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H. Schlereth

Paul Scherrer Institut
CH-5232 Villigen PSI, Switzerland

and

Institute for Theoretical Physics
University of Zürich, Switzerland

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Paul Scherrer Institute
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Abstract

From the fact that massive Yang-Mills theory can be rewritten as a gauge invariant nonlinear σ - model we derive in a large N limit ($N = 4$ in the S.M.) the parameters of the S.M. Higgs sector which appears below a typical momentum scale of order 1.2 TeV. The Higgs mass becomes $m_H = \sqrt{3} v \cong 426 \text{ GeV}$. The symmetry breaking effect of the large top mass plays a major role in our formalism.

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$$\Omega(x) = \frac{1}{\sqrt{2}} \left(\frac{\sigma(x)}{v} + i \frac{\vec{T} \vec{\pi}(x)}{v} \right) \quad (6)$$

$$\sigma^2(x) + \vec{\pi}(x)^2 = v^2$$

where $\{\vec{T}\}$ are the three SU(2) generators.

We further add a Yukawa coupling of the top quark to the field $\sigma(x)$ (the coupling to the $(N-1)$ fields $\vec{\pi}(x)$ does not contribute in the approximation we are performing below).

Discussing the ensuing system in the zero gauge coupling limit, $g = 0$, we have then (it should be possible to include the effects of $g \neq 0$ at a later stage perturbatively)

$$Z = \int D\sigma(x) D\vec{\pi}(x) D\alpha(x) DQ_L(x) DQ_R(x) e^{-\frac{1}{T} \int d^4x \left[\frac{1}{2} (\partial_\mu \sigma(x))^2 + \frac{1}{2} (\partial_\mu \vec{\pi}(x))^2 + \frac{1}{2} \alpha(x) (\sigma(x)^2 + \vec{\pi}(x)^2 - v^2) - g_t \cdot \vec{Q}_L(x) t_R(x) \sigma(x) + \vec{Q}_L(x) \not{\partial} Q_L(x) + \vec{Q}_R(x) \not{\partial} Q_R(x) \right]} \quad (7)$$

Here we introduced a coupling constant $\frac{1}{T}$ (below we rescale back to canonical kinetic terms). This parameter is necessary to make the σ - model critical at large distances, that is, to locate the infrared fixed point (IRFP). At this point there is a relation to the parameters of the linear

σ - model whose relevance is discussed below. The Lagrange multiplier $\alpha(x)$ takes care of the constraint in (6), g_t is the Yukawa coupling and $Q_{L,R} = \begin{pmatrix} t \\ b \end{pmatrix}_{L,R}$ denotes the quark doublet in the third family (L, R means left respectively right and $\not{\partial} = \gamma^\mu \partial_\mu$).

We take $T \sim \frac{1}{N}$ and study this system in the leading $\frac{1}{N}$ ($N = 4$ in the S.M.) approximation by constructing the saddle point of the effective action in $(\sigma(x), \alpha(x))$ field space [5]. From the action describing the fluctuations around the saddle point we read off, roughly speaking, the parameters of the Higgs sector. Important technical details are that it will turn out that the mass of the top quark, m_t , is the exclusive source of all symmetry breaking and that the Higgs sector is constructed in the infrared where the linear and the nonlinear σ - model coincide [6].

The Higgs mass becomes then

$$m_H = \sqrt{3} v = 426 \text{ GeV} \quad (8)$$

which, luckily, is substantially away from the "light" Higgs characteristic of supersymmetric theories and requires for its discovery the LHC. For instance in the minimal supersymmetric S.M. [2] the lightest Higgs has $m_h < 130 \text{ GeV}$ [7].

The momentum spectrum of the fluctuations around the saddle point turns out to be limited to roughly $\leq 1 \text{ TeV}$. It is this limitation which brings to life the scalar radial degree of freedom which describes the Higgs.

The high momentum parts of the σ - model variables have in fact been integrated out, and the Higgs is described by large scale "block" variables [8].

Although we do not use the lattice regularization in this paper a look at the lattice phase diagram of compact massive Yang-Mills theory is helpful at this point [9].

To understand the Higgs structure of the standard Model [S.M] [1] the present trend in particle physics is to follow the guidance of higher symmetries, for instance supersymmetry, in the introduction of extended Higgs sectors with more than one Higgs doublet [2]. Whether the introduction of more doublets is really helpful to understand the one doublet really needed in the S.M. may be doubted. Of course this question must be finally answered by experiment [7].

In fact the Higgs in the S.M. is the necessary ingredient for a theory with a massive nonabelian vector particle to be insensitive to very high momenta. In other words to be renormalizable [3]. From this point of view it would be sufficient if the Higgs appeared as a low energy excitation just ensuring that, at least perturbatively, a consistent low energy theory emerges. With this in mind we start in this note from the other extreme, namely with no Higgs to begin with. In the low momentum region, roughly $|p| \leq 1 \text{ TeV}$, this scalar excitation will however effectively appear. In detail we develop this idea as follows.

It is well known [4] that massive Yang-Mills theory represented by the partition sum

$$Z = \int DA_\mu(x) e^{-\int d^4x \left[\frac{1}{4} F_{\mu\nu}^A(x)^2 + \frac{1}{2} m^2 A_\mu^A(x)^2 \right]} \quad (1)$$

(A runs over the adjoint representation of some gauge group, say SU(M)) can be rewritten in a gauge invariant way by the addition of further variables $\Omega(x)$ without changing the physical content

$$Z = \int DA_\mu(x) D\Omega(x) e^{-\int d^4x \left[\frac{1}{4} F_{\mu\nu}^A(x)^2 + \frac{1}{2} v^2 |D_\mu \Omega(x)|^2 \right]} \quad (2)$$

With an eye to the weak gauge bosons, we have parametrized the mass as

$$m = \frac{1}{2} g v \quad (3)$$

such that g is considered the weak gauge coupling constant and v is identified with the Fermi scale

$$v = 246.2 \text{ GeV} . \quad (4)$$

In the form (2) the theory is invariant under an SU(M) gauge transformation $h(x)$

$$\begin{aligned} A_\mu(x) &\rightarrow h(x) A_\mu(x) h(x)^{-1} - \frac{i}{g} \partial_\mu h(x) \cdot h(x)^{-1} \\ \Omega(x) &\rightarrow h(x) \Omega(x) \end{aligned} \quad (5)$$

since the covariant derivative is defined as $D_\mu \Omega(x) = \partial_\mu \Omega(x) - \frac{g}{2} A_\mu(x) \Omega(x)$.

The matrix $\Omega(x)$ represents the degrees of freedom of an $O(N)$ ($N = M^2$) nonlinear σ -model, $\Omega^\dagger(x)\Omega(x) = 1$. Specializing to $M = 2$ we write

Using dimensional regularization ($d = 4 - \varepsilon$) the coupling constant is given by [11]

$$\frac{1}{g^2} = \frac{2}{3} N_f \frac{1}{(4\pi)^2} \Gamma\left(\frac{\varepsilon}{2}\right) \quad (13)$$

We are interested in the case $N_c = 2$ such that it is natural to try to identify g with the tree level weak coupling constant. Putting for the weak boson mass $m^{\text{tree}} \cong 77 \text{ GeV}$ we find from (3,4)

$$g^{\text{tree}} \cong 0.63. \quad (14)$$

Further from (12) we conclude that for $N_c = 2$ there must be at least

$$N_f = 12 \quad (15)$$

flavours. This is just the number of weak lefthanded doublets for three quark/lepton families. It may suggest a deeper relation between the heavy and the light fermions in (10).

For this ‘minimal’ model we find from (13) and (14) for the cut-off (we may translate $\frac{2}{\varepsilon} = \ln\left(\frac{\Lambda}{M_0}\right)^2$ with a reference scale M_0 . A reasonable choice is $M_0 \cong v$.)

$$\varepsilon = \frac{(g^{\text{tree}})^2}{\pi^2} \cong 0.04. \quad (16)$$

In performing this calculation from (13) we have assumed that in the complete formula from dimensional regularization [see second of [11]] the $\frac{2}{\varepsilon}$ term is the dominant one. Note that $\frac{2}{\varepsilon} \cong 50$.

Let us now assume (although this is not necessary, as already stressed) that (2) respectively (7,9) originate indeed from this model. Then it is natural to use the same cut-off given by (16) to regularize all logarithmic loop divergencies in the theory.

We will carry out this program in the following, but will also stress that our basic results for the Higgs sector can be obtained without it, proceeding directly from (9).

Once the heavy fermion part in (10) has been effectively replaced by (11) we can calculate perturbatively in g (after rescaling $V_\mu^A \rightarrow g V_\mu^A$) the charged weak boson mass arising from the transverse part of the top-bottom vacuum polarization

$$m_W^2 = \Pi_1(p^2) \Big|_{p^2 = m_W^2} \quad (17)$$

(for notation etc. see [12]).

If we then regularize the logarithmic divergence at the r.h.s. of (17) with the cut-off (16) and keep only the leading $\frac{2}{\varepsilon}$ part we find the approximate relation

$$m_W \cong \frac{g^{\text{tree}}}{\sqrt{2}} m_t. \quad (18)$$

So with m_t in the CDF - range [13]

$$m_W \cong 70 \text{ GeV} \quad (19)$$

As one approaches the transition line from the Higgs to the confinement phase there is a light (compared to the cut-off scale) scalar degree of freedom in the spectrum which for its gauge invariant description naturally requires the introduction of linear σ - model variables. It is important to construct the Higgs sector explicitly in this region as an infrared companion of massive Yang-Mills theory [10].

This is basically what we intend to do. However we deviate from the picture just sketched by giving a dominant role to the massive top quark.

If we integrate out in (7) the $(N-1)$ fields $\bar{\pi}(x)$ and keep in the fermion sector only the top contribution we obtain

$$Z = \int D\sigma(x) D\alpha(x) D\bar{t}(x) Dt(x) e^{-\frac{1}{T} \int d^4x \left[\frac{1}{2} (\partial_\mu \sigma(x))^2 - \frac{1}{2} \alpha(x) v^2 + \frac{1}{2} \alpha(x) \sigma^2(x) + \frac{T}{2} (N-1) \text{tr} \ln (\square + \alpha(x)) - g_t \bar{t}(x) t(x) \sigma(x) + \bar{t}(x) \partial t(x) \right]} \quad (9)$$

All our subsequent results follow from this form.

Before elaborating on (9) we will briefly sketch a derivation of (2) including the Yang-Mills term from an underlying purely fermionic theory.

We start from the following theory which contains besides the (t,b) - quarks also heavy vectorlike (Dirac) fermions $\psi^{i,a}(x)$ with i running up to N_f 'flavours' and a up to N_c 'colors' in the fundamental representation of some $SU(N_f)$ respectively $SU(N_c)$

$$Z = \int DV_\mu^A(x) \int D\bar{\psi}(x) D\psi(x) DQ_L(x) DQ_R(x) e^{-\int d^4x [\bar{\psi}(x) (\partial + M) \psi(x) + i V_\mu^A(x) \bar{\psi}(x) \gamma_\mu T^A \psi(x) + \bar{Q}_L(x) \partial Q_L(x) + \bar{Q}_R(x) \partial Q_R(x) + i V_\mu^A(x) \bar{Q}_L(x) \gamma_\mu T^A Q_L(x) + m_t \bar{t}(x) t(x) + m_b \bar{b}(x) b(x)]} \quad (10)$$

M denotes the mass of the Dirac fermions, $\{T^A\}$, $A = 1 \dots, N_c^2 - 1$, are the generators of $SU(N_c)$ and $V_\mu^A(x)$ is a spin one field in the adjoint representation of this group. The b-mass, m_b , is neglected in the following. Integrating out in (10) the heavy fermions on momentum scales $\frac{p^2}{M^2} < 1$ [11] one finds for $V_\mu^A(x)$ the kinetic term of Yang-Mills theory

$$S = \frac{1}{4g^2} \int d^4x F_{\mu\nu}^A(x)^2 \quad (11)$$

with

$$F_{\mu\nu}^A(x) = \partial_\mu V_\nu^A(x) - \partial_\nu V_\mu^A(x) + f^{ABC} V_\mu^B(x) V_\nu^C(x)$$

if

$$N_f > \frac{11}{2} N_c \quad (12)$$

Then from (23b)

$$g_t < \bar{t}t > = v^3. \quad (26)$$

This has to be understood as an equation for the top mass. The requirement on m_W can only be implemented if (26) is consistent with m_t in the CDF-range (we accept this here as the relevant constraint [13]). In the lowest approximation (26) reads $\left(g_t = \frac{m_t}{v}\right)$

$$v^4 = 12 m_t^2 \frac{1}{(2\pi)^4} \int \frac{d^4 p}{p^2 + m_t^2}. \quad (27)$$

(The factor 12 comes from the trace over the Dirac spin and the SU(3) color multiplicity).

If we evaluate the integral at the r.h.s. by dimensional regularization with a cut-off ε close to 4 ($\varepsilon = 4 - d$) we obtain

$$v^4 = m_t^4 \frac{3}{\varepsilon} \frac{1}{2\pi^2} \quad (28)$$

It was again assumed that the $\frac{2}{\varepsilon}$ term is the dominating one, which is justified because $\ln \frac{m_t}{v}$ is small, see (29).

Then the cut-off (16) implies from (28)

$$m_t \cong 176 \text{ GeV}. \quad (29)$$

This is well within the CDF range $m_t = 174 \pm 10_{-12}^{+13}$ GeV [13]. This result may encourage us to accept (26) (with the appropriate regularization), as a basic relation between top mass and Fermi scale and to hope for a deeper understanding within a (not yet existing) complete theory of quark and lepton masses.

This discussion gives for the saddle point in $(\sigma(x), \alpha(x))$ - space from (23a, 23b, 25, 26)

$$\bar{\alpha} = v^2, \quad < \sigma > = v. \quad (30)$$

Note that the physical requirement on m_W expressed in (25) fixes the renormalization constant in (22) completely. If we did not want to start from (10) but proceed directly from (9) we could still insert into (23b) the physical values of m_t and v and calculate the cut-off such that (19) is valid. Clearly we would then find $\varepsilon \cong 0.04$, and so would still have the selfconsistent relation (18,19) with g^{tree} taken from (14). We find it interesting that the cut-off needed here is just the one calculated from (13) and (14). Moreover the start from (10) has at least in principle observable consequences. Namely in the case that the mass M of the Dirac fermions is within the reach of the LHC. In fact M cannot be arbitrarily large if we insist that the $\frac{2}{\varepsilon}$ -

term in (13) dominates the combination $(-\gamma + \ln 4\pi - \ln \frac{M^2}{\mu^2})$, which for this reason is not

included in (13). γ denotes Mascheroni's constant. A reasonable choice for the reference scale is $v \leq \mu \leq 1 \text{ TeV}$ (for the upper bound see below). Therefore lower values for M include the range 250 GeV - 1 TeV. Models for the interactions of these Dirac fermions and their eventual signatures will be discussed elsewhere.

(in fact $m_W \cong 77 \text{ GeV}$ for $m_t = 176 \text{ GeV}$).

We find it interesting that (16,17) is consistent with m_W^{tree} used above to calculate g^{tree} and subsequently the cut-off from (13).

We go now back to (9) and study the leading large N contribution by constructing the saddle point solutions obeying

$$\frac{\partial V}{\partial \sigma(x)} = 0, \quad \frac{\partial V}{\partial \alpha(x)} = 0 \quad (20)$$

of the potential $V = V(\alpha(x), \sigma(x))$ defined by

$$Z = \int D\sigma(x) D\alpha(x) D\bar{t}(x) Dt(x) e^{-\frac{1}{T} \int d^4x \left[\frac{1}{2} (\partial_\mu \sigma(x))^2 + \bar{t}(x) \partial_t t(x) + V(\alpha(x), \sigma(x)) \right]} \quad (21)$$

If we renormalize by the reparametrization [5]

$$T = Z \tilde{t} \equiv \frac{\tilde{t}}{1 - \frac{N-1}{d-2} \Gamma\left(\frac{d}{2}\right) \Gamma\left(2 - \frac{d}{2}\right) \cdot \tilde{t}} \quad (22)$$

(a factor $\frac{2\pi^{\frac{d}{2}}}{(2\pi)^d}$ is absorbed in the new coupling parameter $\tilde{t}$) we obtain from (20, 21, 9)

$$\bar{\alpha} = \frac{g_t \langle \bar{t} t \rangle}{\langle \sigma \rangle} \quad (23a)$$

$$v^2 - \langle \sigma \rangle^2 = \frac{2}{d-2} (N-1) \Gamma\left(\frac{d}{2}\right) \Gamma\left(2 - \frac{d}{2}\right) \cdot \frac{\tilde{t}}{(2\pi^{\frac{1}{2}})^d} \left(v^2 - \frac{g_t \langle \bar{t} t \rangle}{\langle \sigma \rangle} \right) \quad (23b)$$

where (23a) was already used to obtain (23b). Also the factor $\frac{2\pi^{\frac{d}{2}}}{(2\pi)^d}$ is extracted from $\tilde{t}$ in (23b).

In the Yukawa coupling term we have approximated $\bar{t}t$ by $\langle \bar{t}t \rangle$. The other terms will finally (when $\sigma(x)$ is established as the S.M. Higgs) give the S.M. perturbative radiative corrections from the Higgs-top sector.

Further in (23a, b) $\sigma(x)$ and $t(x)$ denote renormalized fields

$$\sigma(x) \rightarrow \sqrt{Z} \sigma(x), \quad t(x) \rightarrow \sqrt{Z} t(x) \quad (24)$$

and $\bar{\alpha}$, $\langle \sigma \rangle$ the saddle point values of $\alpha(x)$ respectively $\sigma(x)$.

We solve (23b) under the requirement that the form of the weak boson mass, m_W , be consistent with arising from spontaneous symmetry breaking. If $\sigma(x)$ is going to be identified with the radial Higgs field the unique solution is

$$\langle \sigma \rangle = v. \quad (25)$$

It is easy to estimate that only for momentum scales of order of 4TeV the last (nonlocal) term in (33) becomes comparable to $\frac{2}{\epsilon}$. Therefore we neglect this term in the following. Now we construct the Higgs sector in the infrared by fixing $\tilde{t}$ such that the nonlinear σ - model coincides with the linear one [for large N , see the last reference of [5]].

After appropriate rescaling of the fields by $\tilde{t}$ the linear $O(N)$ σ - model action

$$S = \int d^4x \left[\frac{1}{2} (\partial_\mu \sigma(x))^2 + \frac{1}{2} (\partial_\mu \bar{\pi}(x))^2 + \frac{1}{2} m_o^2 (\sigma(x)^2 + \bar{\pi}(x)^2) + \frac{g_o}{4!} (\sigma^2(x) + \bar{\pi}(x)^2)^2 \right] \quad (34)$$

can be cast in the form

$$S = \frac{1}{\tilde{t}} \int d^4x \left[\frac{1}{2} (\partial_\mu \sigma(x))^2 + \frac{1}{2} \alpha(x) \sigma(x)^2 - \frac{3}{2} \frac{\tilde{t}}{g_o} (\alpha(x) - m_o^2)^2 \right. \\ \left. + \frac{N-1}{2} \tilde{t} \operatorname{tr} \ln(\square + \alpha(x)) - g_t \tilde{t}(x) t(x) \sigma(x) + \tilde{t}(x) \partial t(x) \right] \quad (35)$$

where g_o and m_o are a bare coupling and mass parameter respectively. In (35) the $(N-1)$ fields $\bar{\pi}(x)$ are integrated out, and the top quark part is included. In this form the linear model can be treated in the same way as the nonlinear one starting from (9). To perform the large N limit we assume $g_o \sim 1/N$. It is then straightforward to find the conditions on the coupling parameters

$$\frac{6}{g} = \frac{N-1}{2} \Gamma\left(\frac{\epsilon}{2}\right) \quad (36a)$$

$$\tilde{t} = \frac{1}{\frac{(N-1)}{(4\pi)^2} \cdot \frac{2}{\epsilon} \left(1 - \frac{m^2}{v^2}\right)} \quad (36b)$$

where g and m are a renormalized coupling and mass parameter of the linear model. In (36a)

the factor $\frac{2\pi^{\frac{d}{2}}}{(2\pi)^d}$ is absorbed in $\frac{6}{g}$. We see that the two models coincide at the IRFP (36a) of

the linear model if $\tilde{t}$ is given by (36b). We define the S.M. Higgs sector as the scalar theory that arises at this point. The physical reason is that in the gauged case, $g \neq 0$, this leads to a renormalizable Yang-Mills Higgs system in the broken phase with a *linear* realization of the gauge symmetry³. The details of the gauged case remain to be worked out.

Keeping in line with the physics of the nonlinear model the only mass parameter that should appear is the top mass. Therefore we put $m^2 = 0$. Then from (36b) ($\frac{2}{\epsilon} = 50$; $N = 4$, the S.M. case)

$$\tilde{t} = \frac{1}{\frac{N-1}{(4\pi)^2} \cdot \frac{2}{\epsilon}} = 1.05 \cong 1. \quad (37)$$

Next, after use of (37), we integrate out $\delta \alpha(x)$ from (33) and obtain with a change of notation $\delta \sigma(x) \equiv H(x)$ for Higgs and a rescaling of the fields by $1/\sqrt{\tilde{t}}$

$$S_{Higgs} = \int d^4x \left[\frac{1}{2} (\partial_\mu H(x))^2 + \frac{1}{2} (v\sqrt{3})^2 H(x)^2 + v H(x)^3 + \frac{1}{4} H(x)^4 \right] \quad (38)$$

³ At long distances it should not matter whether in (2) we couple a linear or a nonlinear σ - model to the gauge field. Then the nonlinear σ - model should determine the parameters of the corresponding linear one and vice versa.

The Higgs sector is now completely specified. In particular we read off the Higgs mass

$$m_H = \sqrt{3} v = 426 \text{ GeV}. \quad (39)$$

Equation (37) is equivalent to the Landau pole $T = \infty$ as follows from (22). Away from this point $\tilde{r}$ is smaller than the value given by (37). This is to keep T positive and implies that (39) is the minimum value of the Higgs mass. This is consistent with the idea that the light Higgs shows up in the infrared region.

The Higgs potential in (38) is classically stable: $V'(H) = 0$ has only $H(x) = 0$ as real solution. Since $m_H < 2m_\pi$ ($m_\pi = v$, as can be gleaned from (9) and (30)) the Higgs particle is stable in the approximation that led to (33). The region of validity of (38) can be roughly estimated from the fact that (31) is only reasonable for $\delta\sigma(x) \equiv H(x) < v$.

Turning this into a constraint on the momentum cut-off Λ_H

$$\langle H(0)^2 \rangle \equiv \frac{1}{(2\pi)^4} \int^{\Lambda_H} \frac{d^4 k}{k^2} < v^2 \quad (40)$$

we find

$$\Lambda_H < 4\pi v \equiv 3 \text{ TeV} \quad (41)$$

This is the typical scale of a σ - model with " f_π " = v . However a stronger constraint than this can be derived because the expansion around the constant values (30) also requires

$$\frac{1}{2} \cdot (\partial_\mu H)^2 < v^4. \quad (42)$$

In analogy to (40) this gives⁴

$$\Lambda_H \leq 1.2 \text{ TeV}. \quad (43)$$

In conclusion we remark that certain corrections are not yet included in the calculation of the Higgs parameters. There are the next to leading $\frac{1}{N}$ corrections to the saddle point. A second task is to include the gauge coupling and to define the S.M. radiative corrections. This will be treated in subsequent work. Finally one should end up with predictions of m_W , m_t and m_H in terms of v . The calculation of the Higgs parameters is universal in the sense that no specification of "new physics" in terms of higher derivative contributions in the σ - model was made. In general such terms are present, but are suppressed below $\approx 1.2 \text{ TeV}$. Beyond this scale universality is lost and it starts to matter where (9) comes from. Note also that the value of m_{Higgs} is not strongly dependent on the precise value of the cut-off ε as long as $\frac{2}{\varepsilon}$ is the leading term. This is clear from (33) and (36) where $\frac{2}{\varepsilon}$ cancels. In particular for $\varepsilon \rightarrow 0$ ($\Lambda \rightarrow \infty$) and v fixed we obtain instead of (38) a free massive scalar theory with the mass as in (39) and $m_W, m_t \rightarrow 0$. The effective gauge and Yukawa couplings vanish logarithmically. This is in accord with „triviality“ [16]. Also ε can be determined without commitment to detailed new physics by requiring consistency in (17). We have however also discussed an extended framework by starting from (10) and calculating ε from the weak SU (2) gauge charge.

⁴ (39) and (43) are consistent with stability bounds derived within the S.M. [17].

As for the present experimental situation it is known that within 2σ no determination on m_H can be made [7]. A recent analysis [15] including LEP + SLD + CDF data presents a value of roughly $m_H \approx 500$ GeV at 95 % C.L. This is in the same ball park as our value. However it is not clear how much is due to the disagreement between LEP and SLD.

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