SALAM-WEINBERG SYMMETRY BREAKING
WITH SUPERHEAVY HIGGS PARTICLES

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ABSTRACT

We discuss here the possibility of the breaking of the Salam-Weinberg symmetry by Higgs particles which are superheavy. The symmetry-breaking is associated with a nonzero vacuum expectation value of fermion condensates. This mechanism, if operative in nature, will imply the absence of Higgs particles at the weak scale.

At present Salam-Weinberg theory appears to be the best theory of electroweak unification, and is in complete agreement with the experiments. The conventional understanding here of the weak symmetry breaking is through Higgs mechanism, where the Higgs scalars are expected to be of low mass and are being looked for in the \( p\bar{p} \) collider and elsewhere. The dynamical mechanism merely consists of minimising a potential with respect to the possible vacuum expectation values (vevs) of the scalars. One then obtains that the masses of the scalars are of the same order as the vevs, and this is the reason for the search of light Higgs scalars. We wish to point out here, however, that an alternative mechanism exists where the Higgs particles could be superheavy. This will naturally imply the absence of observed Higgs particles at the weak scale directly or through radiative corrections and is thus relevant in the context of the present experimental search.

Before constructing explicit models, let us discuss the basic features of the dynamical mechanism involved here. Consider

\[
(\mathbb{1} + m_p^2) \varphi = g J_p,
\]

where \( \varphi \) stands for a Higgs scalar of mass \( m_\varphi \). Further, let \( \langle J_p \rangle \neq 0 \). We then obtain from equation (1) that effectively we shall have \( \langle \varphi \rangle = g \langle J_p \rangle / m_p^2 \). Thus \( m_\varphi \simeq 3 \times 10^{15} \text{GeV} \) and \( \langle \varphi \rangle = \mu \), with \( \mu \simeq 10^{11} \text{GeV} \) can yield \( \langle \varphi \rangle \simeq 10^{-2} \text{GeV} \). Here the following may be noted. We have needed a nonzero vev for the condensate. Further, the small value of \( \langle \varphi \rangle \) was due to the interplay of two large scales, one which gives a superheavy mass to the Higgs particles, and the other which generates the nonzero vev for the condensates.

We can have a nonzero vacuum expectation value of the condensate with a minimisation of the effective potential only when we have higher dimension operators. The theory thus will be nonrenorm-
We may have e.g. a supergravity theory which contains higher dimension operators, or, a gauge theory where the heavy modes have been integrated out, or both. We shall here base our illustration with an SU(5) supergravity model. The effective potential is really the negative of the Lagrangian, where the Lagrangian contains the higher dimension operators suppressed by Planck mass. The chiral supermultiplets taken here will be the 24-plet denoted by $\Sigma^2$ and the 5 and 5* denoted by $H^5$ and $H^{5*}$, respectively. We shall so arrange that the above Higgs particles will all be heavy, and ultimately examine the symmetry breaking mechanism as envisaged in equation (1). We may anticipate that e.g. \( \langle \Sigma^2_{\mu\nu} \rangle \neq 0 \) has the correct quintet structure to enable us to get a nonzero vacuum expectation value $\langle H^5 \rangle$ with requisite Salam-Weinberg quantum numbers.

We now proceed to construct the model. We take the SU(5) invariant superpotential as\(^3\)

\[
W = \lambda_1 \left[ \frac{1}{2} \text{tr} \Sigma^2 + \frac{1}{2} \text{tr} \Sigma^4 \right] + \lambda_2 H^5 \Sigma^2 + \lambda_3 M^2 H^5 \Sigma^2 + \lambda_4 m Z + C.
\]

As usual, $\Sigma^2$ breaks the SU(5) symmetry to the Salam-Weinberg symmetry and preserves global supersymmetry at the grand unification scale $M$. $Z$ is the Polonyi singlet\(^4\) and $m$ is the mass scale associated with the Polonyi sector. $C$ is the additive counter-term so as to make the cosmological constant vanish. The term $\lambda_2 H^5 \Sigma^2$ is taken in (2) so as to make all the Higgs fields superheavy. The scalar potential is given as\(^5\)

\[
V = \frac{1}{4} \varepsilon \left( \frac{1}{2} W + \frac{1}{2} \text{tr} \Sigma^2 \right) + \frac{1}{4} \left( \frac{1}{2} W \right)^2 + \frac{1}{4} \left( \frac{1}{2} W \right)^2 - \frac{1}{2} \phi^2 / W^2 \right] + \frac{1}{2} \text{tr} B^2,
\]

where $E=\exp(\sqrt{2}/\lambda_2)$ and $\phi'$ is any scalar corresponding to a chiral superfield and $B\lambda_2$ is the corresponding $B$ term. In the leading order this potential has a minimum when\(^6\)

\[
\langle \Sigma^2 \rangle = M \left( 2 \delta^2 - 5 \delta^2 \delta_5 \right), \quad (4)
\]

which breaks SU(5) symmetry. Using this solution, the masses of $H$ and $H'$ are given as $m_H^2 = 9M^2$ and $m_{H'}^2 = 4M^2$ where $a$ and $i$ stand for the SU(2)$_L$ and SU(3)$_C$ parts of SU(5). We shall now see that the fermion condensates which couple to the superheavy Higgs particles can yield the mechanism described above. For this purpose we replace the minimisation of $V$ by the extremisation of the Lagrangian which naturally includes possible nonzero vevs for the fermion condensates, similar to the attempts made elsewhere for supersymmetry breaking through gaugino condensates.\(^7\) We note that no superstrong interaction is envisaged; only a classical extremisation is here involved to obtain nonzero values for the condensates. For this purpose, the relevant expression for the fermionic part of the Lagrangian is given as\(^8\), with $i,j$ standing for all the chiral fields,

\[
\mathcal{L}_F = -\frac{1}{2} \text{tr} \left( \frac{1}{2} \phi^2 \phi' \right) \left[ W_{ij} + \frac{1}{2} \left( \phi^2 \phi' \right) W_{iijj} \right] + \frac{1}{2} \text{tr} \left( \frac{1}{2} \phi^2 \phi' \right) X_{ij} X_{ij} + \phi' \epsilon.
\]

From the above we take the SU(2)$_L$ doublet components of the fermion condensates with appropriate quantum numbers and write the expression on the right-hand side of equation (5) as

\[
\mathcal{L}_F = -\frac{1}{2} \text{tr} \left( \frac{1}{2} \phi^2 \phi' \right) \left[ H^4 J^4 + \frac{1}{2} \left( \phi^2 \phi' \right) X_{ij} X_{ij} + \phi' \epsilon. \right.
\]

where for the leading terms we have substituted

\[
J^4 = \lambda_2 \bar{X}_5 \gamma_5 X_H + \cdots + \lambda_2 \bar{X}_n X_H + \cdots \quad (7)
\]

and

\[
J'_4 = \lambda_1 \bar{X}_5 \gamma_5 X_H + \cdots + \lambda_2 \bar{X}_n X_H + \cdots \quad (8)
\]

We recall that the scalar fields $H^4$ and $H'_4$ are superheavy, and thus we shall eliminate them in favour of the condensates $J^4$ and $J'_4$. This is done with the obvious replacement

\[
H^4 \rightarrow \frac{1}{m_H^2} \exp(-\frac{1}{a} \phi^2 \phi'), \quad (9a)
\]
and
\[ R^4 \rightarrow \frac{g^4}{v^4} \exp(-\frac{1}{4} \phi^4 \phi). \]  

(9b)

In equations (7) and (8) we have explicitly written only the condensates with heavy fermion fields, so that no chiral symmetry breaking is involved. Through equations (9), which are another way of considering equation (1), we obtain a nonzero vev for the Higgs fields which is damped by the large Higgs mass. The mechanism here thus is completely different from that of technicolour. Further, we also have \[ \langle \overline{\psi}(x) \gamma^\mu \psi(x) \rangle = 0 \] and \[ \langle \overline{\psi}(x) \gamma^\mu \gamma^5 \psi(x) \rangle = 0, \] so that it is not possible to construct any mechanism to generate a mass for the gauge bosons directly through fermions without going through the Higgs particles.

We now make the substitutions \( f = \kappa M \) and \( f' = \kappa M \) and write the Lagrangian in a dimensionless form by multiplying by the appropriate powers of \( \kappa \). In fact, \( \kappa \) will be in the weak scale, so that we have \( f' = 10^{-5} \ll 1 \). We can thus use \( f' \) as an expansion parameter, while considering the vacuum expectation values of the Higgs fields as well as of the condensates. We then have the bosonic and the fermionic parts of the Lagrangian given as

\[ \mathcal{L} = \frac{1}{2} (\partial \Sigma)^2 + \frac{1}{4} \Sigma^2 + \frac{1}{4} Z^2 + \frac{1}{4} Z^2 \]

(10)

\[ + \frac{1}{2} \left( \lambda \phi^4 \phi^4 - \frac{3}{8} \right) \left[ \left( \lambda \phi^4 \phi^4 \right)^2 + \left( \lambda \phi^4 \phi^4 \right)^2 \right] \]

(11)

where

\[ \alpha = \exp(-\frac{1}{2} \phi^4 \phi). \]

and

\[ \mathcal{L}_F = \frac{1}{2} \left\{ \left( \lambda \phi^4 \phi^4 \right)^2 + \left( \lambda \phi^4 \phi^4 \right)^2 \right\}. \]

We now use equations (10) and (12) for the purpose of extremisation using \( f' \) as a small parameter. We thus include the \( f'^2 \) order contributions to the vacuum expectation values of \( \Sigma \) and \( Z \), in addition to the zeroth order contributions. For the condensates, we further substitute

\[ \frac{\langle \overline{\Sigma} (x) \rangle}{m^2} = \eta \phi + \ldots \]

(13)

and

\[ \frac{\langle \overline{Z} (x) \rangle}{m^2} = \eta \phi + \ldots. \]

(14)

We also expand the additive constant \( C = C_0 + C_2 \phi^2 \) adjusted so that the cosmological constant vanishes, and take \( \langle W \rangle = \frac{1}{2} \phi \)

so that gravitino has a mass in the weak scale. The expansion for the condensates has been taken as in equations (13) and (14) so that the right-hand sides exactly correspond to the conventional vacuum expectation values of the Higgs fields.

While considering the extremisation equations we find that for the consistency of the above type of solutions we need a fine tuning given by \( \delta \alpha \phi^2 / (9 \phi^2) = 1 \), which is a link up of the grand unification mass scale with a Yukawa coupling constant, and might indicate why the Yukawa coupling constants may be small. In this case we also obtain that \( \eta \phi \) in equation (14) is zero. \( Z_0 \) for the Polonyi singlet becomes independent of the coupling parameters and is given by \( Z_0^2 = 8 \pm 4 \). The possibility of a large value of \( Z_0 \) may be noted, which can make \( \alpha \) in equation (11) quite small. We also need a constraint \( \lambda^2 > \alpha^2 \phi^4 \), which is easily satisfied for a small \( \alpha \). In fact, with

\[ U = \lambda \eta \phi \frac{\phi^2 \phi^2}{\phi^2}, \]

we obtain \( m_q = \frac{\phi}{\sqrt{2}} \phi \) and \( m_z = m_\psi / \cos \theta_q \). Thus, all the results of the Salam-Weinberg theory remain unaltered except for the fact that there will be no Higgs particles at the weak scale. Since the Higgs particles are superheavy, their presence will also not be felt through radiative corrections.
We may next consider another model which is motivated by
superstrings. For such models, one starts with a ten-dimen-
sional manifold having the structure $M^4 \times K$, where $M^4$ is the
four-dimensional Minkowski space and $K$ is a compact six-
dimensional manifold. The mechanism of compactification
ensures N=1 supergravity in four-dimensions. Also, the resi-
dual gauge symmetry is expected to be $E_6$ or a subgroup of $E_6$
with the 27-plet yielding quarks and leptons along with some
other coloured particles. Witten calculated the N=1 super-
gravity content of the four-dimensional theory, which consists
of the following. The Kahler potential is given as, with $k=1$,

$$K = -\ln(S + S) - 3 \ln(T + T - 2 C \times C).$$

In the above, $S$ and $T$ are gauge singlet complex scalar fields
expressed in terms of the dilaton field $\varphi$, the scalar field
$\sigma$ arising from $\varphi^2$ in the metric in ten dimensions, the field
$Q$ arising from the two-form of the ten-dimensional supergra-
vity corresponding to the compactification dimensions, and the
field $D$ arising in writing the Minkowski part of the field
strength of ten dimensions as a gradient and an antisymmetric
Kronecker index. The gauge part of the Lagrangian is also
expressed through the function

$$f_{\alpha\beta} = \frac{1}{2} \delta_{\alpha\beta} S,$$

where $\alpha$ and $\beta$ represent the indices for the adjoint represent-
ation of the gauge group. The third element of supergravity, the
superpotential, is given by the invariant cubic expression in terms of the 27-plet matter field.

It may, however, be recognised that the mechanism of compacti-
fication is as yet not understood. Thus, the symmetry group
after compactification is often taken as an input assumption
for subsequent phenomenology. For our illustration we
shall adopt this philosophy here and consider the symmetry
group $SU(3) \times SU(2) \times SU(1)$, which is a rank five
subgroup of $E_6$. The 27-plet is here written as

$$27 \rightarrow (3,2,1/3,1/4)_Q + (1,2,-1/4)_Q + (1,2,1,-1/2)_Q +
(1,2,-1,1/2)_H + (3,1,2/3,1/2)_G + (3,1,1,-2/3,1/2)_C +
(3,1,2,1/4)_C + (1,1,0,1/4)_N + (1,1,0,1)_N.$$

We note that in the above $U(1)_R$ symmetry is broken along with
compactification. We take the broken form of the superpotential as

$$W_1 = \lambda_h H H Q + \lambda_d H D + \lambda_3 H E +
\lambda_{14} H H N + \lambda_g 3 N^2.$$

In addition, we also take a Polonyi-like superpotential

$$W_2 = m^2 (S + B_0),$$

where, as before, we assume $m^2 K = \rho^2$ to be small. With this,
the Lagrangian is constructed, which is obviously highly compi-
lcated. We next consider the nonzero vev structure of the scalar
fields as $<T> = t^L N^L$ along with a higher order vev
structure given as $<T^2> = t^L \rho^2, <T^2 \sigma^2> = \sigma^2 \rho^2$ and $<W> = \sigma^2 \rho^2$.

The last but one term above is the condensate vev. The Higgs
particles here have a mass of the Planck scale. The vev struc-
ture is so taken that the supersymmetry breaking is at the weak
scale. The calculations, which are involved, yield four equations
in the parameters $t, \rho, \sigma$, and $\rho / \sigma$. The extremisa-
tion has been done using the $\sigma^2$ as a small parameter. Peculiarly
only the ratios of some vacuum expectation values occur. Some
of the masses are written down in Ref.15. The coloured objects
$g$ and $g_0$ are seen to be superheavy. Similar symmetry groups
have also been considered for baryogenesis in cosmology.
In the above we have considered supergravity-based models which are naturally nonrenormalisable. As mentioned in the beginning, we could also have generated the higher dimensional operators from purely gauge theories which would be renormalisable, and followed the above mechanism for symmetry breaking. When we have fermion bilinear condensates, will the theory continue to be renormalisable? The answer to this question appears to be in the affirmative. For a moment let us consider renormalisability of spontaneously broken gauge theories when scalar Higgs fields have nonzero vevs. A simple way of doing this is to consider the vertex function of the spontaneously broken theory as an infinite sum of the vertex functions of the symmetric theory when there are extra external legs which escape to the vacuum. As stated, all such contributions are added along with appropriate renormalisation of the vevs. Then the renormalisability of the symmetric theory can be used to show the finiteness of the spontaneously broken theory. One can employ the same trick here when the vertex function of the spontaneously broken theory is expressed in terms of an infinity of vertex functions of the symmetric theory with extra pairs of fermion legs escaping to vacuum, and, again, use the renormalisability of the symmetric theory. Thus, if the original theory is renormalisable, then the resulting spontaneously broken symmetric theory with fermion condensates will also continue to be renormalisable.

The above feature has some phenomenological repercussions. Since the Higgs particles are superheavy, the slope of the curve for renormalisation group equations while going from the Salam-Weinberg scale to the grand unification scale will be larger, and thus proton will be more stable than otherwise envisaged. However, at the intermediate scale of condensates, given by \( \mu = 10^{11} \) GeV, new physics may be involved, the effect of which is not clear. Thus the nonobservation of the Higgs scalars at the weak scale may not only be consistent with Salam-Weinberg theory, but may also give rise to interesting theoretical possibilities. In contrast we may consider the renormalisability of superstring motivated supergravity models, which is to be traced back to the anticipated finiteness of the superstring phenomenology in ten dimensions. This probably will remain fuzzy as long as our understanding of compactification is incomplete.

The basic idea here is the interplay of the scale generating the Higgs mass with the scale generating the condensate vacuum expectation value to obtain the weak symmetry breaking scale. Thus it could be even relevant in preonic models where such multiple scales are present. Another remark that may be worthwhile is the possibility of generating intermediate symmetries starting from \( E_6 \) since we are considering such a small group as \( G_2 \). But, the theory with condensates and nontrivial Kahler metric is so complicated that such an approach appears to be difficult.

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References


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20. Ref. 15 and paper submitted for publication.