^4_q \ln m^2_q] \ . \tag{11}$$

Taking into account that the second term in (11) is sufficiently larger than the first, the potential (10) can be read as

$$V_{eff}(\chi) = -\frac{1}{2} m^2_0 \chi^2 + [D \ln(\chi^2/M^2) + \frac{g^2_0}{4!N}]\chi^4 \ , \tag{12}$$

where M is a new mass parameter, $D \equiv m^4_{top}/(4\pi < \chi >^2)^2$ and we consider the theory with the local minimum potential scale $< \chi >$ of order m_{top}. Using the facts that the quartic χ^4 term coupling is very weak and a local minimum of the effective potential (12) is expected at $<\chi>$, we can get the following relation

$$<\chi>^2 [\ln \frac{<\chi>^2}{M^2} + \frac{1}{2}] = \frac{m_0^2}{4D} \quad (13)$$

and the mass squared of the χ-boson looks as

$$m_\chi^2 = 8D <\chi>^2 [\ln \frac{<\chi>^2}{M^2} + \frac{3}{2}] \,. \quad (14)$$

The remaining question is whether the symmetry breaking considered here is allowed. It is clear, if the scale M is less than $e^{3/4} <\chi>$, it is definitely ruled out. If $M > e^{3/4} <\chi>$, it will be probably true. Combining both (13) and (14), one can get an upper bound of the χ-boson mass in the unstable phase $m_\chi < m_{top}^2/(2\pi <\chi>)$, which leads to estimation of $m_\chi < 19.5\ GeV$ for $m_{top} = 174\ GeV$ at $<\chi> = 247\ GeV$. For the stable phase,

$$m_\chi \geq \frac{3}{4\pi <\chi>} (\sum_{V:W^\pm,Z} m_V^4)^{1/2} = 6.8\ GeV \ ,$$

if $m_{top} \sim 100\ GeV$.

4 The One-Loop Potential Effect

Physically, the issue of whether or not the weak vacuum is necessarily unstable will depend on the scheme of calculation of the effective potential. The renormalization group included potential is not considered here since we are not interested in a very large set of scalar fields $\chi^j(x) = (\chi^1, ..., \chi^N)$. We shall use the effective potential (obeying the Lagrangian density (1)) in the one-loop approximation [2]:

$$V_{eff}(\chi) = V_0(\chi) + V_1(\chi) \ , \quad (15)$$

$$V_0(\chi) = -\frac{1}{2} m_0^2 \chi^2 + \frac{1}{4!N} g_0^2 \chi^4 \ ,$$

$$V_1(\chi) = \frac{1}{64\pi^2} [G\chi^4 \ln \frac{\chi^2}{M^2} + (3g_0^2\chi^2 - m_0^2)^2 \ln \frac{3g_0^2\chi^2 - m_0^2}{M^2}$$

$$+3(g_0^2\chi^2 - m_0^2)^2 \ln \frac{g_0^2\chi^2 - m_0^2}{M^2}] , \qquad (16)$$

where $G \equiv (3/8)(g_1^4 + 2g_1^2g_2^2 + 3g_2^4 - 4g_Y^4)$, $g_Y = \sqrt{2}m_{top}/v$ and M is the renormalization scale which can be chosen to be v in the calculations. The contribution of the heavy top quark to $V_1(\chi)$ at $\chi >> v$ is negative and for $m_{top} > (m_W^4/2 + m_Z^4/4)^{1/4}$ the potential (15) is unbounded from below. A spontaneous generation of an infinitely large χ-field should take place:

$$V_1(\chi >> v) \sim \frac{1}{64\pi^2}(G + 9g_0^4)\chi^4 \ln \frac{g_0^2\chi^2}{M^2} .$$

Higher order perturbative contributions can modify the expression (16) at the asymptotically large values of $\chi \sim v \exp(const/g_0^2)$ [12]. But we shall operate with the field scale $\sim v$. The effective potential (15) has a local minimum at the point $< \chi >$ of the scalar field vacuum expectation value given by

$$\begin{aligned}(\frac{<\chi>}{4\pi})^2[G(\ln\frac{<\chi>^2}{M^2} + \frac{1}{2})\\
+3g_0^2(3g_0^2 - \frac{m_0^2}{<\chi>^2})(\ln\frac{3g_0^2<\chi>^2 - m_0^2}{M^2} + \frac{1}{2})\\
+3g_0^2(g_0^2 - \frac{m_0^2}{<\chi>^2})(\ln\frac{g_0^2<\chi>^2 - m_0^2}{M^2} + \frac{1}{2}) + \frac{8\pi^2}{3N}g_0^2] = m_0^2 \qquad (17)\end{aligned}$$

and the χ-boson mass squared is

$$\begin{aligned}m_\chi^2 = \frac{1}{2}(\frac{<\chi>}{2\pi})^2[G(\ln\frac{<\chi>^2}{M^2} + \frac{3}{2}) + \frac{8\pi^2}{3N}g_0^2\\
+9g_0^4(2 + \ln\frac{3g_0^2<\chi>^2 - m_0^2}{M^2} + \frac{1}{3}\ln\frac{g_0^2<\chi>^2 - m_0^2}{M^2})] . \qquad (18)\end{aligned}$$

To obtain the lower bound restrictionon m_χ^2, we express the potential (15) at the point (17) taking into account (18). The N-independent potential has the form:

$$\begin{aligned}V_{eff}(\chi) = -\frac{1}{8}m_\chi^2 <\chi>^2 + \frac{G}{128\pi^2} <\chi>^4\\
+\frac{m_0^4}{64\pi^2}(\ln\frac{3g_0^2<\chi>^2 - m_0^2}{M^2} + 3\ln\frac{g_0^2<\chi>^2 - m_0^2}{M^2})\end{aligned}$$

$$-\frac{3}{32\pi^2} g_0^2 < \chi >^2 (m_0^2 - g_0^2 < \chi >^2) \ . \tag{19}$$

It is clear that the effective potential (19) turns upward, leaving a new stable point (new vacuum) on a large enough field scale χ_1 which is exponentially larger than the weak scale, $\chi_1 \sim < \chi > \exp(const/g_0^2)$. This new scale could be extended beyond the Planck scale or the scale where the considered model might apply. We do not know the natural receipt of the stabilization of the effective potential (19) except some manipulation or speculation on the bare parameters.

Now, suppose that the renormalization scale is compared to $< \chi >$. Then, the lower bound on m_χ^2 is:

$$m_\chi^2 \geq \frac{1}{2\pi^2}\{\frac{1}{8}G < \chi >^2 - \frac{3}{2} g_0^2 m_0^2 (1 - \frac{g_0^2 < \chi >^2}{m_0^2})$$

$$+ \frac{m_0^4}{< \chi >^2}[\ln \frac{m_0^2}{< \chi >^2} + \frac{1}{4} \ln(1 - \frac{6 g_0^2 < \chi >^2}{m_0^2}(1 - \frac{2 g_0^2 < \chi >^2}{m_0^2}))]\} \ .$$

If the typical scalar mass in the effective potential (15) is of order m_0, then $< \chi > \sim m_0/g_0$, while the χ-boson mass will be of order m_0. Neglecting the contribution which is proportional to $\sim g_0^4$, one can obtain the lower bound on the χ-boson mass: $m_\chi \geq 110.5\ GeV$ for $g_0^2 \sim 0.025$ [13] at small values of χ, while for $g_0^2 \sim 0.001$ the restriction is $m_\chi \geq 22.1\ GeV$.

In the case when the bare mass parameter squared $m_0^2 \sim \xi m_\chi^2 (\xi \sim O(1))$, the mass upper bound looks as

$$m_\chi \leq \{\frac{1}{\xi}[\frac{1}{2} + (\frac{2}{\xi})^{1/2} \pi (1 + \frac{\xi}{16\pi^2}(1 + 14 g_0^2))]\}^{1/2} < \chi > \ .$$

5 Conclusion

The naive consideration based on breaking the symmetry (3) by the $g_0^2\chi^4$-term leads to the fact that the values of $m_\chi < 15\ GeV$ are unnatural, which coincides with the ALEPH data [9]. For the top quark mass which is of the order of masses of the gauge bosons, there is the stable vacuum if there is only a narrow window of the top quark mass allowed: $78.5\ GeV < m_{top} < 103\ GeV$. For the scale where $\Lambda < m_\chi [1 - m_\chi^2/(2v^2 R)]^{-b}$, the instability

appears for $m_{top} > 103\ GeV$, and the bound for small m_χ does not depend on the Λ-scale at all. The Λ-independent bound for χ- boson mass (9) is true and it is provided by the arbitrary top quark masses only under the Λ-restricted condition (8).
The one-loop effective potential consideration for χ-boson mass leads to the lower bound $m_\chi \geq 110.5\ GeV$ for $g_0^2 \sim 0.025$ [13], while $m_\chi(g_0^2 \sim 0.001) \geq 22.1\ GeV$. For large enough top quark mass $[G(m_{top} = 174\ GeV) = -1.42]$, the condition

$$x[x \ln x - 2\pi^2/\xi - (7/2)g_0^2] + G/8 \leq 0$$

is true always for arbitrary m_χ and ξ ($x = \xi m_\chi^2/ < \chi >^2$). The χ-boson mass will not exceed the 1 TeV level in the case when $\xi \sim 0.5$ and $m_{top} \sim 174\ GeV$, but increasing of ξ leads to the effect that χ-boson might be lighter ($m_\chi \cong 549\ GeV$ at $\xi = 1$). The above-mentioned upper bound for small ξ coincides with the estimations of the Higgs mass upper bound in SM, $M_{Higgs} < 1\ TeV$ [14].

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