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INTERNATIONAL CENTRE FOR THEORETICAL PHYSICS

INDUCED HIGGS COUPLINGS AND SPONTANEOUS SYMMETRY BREAKING

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INDUCED HIGGS COUPLINGS AND SPONTANEOUS SYMMETRY BREAKING *

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ABSTRACT

It is shown that spontaneous symmetry breaking can arise in a non-Abelian gauge theory free of quartic scalar couplings only if fermions are present in the theory. A sufficiency condition is developed for positivity of the induced Φ^4-potential as $\Phi \to \infty$. The same condition guarantees the existence of asymptotically free positive-eigenvalue solutions to the renormalization group equations for running coupling constants. Correspondence is established between "eigenvalue" and induced-potential approaches toward total asymptotic freedom.

Consider a gauge theory of scalar (ϕ) and vector bosons. Perturbative renormalizability implies the existence of a term in the scalar field potential in order to cancel the logarithmic infinity induced by the gauge interaction (Fig.1). However, it has been recently demonstrated that self-energy and vertex insertions in a given scalar-amplitude graph (such as those in Fig. 1) are mimicked by replacing $g_{renormalized}$ with the running coupling constant $\bar{g}$.[1)] Upon performing this substitution, one obtains a renormalization-group-improved perturbation theory in which the four-scalar amplitudes of Fig.1 are no longer infinite. Therefore a ϕ^4 term in the scalar field potential is no longer required to ensure renormalizability.[1)]

The idea of dispensing with primitive quartic scalar couplings in the potential is most appealing, for such couplings are not asymptotically free in theories of physical interest.[2)] The only ϕ^4 terms that occur in a theory free of primitive ϕ^4 couplings are those induced by the gauge interaction and by Yukawa interactions. In Ref.1 it is suggested that these induced couplings, rendered finite by incorporation of running coupling constants, can generate spontaneous symmetry breaking. We examine these issues in greater detail below.

Consider the case of a scalar field ϕ_i transforming as a triplet under SU(2) $[\Phi^2 \equiv \phi_1^2 + \phi_2^2 + \phi_3^2]$. We choose to work in Landau-gauge in order to eliminate trilinear scalar-gauge-gauge-couplings to zero-momentum scalar field lines (such as in Figs. 1a and 1b). To one loop order, graphs of the form of Fig. 1c induce non-zero Φ^4 couplings through the quadrilinear scalar-scalar-gauge-gauge field interaction

$$\begin{aligned} \Delta \mathcal{L}_g &= \frac{g^2}{2}\left[\epsilon_{ijk}\,\phi_j W_k\right]\left[\epsilon_{ilm}\,\phi_l W_m\right] \\ &= \frac{g^2}{2}\,\Phi^2\left(W_I^2 + W_{II}^2\right) \equiv \frac{1}{2} M_{ab}^2(\Phi)\, W_a W_b \,. \end{aligned} \tag{1}$$

[W_I and W_{II} are orthonormal to the uncoupled gauge field $W_{III} \equiv (\phi_1 W_1 + \phi_2 W_2 + \phi_3 W_3)/\Phi$, where $W_{1,2,3}$ correspond to τ-matrices.] In Fig. 1c, W_I and W_{II} exchanges lead to equal contributions to the induced Φ^4 potential. The sum of these contributions is[3),4)]

$$\begin{aligned} \Delta V_g &= 3i\left(\frac{1}{4}\right)\int \frac{d^4k}{(2\pi)^4}\, Tr\left[\frac{M^4(\Phi)}{(k^2 - m_W^2 + i\epsilon)^2}\right] \\ &= \frac{3i}{2}\left[\int \frac{d^4k}{(2\pi)^4}\, \frac{\bar{g}^4(-k^2)}{(k^2 - m_W^2 + i\epsilon)^2}\right] . \end{aligned} \tag{2}$$

The bracketed integral is manifestly UV-finite. In Ref.1 arguments are presented invoking a spectral representation of $\bar{g}$ to show that the bracketed integral is IR-finite as well.[5)] Upon Wick rotation to Euclidean $\mathcal{k}$ $[k^2 = (k_0^2 - \vec{k}^2) \to -(k_4^2 + \vec{k}^2) \equiv -\mathcal{k}^2;\; d^4k \to i\pi^2 \mathcal{k}^2 d\mathcal{k}^2]$ we find that

$$\begin{aligned} \Delta V_g &= -\frac{3\Phi^4}{32\pi^2}\left[\int_0^\infty \frac{\mathcal{k}^2 d\mathcal{k}^2\; \bar{g}^4(+\mathcal{k}^2)}{(\mathcal{k}^2 + m_W^2)^2}\right] \\ &\equiv \frac{-3}{32\pi^2}\, \Phi^4\, I[\bar{g}, m_W] \,. \end{aligned} \tag{3}$$

Since the integrand of $I[\bar{g}, M_W]$ is positive, ΔV_g is negative. In fact, ΔV_g is negative for any gauge model, as $Tr[M^4]$ is always greater than zero. Had we not substituted $\bar{g}$ for g, ΔV_g would have corresponded to a negative infinity to be absorbed in primitive $\lambda\Phi^4$ terms. If ΔV_g is rendered finite by substitution of $\bar{g}$ for g, a primitive $\lambda\Phi^4$ term is no longer required to absorb infinities. However, in the absence of such a term, the induced coefficient of Φ^4 is negative, and spontaneous symmetry breaking is impossible.

This situation can be remedied by introducing a sufficient number of fermions into the theory. The graphs shown in Fig. 2 induce positive Φ^4 couplings that are opposite in sign to those of Fig. 1c because of the closed fermion loops. It is non-trivial, however, to change the sign of the induced potential by adding fermions to the theory. Suppose, for example, that we almost saturate asymptotic freedom of the SU(2) gauge coupling by incorporating 10 fermion doublets f^A in the theory. For simplicity, we assume that there is only one free Yukawa parameter h and a Yukawa interaction given by

$$\Delta\mathcal{L}_Y = -h \sum_{A=1}^{10} \bar{f}_\alpha^{\,A} (\tau^a/2)_{\alpha\beta} f_\beta^{\,A} \phi^a \tag{4}$$

This Yukawa interaction can be expressed as

$$\Delta\mathcal{L}_Y = -\bar{F}_I \; m_{IJ}(\phi_i) F_J \quad , \tag{5}$$

where F_I is a column vector containing all 20 fermions.
The induced Φ^4 coupling from graphs in Fig. 2 is given by[3]

$$\Delta V_F = -i\left(\tfrac{1}{4}\right)\int \frac{d^4k}{(2\pi)^4} \; Tr\left[\frac{m^4(\phi_i)}{(k^2 - m_F^2 + i\epsilon)^2}\right], \tag{6}$$

where $m_{IJ}^2 = h^2 \Phi^2 \delta_{IJ}/4$, and where the trace is understood to be over Dirac indices as well[3].

ΔV_F will be finite provided h is replaced by the running Yukawa coupling constant $\bar{h}(-k^2)$:

$$\begin{aligned} \Delta V_F &= \frac{\Phi^4}{16\pi^2}\left(\frac{5}{4}\right)\int_0^\infty \frac{k^2 dk^2\, \bar{h}^4(+k^2)}{[k^2 + m_F^2]^2} \\ &= \frac{\Phi^4}{16\pi^2}\left(\frac{5}{4}\right) I[\bar{h}, m_F]. \end{aligned} \tag{7}$$

We now wish to compare the magnitudes of ΔV_F and ΔV_g. In Appendix I the renormalization group equations of $\bar{g}$ and $\bar{h}$ are shown to be of the form

$$\begin{aligned} \frac{1}{2}\frac{d\bar{g}^2}{dt} &= -b\,\bar{g}^4 \\ \frac{1}{2}\frac{d\bar{h}^2}{dt} &= A\,\bar{h}^4 - B\,\bar{h}^2\bar{g}^2 \quad , \end{aligned} \tag{8}$$

where A,B and b are all greater than zero. The solution for $\bar{h}(t)$ may be written in the form

$$\bar{h}^2(t) = \bar{g}^2(t)\left[\frac{R\xi}{R + u(t)(\xi - R)}\right], \tag{9}$$

where $R \equiv \bar{h}^2(0)/\bar{g}^2(0)$, $\xi \equiv (B-b)/A$
and $u(t) = [\bar{g}^2(0)/\bar{g}^2(t)]^{(B-b)/b}$. The gauge coupling constant is asymptotically free provided $b > 0$.
In Appendix I, the Yukawa coupling is shown to be asymptotically free provided $\xi > 0$ and $0 \leq R \leq \xi$. Since the fermion mass m_F obtained through spontaneous symmetry breaking is related to the gauge boson mass by $m_F^2 = R m_w^2/4$,[6]

$$\begin{aligned} I[\bar{h}, m_F] &= \int_0^\infty k^2 dk^2\; \bar{h}^4(+k^2)/(k^2 + m_F^2)^2 \\ &= \int_0^\infty k^2 dk^2\; \bar{g}^4(+k^2)\left[\frac{R\xi}{[R + u(k^2)(\xi - R)][k^2 + R m_w^2/4]}\right]^2. \end{aligned} \tag{10}$$

The expression in brackets [Eq. (10)] increases monotonically as R approaches its maximum allowed value of ξ [note that $u(t) > 1$ for all t]. Therefore

$$I_{max}[\bar{h}, m_F] = 16\int_0^\infty k^2 dk^2\, \frac{\bar{g}^4(+k^2)(\xi/4)^2}{(k^2 + \xi m_w^2/4)^2}. \tag{11}$$

Since $\xi/4 < 1$ in our example (Appendix I), we see that

$$1/(k^2+m_w^2) > (\xi/4)/(k^2+\xi m_w^2/4) > (\xi/4)/(k^2+m_w^2), \qquad (12)$$

in which case

$$16\, I[\bar g, m_w] > I_{max}[\bar h, m_F] > \xi^2 I[\bar g, m_w]. \qquad (13)$$

Therefore the total induced potential $V^{ind} = \Delta V_g + \Delta V_F$ has a maximum possible value in the range

$$(37/32\pi^2)\, I[\bar g, m_w] > V^{ind} > [(-3/2 + 5\xi^2/4)/16\pi^2]\, I[\bar g, m_w]. \qquad (14)$$

Since the factor $(-3/2 + 5\xi^2/4)$ is negative in our example, we have failed to prove that V^{ind} is necessarily positive, despite the fact that the addition of even one more fermion doublet destroys asymptotic freedom for the gauge coupling constant.[7]

Under what circumstances will the induced potential be guaranteed to be positive? Suppose the SU(2) model we have been considering contained a primitive scalar field potential $\lambda\Phi^4/4!$. The renormalization group equation for $\bar\lambda$ is found to take the form[8]

$$16\pi^2 \dot{\bar\lambda} = \alpha\bar\lambda^2 - \beta\bar\lambda\bar g^2 + \gamma\bar\lambda\bar h^2 + \rho\bar g^4 - \sigma\bar h^4 \qquad (15)$$

[$\alpha, \beta, \gamma, \rho, \sigma$ are all positive]. If $\bar h^2 = \xi\bar g^2$, its maximum allowed value, and if $(\rho - \xi^2\sigma) < 0$, then $d\bar\lambda/dt$ is negative for arbitrarily small positive values of $\bar\lambda$,

and a "positive-eigenvalue" asymptotically free solution $\bar\lambda(t) = k\bar g^2(t)$ [$k>0$] is guaranteed to exist.[9] Note that the coefficients ρ and σ are proportional to the coefficients of the divergent part of diagrams in Figs. 1c and 2, respectively. In our example, $\rho = 72$, $\sigma = 60$, and in which case $(\rho - \xi^2\sigma) = 48(3/2 - 5\xi^2/4) > 0$:[8] a positive-eigenvalue solution is *not* guaranteed to exist. The point we wish to make is that the criterion that would guarantee a positive induced potential [$(-3/2 + 5\xi^2/4) < 0$] is *the same* as the criterion guaranteeing a positive-eigenvalue asymptotically free solution for $\bar\lambda$.[10] This is hardly surprising, as the same combinatorics manifestly appear from Figs. 1c and 2. In terms of the notation of Ref.8, in which scalar fields ϕ_i transform according to

$$D_\mu\phi_i = \partial_\mu\phi_i + ig\,\theta^a_{ij}\phi_j W^a_\mu \qquad (16)$$

and the Yukawa coupling to the set of fermion fields F is given by

$$\Delta\mathcal{L}_Y = -\bar F h_i F\phi_i , \qquad (17)$$

a sufficient condition for the existence of a positive induced coefficient of $\phi_i\phi_j\phi_k\phi_l$ is that[11]

$$3A_{ijkl} - 48\mathcal{H}_{ijkl} < 0, \qquad (18)$$

where

$$
\begin{aligned}
A_{ijkl} = & \{\theta^a, \theta^b\}_{ij}\{\theta^a, \theta^b\}_{kl} \\
& + \{\theta^a, \theta^b\}_{ik}\{\theta^a, \theta^b\}_{jl} \\
& + \{\theta^a, \theta^b\}_{il}\{\theta^a, \theta^b\}_{jk} ,
\end{aligned}
\tag{19}
$$

and where the matrix

$$
\begin{aligned}
\mathcal{H}_{ijkl} = \frac{1}{6} \operatorname{Tr} \big[& \rho_i \rho_j \{\rho_k, \rho_l\} \\
& + \rho_i \rho_k \{\rho_j, \rho_l\} \\
& + \rho_i \rho_l \{\rho_j, \rho_k\} \big]
\end{aligned}
\tag{20}
$$

is expressed in terms of Yukawa-coupling-constant eigenvalue solutions $\bar{h}(t) \equiv \rho_i \, \bar{g}(t)$.

The criterion of Eq. (18) is also sufficient for the existence of a positive-eigenvalue solution for the running coefficient of $\phi_i \phi_j \phi_k \phi_l$ within any primitive potential present in the theory.[11] We claim, however, that the real significance of theories containing such positive-eigenvalue solutions, which can hardly be expected to be radiatively stable, is that such theories are also likely to contain positive induced potentials after running coupling constants have been incorporated to allow removal of all primitive Φ^4 couplings.

There is an important qualitative difference between the positive-eigenvalue and the induced potential approaches to total asymptotic freedom; only the latter approach is successful for a continuous domain of input parameters - the

initial values R for the Yukawa couplings. Suppose, for example, that $(\rho - \xi^2 \sigma) < 0$, thereby guaranteeing that Eq. (15) has a positive eigenvalue solution for $\bar{\lambda}$ when the Yukawa coupling h is at its limiting eigenvalue solution $\bar{h}^2(t) = \rho \bar{g}^2(t)$. If $R = \bar{h}^2(0)/\bar{g}^2(0)$ is chosen to be anything less than ρ, the eigenvalue solution disappears, as $\bar{h}(t)$ now falls faster with t than $\bar{g}^2(t)$ [Eq. (9)].[8] The corresponding induced Φ^4-potential is positive provided $I[\bar{h}, m_F] > \rho^2 I[\bar{g}, m_W]$. However, we have already shown that this is true when $R = \rho$ [Eq. (13)], and $I[\bar{h}, m_F]$ is manifestly a continuous function of R[Eq. (10)].[12] Therefore, if the difference between $I_{max}[\bar{h}, m_F]$ and $\rho^2 I[\bar{g}, m_W]$ is finite, there will exist some domain of $R < \rho$ for which $I[\bar{h}, m_F]$ remains greater than $\rho^2 I[\bar{g}, m_W]$. The induced potential remains positive over this domain of R. Thus the presence of positive-eigenvalue solutions to a given theory serves to indicate the possible existence of an asymptotically free theory <u>of a non-eigenvalue character</u> in which positive quartic scalar field couplings are entirely induced.

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Appendix I

The Running Yukawa Coupling Constant

The Yukawa interaction between f fermion doublets and one fermion triplet ψ^A may be expressed as[8]

$$\Delta\mathcal{L} = -h_{AB}\,\bar\psi^A_\alpha\,(\tau^a/2)_{\alpha\beta}\,\psi^B_\beta\,\phi^a . \qquad (A.1)$$

In this equation there are f(f + 1)/2 independent h_{AB}'s. The associated running coupling constants satisfy [8]

$$16\pi^2 d\bar h_{AB}/dt = \frac{1}{4}\bar h_{AC}\bar h_{CD}\bar h_{CB} + \bar h_{IJ}\bar h_{JI}\bar h_{AB} - \frac{9}{2}\bar h_{AB}\bar g^2 . \qquad (A.2)$$

In the text we consider the case of a single independent Yukawa parameter such that

$$h_{AB} = h\,\delta_{AB} . \qquad (A.3)$$

The running value of h is then seen to satisfy the following differential equation:

$$\frac{1}{2}\frac{d\bar h^2}{dt} = A\bar h^4 - B\bar h^2\bar g^2 , \qquad (A.4)$$

where $A = [(1+4f)/4]/16\pi^2$ and $B = (9/2)/16\pi^2$.

The running gauge coupling constant equation is [7]

$$\frac{1}{2}\frac{d\bar g^2}{dt} = -b\bar g^4 , \qquad (A.5)$$

where $b = (21-2f)/48\pi^2$.

The solution to Eqs. (A.4,5) is found by setting $\bar h^2(t) = k(t)\bar g^2(t)$, in which case [8]

$$\frac{1}{2\bar g^2}\frac{dk}{dt} = Ak^2 - (B-b)k . \qquad (A.6)$$

Eq. (A.6) is integrated using Eq. (A.5):

$$\int_{k(0)}^{k(t)}\frac{dk}{k[Ak-(B-b)]} = 2\int_0^t \bar g^2(t')dt' = -\frac{1}{b}\log\frac{\bar g^2(t)}{\bar g^2(0)} , \qquad (A.7)$$

and the solution

$$k(t) = R\rho/[R + v(t)(\rho - R)] \qquad (A.8)$$

is obtained in which $v(t) = [\bar g^2(0)/\bar g^2(t)]^{(B-b)/b}$, $\rho = (B-b)/A$, and $R = k(0) = \bar h^2(0)/\bar g^2(0)$.

The gauge coupling is asymptotically free provided $b > 0$, in which case $f \leq 10$. The Yukawa coupling is asymptotically free provided $\lim_{t\to\infty}\bar h(t) \to 0$. A sufficient condition for this to be the case is that $\lim_{t\to\infty} k(t) \to 0$. Consider the RHS of Eq. (A.6):

$$(2A\bar g^2)^{-1}\dot k = k(k-\rho) \equiv F(k) . \qquad (A.9)$$

Eq. (A.9) has critical points at k=0 and $k=\rho$. The criterion for stability is that $(dF/dk|_{k_{crit}}) < 0$. Now $(dF/dk|_{k=0}) = -\rho$ and $(dF/dk|_{k=\rho}) = +\rho$. Suppose $\rho < 0$. Since $k \equiv \bar h^2/\bar g^2$ is positive, the $k=\rho$ critical point is unphysical. Therefore, physical values of k have stability properties determined by the k=0 critical point, which is unstable. However, if $\rho > 0$, the k=0 critical point is stable. k(t) will go to zero as $t \to \infty$ provided $0 \leq k(t) < \rho$. Therefore k(t) is asymptotically free provided $\rho > 0$ and $0 \leq R < \rho$. The case $R=\rho$ corresponds to k being constant over all t. Note for the example of SU(2) that $\rho = (8f-30)/(12f+3)$. Therefore to have an asymptotically free Yukawa coupling, f must be greater than or equal to 4.

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2) V. Elias, Phys. Rev. D20, 262 (1979).

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4) Our combinatoric and sign conventions for diagrammatic contributions to the potential are consistent with those of Ref. 3. Since the primitive potential $\lambda\phi^4/4!$ corresponds to a vertex of $-i\lambda$, an "induced vertex" $+\Lambda$ corresponds to an induced potential $i\Lambda\phi^4/4!$; we determine the sign of the induced potential by multiplying the diagrammatic coefficient of ϕ^4 by $+i$. The sign of the integral in Eq. (2) is opposite to the sign appearing in Eq. (11) of Ref. 1, which is in error.

5) The form of $\bar{g}(-k^2)$ in the IR region is the subject of some speculation -- see R. Anishetty, M. Baker, S.K. Kim, J.S. Bell and F. Zachariesen, Univ. of Washington preprint RLO-1388-789.

6) The existence of a suitably defined renormalization point μ such that $t = \log k/\mu$ is implicit throughout. Thus $m_W^2 = \bar{g}^2(0)\langle\phi\rangle^2$ and $m_F^2 = \bar{h}^2(0)\langle\phi\rangle^2/4$.

7) D.J. Gross and F. Wilczek, Phys. Rev. D8, 3633 (1973).

8) T.P. Cheng, E. Eichten, and L.-F. Li, Phys. Rev. D9, 2259 (1974).

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10) Note that a negative diagrammatic ϕ^4 coefficient to the divergent part of an integral corresponds to a positive term in the appropriate renormalization group equation β-function. These signs are sorted out properly in D. Bailin, Weak Interactions, (Sussex University Press, London, 1977), pp. 345-350.

11) See Eq. (2.8) of Ref.8, as corrected in footnote 9 of Chang[Op.Cit Ref.9].

12) $I[k, m_F]$ is displayed as a function of R in Eq. (10). If $I[q, m_F]$ is finite, then $dI[k, m_F]/dR$ is also finite for $0<R\leq\xi$. Hence $I[k, m_F]$ is a continuous function of R over this domain.

Figure Captions

Figure 1

One loop contributions to the ϕ^4-potential induced by interaction with the gauge field multiplet.

Figure 2

One loop contributions to the ϕ^4-potential induced by the Yukawa interaction.

FIG.1a

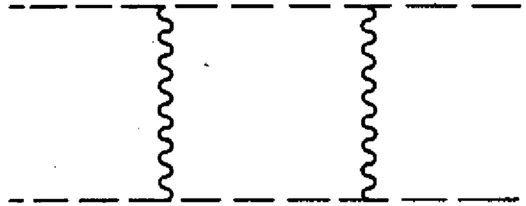


FIG.1b

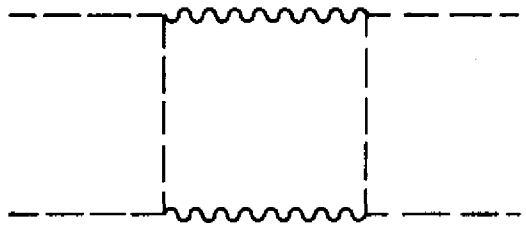


FIG.1c

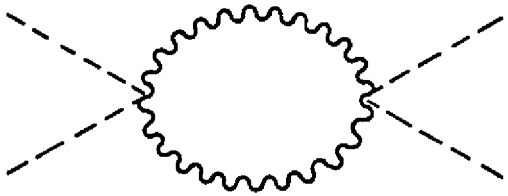
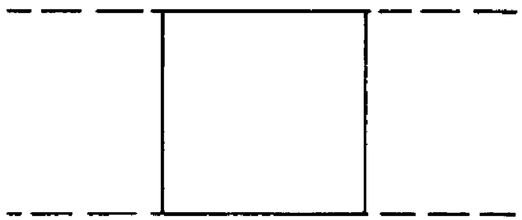


FIG.2



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