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A Soluble Model of "Higgs boson" as a Composite*†

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Abstract

Higgs boson may turn out to be a composite. The theoretical description of such a composite is illustrated by an example of a soluble model.

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1. Introduction

The $\sigma$-model [1] is a successful phenomenological theory of low energy particle physics. Yet, the $0^+ \sigma$-particle itself has never been identified experimentally [2]. One of the possible reasons for this failure might be that the spin-parity transformation of $\sigma$ is the same as that of $\sigma^2$, $\sigma^3$, $\cdots$, $\sigma^n$, $\cdots$. (1.1)

Thus, what is a $\sigma$-field in the idealized theoretical model may appear experimentally as a "composite" due to the possible mixture of (1.1). Today, a major focus of high energy physics is to search for Higgs boson. It might be that for similar and other reasons, Higgs boson [3] could also turn out to be a "composite" [4-8]. Neither the experimental identification of $\sigma$ nor that of Higgs boson would correspond to the usual simple theoretical description of a single pole in the complex energy plane. In this paper, we explore the theoretical description of such a composite from a more elementary perspective, by examining the generalization of a soluble model [9]. The structure of the model is given in Section 2, and its solution in Sections 3 and 4.

This paper is dedicated to the memory of Gunard Källen, who (besides others) has made important contributions [10-15] to the original soluble model.

2. A Generalized Soluble Model

We generalize the original $V \rightleftharpoons N\theta$ model by retaining the same fixed Fermion states $V$ and $N$, but replacing the single $\theta(r)$ field by three boson fields $A(r)$, $B(r)$ and $C(r)$. The Hamiltonian $H$ in the new model is given by

$$H = H_0 + H_1 + H_2$$ (2.1)
where
\[ H_0 = m_0 V^\dagger V + \sum_k (\lambda_k a_k^\dagger a_k + \mu_k b_k^\dagger b_k + \nu_k c_k^\dagger c_k). \] (2.2)

For convenience, the entire system is enclosed within a sphere of large radius \( R \). The \( s \)-wave part of the annihilation field operators \( A(r), B(r) \) and \( C(r) \) are given in terms of their annihilation operators \( a_k, b_k \) and \( c_k \) by

\[
A(r) = \sum_k (4\pi R\lambda_k)^{-\frac{1}{2}} u_k r^{-1} (\sin kr) a_k,
\]
\[
B(r) = \sum_k (4\pi R\mu_k)^{-\frac{1}{2}} v_k r^{-1} (\sin kr) b_k
\]

and

\[
C(r) = \sum_k (4\pi R\nu_k)^{-\frac{1}{2}} w_k r^{-1} (\sin kr) c_k
\]

with
\[
\lambda_k = (k^2 + \alpha^2)^{\frac{1}{2}},
\]
\[
\mu_k = (k^2 + \beta^2)^{\frac{1}{2}},
\]
\[
\nu_k = (k^2 + \gamma^2)^{\frac{1}{2}}
\]

and \( \alpha, \beta, \gamma \) the masses of bosons \( a, b \) and \( c \). The functions \( u_k, v_k, w_k \) are convergence factors, which may all be chosen to be 1 for \( k < k_{max} \) and 0 otherwise. In (2.3) all summations extend over

\[
k = n\pi/R
\]

with \( n = 1, 2, 3, \cdots \). At equal time, we have the anti-commutation relations

\[
\{V, V^\dagger\} = \{N, N^\dagger\} = 1
\]

and the commutation relations

\[
[a_k, a_{k'}^\dagger] = [b_k, b_{k'}^\dagger] = [c_k, c_{k'}^\dagger] = \delta_{kk'}.
\] (2.7)
If one wishes, (2.6) can also be changed into commutation relations, and $V$ and $N$ would then be bosons.

In (2.2), we set the mass of $N$ to be zero, and the "bare" mass of $V$ to be $m_0$. The interaction Hamiltonians $H_1$ and $H_2$ are given by

$$H_1 = g (V \dagger NC(0) + N \dagger VC \dagger (0)) \quad (2.8)$$

and

$$H_2 = f (V \dagger NB \dagger (0) A(0) + N \dagger VA \dagger (0) B(0)) \quad (2.9)$$

The $g$-coupling governs the transition

$$V \rightleftharpoons Nc \quad (2.10)$$

and the $f$-coupling gives rise to the scattering

$$Na \rightleftharpoons Vb \quad (2.11)$$

Thus, when $f = 0$ the Hamiltonian is identical to that of the original $V \rightleftharpoons N\theta$ model, with $\theta$ replaced by $c$.

Throughout the paper, we assume $V$ to be unstable through $V \rightarrow Nc$ when $R \rightarrow \infty$. Since the mass of $N$ is set to be zero in the model, $V$ is unstable if its physical mass $m$ is larger than $\gamma$, the mass of $c$; i.e.,

$$m > \gamma \quad (2.12)$$

Thus, in a collision of $Na$, beside the elastic scattering

$$Na \rightarrow Na \quad (2.13)$$

we also have the inelastic process

$$Na \rightarrow Vb \rightarrow Nbc \quad (2.14)$$
provided that the total energy $E$ satisfies

$$E > \beta + \gamma ,$$  \hspace{1cm} (2.15)

the threshold energy of the channel $Nbc$.

We shall assume

$$\alpha < \beta + \gamma \quad (2.16)$$

Hence, in the $Na$ channel at low energy when

$$\alpha < E < \beta + \gamma$$ \hspace{1cm} (2.17)

there is only the elastic scattering (2.13); at higher energy when $E > \beta + \gamma$, we have both (2.13) and the inelastic process (2.14).

Consider first the process

$$Nc \leftrightarrow V \leftrightarrow Nc \quad .$$ \hspace{1cm} (2.18)

Denote the corresponding state vector by

$$|Nc\rangle \propto \left[ V^\dagger + g(4\pi R)^{-\frac{1}{2}} \sum_k \nu_k^{-\frac{1}{2}} k w_k (E - \nu_k)^{-1} c_k^\dagger \right] |vac\rangle .$$ \hspace{1cm} (2.19)

One can readily verify that it satisfies

$$H |Nc\rangle = E |Nc\rangle .$$ \hspace{1cm} (2.20)

At a finite $R$, $E$ satisfies the eigenvalue equation

$$h_R(E) \equiv E - m_0 - g^2 \sum_k \frac{k^2 w_k^2}{4\pi \nu_k R} \left( \frac{1}{E - \nu_k} \right) = 0$$ \hspace{1cm} (2.21)

with its derivative

$$h_R'(E) = 1 + g^2 \sum_k \frac{k^2 w_k^2}{4\pi \nu_k R} \left( \frac{1}{E - \nu_k} \right)^2 .$$ \hspace{1cm} (2.22)
always positive.

When $R \to \infty$, $h_R(E)$ becomes

$$h_{\infty}(E) = E - m_0 - g^2 \int_0^\infty \frac{k^2 w_k^2}{4\pi^2 \nu_k (E - \nu_k)} \frac{dk}{E - \nu_k}.$$  \hspace{1cm} (2.23)

The condition for $V$ being unstable is that when $E = \gamma$,

$$h_{\infty}(\gamma) < 0.$$  \hspace{1cm} (2.24)

In this case, we introduce a cut along the real axis from

$$E = \gamma \text{ to } \infty$$  \hspace{1cm} (2.25)

where $\gamma$ is the mass of the $c$-meson. The derivative of $h_{\infty}(E)$ is

$$h'_{\infty}(E) = 1 + g^2 \int_0^\infty \frac{k^2 w_k^2}{4\pi^2 \nu_k (E - \nu_k)^2}.$$  \hspace{1cm} (2.26)

which is positive $\geq 1$, when $E$ is real $< \gamma$. For $E = \nu_k > \gamma$ but just above the cut along the real axis, we have from (2.23)

$$\text{Im} h_{\infty}(\nu_k + io+) = i \left( \frac{g^2}{4\pi} \right) k w_k^2.$$  \hspace{1cm} (2.27)

Thus, on the second sheet near and below the cut (2.25), there is a zero of $h_{\infty}(E)$, which corresponds to the $V$-resonance. It can be shown that the phase shift $\delta$ for $Nc$ scattering (2.18) is related to $h_{\infty}(\nu_k + io+)$ and its complex conjugate by

$$e^{-2i\delta} = \frac{h_{\infty}(\nu_k + io+)}{[h_{\infty}(\nu_k + io+)]^*}.$$  \hspace{1cm} (2.28)

3. Na Sector (General Discussion)

As discussed in the Introduction, assume the idealized case that there does exist a fundamental spin 0 field $\phi$ which is the origin of masses of spin nonzero
particles. In any physical process, there are bound to be effective couplings between $\phi$ and some of its higher power products, such as
\[ \phi^2, \phi^3, \ldots, \phi^n, \ldots \]
Thus, the physical Higgs channel becomes connected to not only a complex pole, but also to a cut in the complex energy plane, or other more complicated analytical structure.

In this section, the $Na$ channel that we shall analyze represents a highly idealized model of "Higgs" as a composite. From reactions (2.10) and (2.11), we see that a state vector $| \rangle$ in the $Na$ sector must also have components in $Vb$ and $Nbc$ channels as well. Thus for $R$ finite, we may write
\[
| \rangle = \left[ \sum_k \psi(k) a_k^\dagger N_k^\dagger + \sum_p \phi(p) b_p^\dagger V_p^\dagger + \sum_{p,q} \chi(p,q) b_p^\dagger c_q^\dagger N_q^\dagger \right] |0 \rangle .
\]
(3.1)
From
\[ H| \rangle = E| \rangle , \]
(3.2)
we find
\[ (E - \lambda_k) \psi(k) = f U_k \sum_p V_p \phi(p) , \]
(3.3)
\[ (E - m_0 - \mu_p) \phi(p) = g \sum_q W_q \chi(p,q) + f V_p \sum_k U_k \psi(k) \]
(3.4)
and
\[ (E - \mu_p - \nu_q) \chi(p,q) = g W_q \phi(q) \]
(3.5)
where $k, p, q$ are all given by (2.5) and $U_k, V_p, W_q$ are related to the $u_k, v_p$ and $w_q$ of (2.3) by
\[ U_k = (4\pi R \lambda_k)^{-\frac{1}{2}} ku_k \]
(3.6)
\[ V_p = (4\pi R \mu_p)^{-\frac{1}{2}} pv_p \]
and
\[ W_q = (4\pi R \nu_q)^{-\frac{1}{2}} qw_q \]
Substituting (3.5) into (3.4), we have

\[
(E - m_0 - \mu_p)\phi(p) = g^2 \sum_q W_q^2 (E - \mu_p - \nu_q)^{-1} \phi(p) + fV_p \sum_k U_k \psi(k) \tag{3.7}
\]

and therefore

\[
\phi(p) = [D_p(E)]^{-1} fV_p \sum_k U_k \psi(k) \tag{3.8}
\]

where

\[
D_p(E) = E - m_0 - \mu_p - g^2 \sum_q W_q^2 (E - \mu_p - \nu_q)^{-1} . \tag{3.9}
\]

From (3.8), we also have

\[
\sum_p V_p \phi(p) = f \left[ \sum_p \frac{V_p^2}{D_p(E)} \right] \sum_k U_k \psi(k) .
\]

Thus, (3.3) becomes

\[
(E - \lambda_k)\psi(k) = f^2 U_k \left[ \sum_p \frac{V_p^2}{D_p(E)} \right] \sum_{k'} U_{k'} \psi(k') . \tag{3.10}
\]

Multiplying both sides by \(U_k/(E - \lambda_k)\) and summing over \(k\), we find that the eigenvalue \(E\) satisfies

\[
1 = f^2 F(E) \sum_p \frac{V_p^2}{D_p(E)} . \tag{3.11}
\]

in which

\[
F(E) = \sum_k U_k^2 (E - \lambda_k)^{-1} . \tag{3.12}
\]

Next, we study the continuum limit. When \(R \to \infty\), the sum

\[
U_k \sum_p V_p = (4\pi R)^{-1} \sum_p (\lambda_k \mu_p)^{-\frac{1}{2}} k u_p v_p . \tag{3.13}
\]

becomes

\[
(4\pi^2)^{-1} \int (\lambda_k \mu_p)^{-\frac{1}{2}} k u_p v p dp . \tag{3.14}
\]

Thus, from (3.3) we have

\[
(E - \lambda_k)\psi(k) = (4\pi^2)^{-1} f k u_k \int \frac{1}{2} v p \phi(p) dp . \tag{3.15}
\]
Likewise, (3.7) leads to

\[
(E - m_0 - \mu_p)\phi(p) = (4\pi^2)^{-1} g^2 \phi(p) \int \nu^{-1}_q (E - \mu_p - \nu_q)^{-1} q^2 w_q^2 dq \\
+ (4\pi^2)^{-1} f p v_p \int (\mu_p \lambda_k)^{-\frac{1}{2}} k u_k \psi(k) dk 
\]

(3.16)

which gives

\[
\phi(p) = (4\pi^2 \mathcal{D}_p(E))^{-1} f p v_p \int (\mu_p \lambda_k)^{-\frac{1}{2}} k u_k \psi(k) dk 
\]

(3.17)

where

\[
\mathcal{D}_p(E) = E - m_0 - \mu_p - (4\pi^2)^{-1} g^2 \int [\nu(E - \mu_p - \nu_q)]^{-1} q^2 w_q^2 dq 
\]

(3.18)

In a collision of Na, in order to describe reactions (2.13) and (2.14), we write \( \psi(k) \) and \( \phi(p) \) as

\[
\psi(k) = \delta(k - k_0) + \tilde{\psi}(k) 
\]

(3.19)

and

\[
\phi(p) = \tilde{\phi}(p) 
\]

(3.20)

in which \( \tilde{\psi}(k) \) and \( \tilde{\phi}(p) \) denote the scattered wave amplitudes. Thus, (3.15) remains valid if we replace \( \psi, \phi \) simply by \( \tilde{\psi} \) and \( \tilde{\phi} \). Hence

\[
(E - \lambda_k)\tilde{\psi}(k) = (4\pi^2)^{-1} f k u_k \int (\lambda_k \mu_p)^{-\frac{1}{2}} p v_p \tilde{\phi}(p) dp 
\]

(3.21)

On the other hand, (3.17) yields

\[
\tilde{\phi}(p) = (4\pi^2 \mathcal{D}_p(E))^{-1} f p v_p \left[ (\mu_p \lambda_0)^{-\frac{1}{2}} k u_0 + \int (\mu_p \lambda_k)^{-\frac{1}{2}} k u_k \tilde{\psi}(k) dk \right] 
\]

(3.22)

with

\[
\lambda_0 = \lambda_k \quad \text{and} \quad u_0 = u_k \quad \text{at} \quad k = k_0 
\]

(3.23)

Define

\[
\mathcal{A} = (4\pi^2)^{-1} \int (v_p p / \mu_p^2) \tilde{\phi}(p) dp 
\]

(3.24)
\[ \mathcal{B} = (4\pi^2)^{-1} \int (u_k k/\lambda_k^{\frac{1}{2}}) \tilde{\psi}(k) dk \] (3.25)

and

\[ \mathcal{C} = (4\pi^2)^{-1}(u_0 k_0/\lambda_0^{\frac{1}{2}}) \] . (3.26)

Hence, (3.21) and (3.22) can be written as

\[ \tilde{\psi}(k) = f \mathcal{A} u_k k/\lambda_k^{\frac{1}{2}}(E - \lambda_k) \] (3.27)

and

\[ \tilde{\phi}(p) = f (\mathcal{B} + \mathcal{C}) v_p p/\mu_p^{\frac{1}{2}} \mathcal{D}_p(E) \] (3.28)

with \( \mathcal{D}_p(E) \) given by (3.18). In above expressions, all integrations over \( p \) and \( k \) extend from 0 to \( \infty \).

Substituting (3.28) into (3.24), we find

\[ I \equiv \mathcal{A}/(\mathcal{B} + \mathcal{C}) \] (3.29)

is given by

\[ I = \frac{f}{4\pi^2} \int_0^\infty \frac{(v_p^2 p^2/\mu_p)dp}{E - m_0 - \mu_p - \frac{g^2}{4\pi^2} \int_0^\infty \frac{q^2 w_q^2 dq}{\nu_q (E - \mu_p - \nu_q)}} . \] (3.30)

Likewise,

\[ J \equiv \mathcal{B}/\mathcal{A} \] (3.31)

becomes

\[ J = \frac{f}{4\pi^2} \int_0^\infty \frac{(u_k^2 k^2/\lambda_k)dk}{E - \lambda_k} . \] (3.32)

Thus

\[ \mathcal{A} = \frac{I}{1 - IJ} \mathcal{C} \] (3.33)

and

\[ \mathcal{B} + \mathcal{C} = \frac{1}{1 - IJ} \mathcal{C} \] (3.34)
From (3.26), (3.30) and (3.32), we have the explicit expressions for $C$, $I$ and $J$. Hence $A$ and $B$ are also known. Equation (3.27) and (3.28) then give scattering amplitudes $\tilde{\psi}(k)$ and $\tilde{\phi}(p)$.

4. Na Sector (Critical $f^2$)

We shall show that when $R \to \infty$ and $f^2$ greater than a critical strength $f_c^2$, there exists a bound state in the Na sector. Write the $R \to \infty$ limit of (3.11) as

$$1 = f^2 \mathcal{F}(E)\mathcal{G}(E)$$ (4.1)

in which

$$\mathcal{F}(E) = \lim_{R \to \infty} F(E)$$

$$= (4\pi^2)^{-1} \int_{0}^{\infty} \left[ k^2 u_k^2 / \lambda_k(E - \lambda_k) \right] dk$$ (4.2)

with $u_k$, $\lambda_k$ given by (2.3) and (2.4). The function $\mathcal{G}(E)$ is similarly related to the last summation in (3.11) by

$$\mathcal{G}(E) = \lim_{R \to \infty} \sum_{p} \frac{V_p^2}{D_p(E)}$$

$$= (4\pi^2)^{-1} \int_{0}^{\infty} \left[ p^2 v_p^2 / \mu_p D_p(E) \right] dp$$ (4.3)

where $D_p(E)$ is given by (3.18), Since $D_p(E)$ is related to $h_{\infty}(E)$ of (2.23) by

$$D_p(E) = h_{\infty}(E - \mu_p)$$ (4.4)

we have from (2.24)

$$D_p(\mu_p + \gamma) = h_{\infty}(\gamma) < 0$$ (4.5)

From (2.4),

$$\nu_q = (q^2 + \gamma^2) > \gamma$$ (4.6)

Thus, (3.18) and (4.4) - (4.5) imply that $D_p(E)$ and its derivative

$$D'_p(E) = \frac{\partial}{\partial E} D_p(E)$$ (4.7)
are continuous and satisfy

\[ D_p(E) < 0 \quad \text{and} \quad D'_p(E) > 0 \] (4.8)

over the range

\[ E < \mu_p + \gamma \] ,

which includes the range

\[ E < \alpha \] (4.10)

in accordance with (2.4) and (2.16). Thus, both \( \mathcal{F}(E) \) and \( \mathcal{G}(E) \) are negative, with negative derivatives; their product is positive and varies from 0 to a finite value as \( E \) increases from \(-\infty\) to \( \alpha \), the mass of \( a \)-meson.

Define a critical \( f^2 \)-coupling by

\[ f_c^2 = [\mathcal{F}(\alpha)\mathcal{G}(\alpha)]^{-1} \] (4.11)

It then follows that there exists a bound state energy \( E_0 \) in the \( Na \) sector with

\[ 1 = f^2 \mathcal{F}(E_0)\mathcal{G}(E_0) \] (4.12)

provided

\[ f^2 > f_c^2 \] (4.13)

For \( f^2 < f_c^2 \), the state turns into a resonance with a complex \( E_0 \). In this case the scattering amplitudes \( \tilde{\psi}(k) \) and \( \tilde{\phi}(k) \) have besides the cuts given by (3.30) and (3.32), also a complex pole at \( E = E_0 \).

It is of interest to note the difference between this pole in the \( Na \) sector and the \( V \)-pole in the \( Nc \) sector. The \( V \)-pole becomes stable in the weak coupling limit when \( g^2 \to 0 \), whereas in the \( Na \) sector the bound state \( E_0 \) becomes stable only in the strong coupling region when \( f^2 > f_c^2 \).

We note that when \( g^2 = 0 \), \( \mathcal{G}(E) \) of (4.3) becomes

\[ \mathcal{G}_0(E) = (4\pi^2)^{-1} \int_0^\infty \left[ p^2 v^2_p / \mu_p(E - m_0 - \mu_p) \right] dp \] . (4.14)
Correspondingly, (4.12) becomes

\[ 1 = f^2 \mathcal{F}(E_0) \mathcal{G}_0(E_0) \]  

(4.15)

with the same \( \mathcal{F}(E) \) of (4.2). Thus, the existence of the pole at \( E = E_0 \) does not depend sensitively on \( g^2 \); instead, it is closely related to the second (and higher) order attractive potential between \( Na \) due to the \( f \)-coupling transitions

\[ Na \rightleftharpoons Vb \rightleftharpoons Na \]  

(4.16)

Its physical origin is quite different from the \( V \)-pole in the \( Nc \) channel of (2.18).
5. Remarks

Consider the case

\[ f^2 < f_c^2 \]  \hspace{1cm} (5.1)

The composite \( Vb \) is unstable, and may serve as a highly simplified model of either the \( \sigma \)-meson or the Higgs boson. Besides the elastic process (4.16), there is also the inelastic reaction

\[ Na \leftrightarrow Vb \leftrightarrow Nbc \]  \hspace{1cm} (5.2)

In order to detect the composite \( Vb \) as a resonance, we require in (4.12) the corresponding pole at

\[ E = E_0 \]  \hspace{1cm} (5.3)

to be not too far from the real axis; hence \( f^2 \) cannot be too small. The amplitude for the continuum background must then also be relatively large.

In any composite model, we may regard the amplitudes \( Na \) and \( Vb \) as the idealized representations of its low and high frequency components of the same composite state vector. A second order transition between these two components would always depress \( Na \) and elevate \( Vb \), as in (4.16). A resonance thus formed would require a strong coupling, and therefore also a large continuum background as in the model. This could be the reason why the \( \sigma \)-meson does not appear as a sharp resonance, and it might also be difficult to isolate the Higgs boson resonance.

We wish to thank N. Christ and E. Ponton for discussions.
Appendix

In the special case when
\[ w_k^2 = k/\nu_k^3 \quad , \]  
(A.1)
the integral
\[ F_\gamma(E) = \int_0^\infty \frac{k^2 w_k^2}{4\pi^2 \nu_k (E - \nu_k)} \, dk \]  
(A.2)
in (2.23) is given by
\[ F_\gamma(E) = \frac{1}{4\pi^2 \gamma} \left[ \frac{1}{z^3} (1 - z^2) \ln \frac{1}{1 - z} - \frac{1}{2z^2} (z + 2) \right] \]  
(A.3)
with
\[ z = E/\gamma \quad . \]  
(A.4)
At \( E = \nu_k + i0^+ \), we have
\[ \text{Im} F_\gamma(E) = -\frac{i}{4\pi \gamma} \left[ \frac{1}{z^3} (z^2 - 1) \right] \]  
(A.5)
in agreement with (2.27).
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[2] C. Amsler et al. (Particle Data G Group), Physics Letters B667, 1(2008) and 2009 partial update for the 2010 edition listed $f_0(600)$ of mass $400-1200$ $MeV$ and full width $600-1000$ $MeV$, which closely resembles the appearance of a composite, similar to $Na \rightarrow Nbc$ of (2.14), that will be analyzed in this paper.

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