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Hidden Local Symmetries in Extended Abbott-Farhi Models

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Abstract

It is shown that the low-energy physics of an extended Abbott-Farhi model with two scalar doublets is described by a non-linear sigma model based on SP(4)/SU(2)xSU(2), which possesses an SU(2) gauge invariance as a hidden symmetry. This raises the interesting possibility of identifying the weak bosons observed at the collider experiments with the composite gauge bosons associated to such a hidden local symmetry.

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The standard theory of electroweak interactions has been very successful in describing physics of quarks and leptons up to $\sim$100 Gev. Nevertheless, there is still considerable effort expanded in constructing composite models for quarks, leptons and weak bosons, since these models may provide us with a better solution of the hierarchy (Fermi-scale) problem /1/. In these models, the weak bosons are spin-one bound states of more fundamental particles (preons) and their mass, i.e. the Fermi scale $G_F^{-1/2}$, is a result of the underlying preon dynamics. However, there is a conceptual difficulty in understanding the observed nature of the weak-intermediate bosons, mainly because of the lack of gauge invariance at the composite level.

By now it is well known /2/ that any non-linear sigma model based on a manifold G/H is equivalent to a theory having a $G \times H_{local}$ symmetry. Here, the H_{local} is a hidden local symmetry, whose corresponding gauge bosons become physical through quantum corrections /3/. With this beautiful basis for composite gauge bosons, Kugo, Uehara and one of the authors (T.Y.) /4/ have recently pointed out that the supersymmetric U(4n+2)/U(4n)xSU(2) model /5/ has a hidden local symmetry which is completely identifiable with the SU(2) gauge group of the standard electroweak theory. The gauge-field nature of composite weak bosons would be immediately understood in these kind of preon theories. However, it seems likely that supersymmetry is broken at scales $\gtrsim 10^2$ GeV /6/ and thus the dynamical basis for these theories might be questionable.

In this short note, we construct a non-supersymmetric preon model whose low-energy dynamics possesses also such an SU(2) hidden local symmetry.

Further, we will discuss some properties of the composite weak bosons related to this hidden local symmetry.

The preon model

Our preon model is basically an extended version of the strongly coupled standard model (Abbott-Farhi model) /7/ with many scalar doublets. Here one has left-handed fermions ψ_L^a ($a = 1 \sim 4n$) which are doublets of the strong $SU(2)_L$ gauge group, where n is the number of quark-lepton families. For each left-handed fermion family, there are three right-handed $SU(2)_L$ singlets u_R, d_R and e_R. For simplicity we will neglect the QCD and electro-magnetic interactions, since these interactions are weak enough at the $SU(2)_L$ confining scale Λ_2 (~ 1 TeV). We introduce, in contrast to the original Abbott-Fahri model, many scalar fields $\phi^{(i)}$ ($i = 1 \sim \ell$) which are $SU(2)_L$ doublets. These scalar fields $\phi^{(i)}$ have Yukawa couplings to the fermions which become the possible sources of masses for the composite quarks and leptons.

Let us discuss the symmetries of this theory, in the limit of all inter-actions besides the strong $SU(2)_L$ confining forces vanishing. In this limit the Lagrangian of our preon theory is simply given by

$$\mathcal{L} = i\bar{\psi}_L^a \not{D} \psi_L^a + i\bar{\psi}_R^a \not{\partial} \psi_R^a + \frac{1}{2}\mathrm{tr}[(D_\mu\Omega)^\dagger(D^\mu\Omega)] - \frac{1}{2}\mathrm{tr}(\Omega^\dagger\Omega\hat{m}^2) - \frac{1}{2g^2}\mathrm{tr}\, G_{\mu\nu}G^{\mu\nu}, \tag{1}$$

where D_μ denotes the usual covariant derivative of the strong $SU(2)_L$ group and $G_{\mu\nu}$ is a strength tensor for the strong gauge field V_μ. (We omit the

$SU(2)_L$ indices, if it does not cause any confusion). Ω is a $2\times2\ell$ matrix which represents the ℓ scalar doublets. To wit:

$$\Omega = \begin{pmatrix} \phi_1^{(1)} & -\phi_2^{(1)*} & \phi_1^{(2)} & \cdots & -\phi_2^{(\ell)*} \\ \phi_2^{(1)} & \phi_1^{(1)*} & \phi_2^{(2)} & \cdots & \phi_1^{(\ell)*} \end{pmatrix}, \tag{2}$$

and $\hat{m}^2$ is $2\ell \times 2\ell$ mass matrix for the scalar bosons;

$$\hat{m}^2 = \begin{pmatrix} m_1 & & & & 0 \\ & m_1 & & & \\ & & \ddots & & \\ & & & m_\ell & \\ 0 & & & & m_\ell \end{pmatrix}^2 . \tag{3}$$

For $m_i = m_j \ll \Lambda_2$, the Lagrangian (1) possesses an exact global symmetry $SU(4n) \times SP(2\ell)$, where the $SU(4n)$ is associated with the global rotation of $4n$ fermions ψ_L^a and $SP(2\ell)$ with the right transformation of scalars, $\Omega \to \Omega h$. We will, hereafter, restrict our discussion to the minimal extension of the Abbott-Farhi model, that is $\ell = 2$.

We shall assume that condensates such as $\langle \phi^{(i)*}\phi^{(i)} \rangle = v_i^2$ form by the strong $SU(2)_L$ forces. Later on we shall try to justify why this happens. If $v_1 \neq v_2$, then the global $SP(4)$ gets broken down to the maximal subgroup $SU(2) \times SU(2)$, producing four massless Nambu-Goldstone (NG) bosons χ^i ($i = 1,2$) (χ^i being complex). This is the most important assumption of our preon dynamics /F.1/.

The low-energy physics of the NG bosons χ^i is well described by a non-linear $SP(4)/SU(2)\times SU(2)$ sigma model, as one can easily demonstrate that this Lagrangian has a hidden local $SU(2)_H$ symmetry /4/. In fact, the effective Lagrangian of the χ's is given, retaining only the lowest derivative terms, by

$$\mathcal{L} = \frac{1}{2} \mathrm{tr} \left[(D_\mu \Sigma)^\dagger (D^\mu \Sigma) \right] , \tag{4}$$

with

$$\Sigma = \begin{pmatrix} v & 0 & \chi_1 & -\chi_2^* \\ 0 & v & \chi_2 & \chi_1^* \end{pmatrix} . \tag{5}$$

The parameter v is a dimension-one constant corresponding to the $\phi^*\phi$ condnesate scale which is assumed to be of order of the confining scale Λ_2. The covariant derivative D_μ is defined as

$$D_\mu = \partial_\mu - i \frac{\tau^i}{2} A^i_\mu \qquad , \text{ with } i = 1\text{-}3, \tag{6}$$

where the A^i_μ are the redundant gauge fields of the hidden local $SU(2)_H$ symmetry. Notice that there is no kinetic term for A^i_μ at the classical level. By integrating over A^i_μ fields in Eq. (4), one easily obtains a non-linear sigma model Lagrangian for the χ's based on the coset space $SP(4)/SU(2)\times SU(2)$.

Because the $SU(2)_L$ is confining, some massless composite fermions are required to satisfy the 't Hooft anomaly matching conditions. We can simply combine the 4n fermions ψ^a_L with some constituent scalar bosons ϕ'_k to make $SU(2)_L$ singlet composite fermions. In the SP(4)-broken phase, one can construct a "constituent" scalar field ϕ'_k, defined as

$$\phi'_k \equiv (\Sigma \Omega^\dagger)_k \quad , \quad k = 1, 2 \quad , \tag{7}$$

which transforms as an ($\underline{2}$, $\underline{2}$) under the local groups $SU(2)_H \times SU(2)_L$ /F.2/.

Notice that any possible dynamical mass for ϕ'_k is SP(4)-invariant, since ϕ'_k is a singlet of SP(4). The left-handed composite fermions $F^a_{Lk} \equiv \psi^a_L \phi'_k$ (or $\psi^a_L \phi'^*_k$) match the SU(4n) anomalies of the fundamental preons ψ^a_L /F.3/. A crucial point here is that the massless composite fermions F^a_{Lk} transform as doublets under the hidden local $SU(2)_H$ symmetries, but as singlets under the SP(4), so that the SP(4) anomaly matching is trivial.

In view of the transformation properties of these massless fermions, the following Lagrangian for F^a_{Lk} is perfectly consistent with all symmetries in the theory:

$$\mathcal{L} = i \bar{F}^a_{Lk} \not{D}_{kl} F^a_{Ll} + O\left(\frac{1}{\Lambda^2} \bar{F} F \bar{F} F\right) . \tag{8}$$

Thus it is sensible to identify F^a_{Lk} with the left-handed quarks and leptons /F.4/.

If the kinetic term of A^i_μ is present, the effective interactions of (8) combined with Eq. (4) have precisely the same form as those in the standard electroweak theory even including the Higgs sector. In fact, one can argue that such a kinetic term may be generated in certain circumstance by quantum effects of the non-linear interactions /3/. If this is indeed the case, we can account for some observed nature of the weak bosons associated with the gauge symmetry; in particular the Weinberg mass formula, $\rho = m_W / m_Z \cos\theta_W = 1$, naturally follows when the electromagnetic interactions are turned on /4/. Moreover, in principle one may be able to calculate the condensate scale $v \sim 250$ GeV (or equivalently m_W) in terms of the scale of the underlying preon dynamics. Therefore, it is very important to attempt to study the dynamics for creating the kinetic term of the redundant gauge fields A_μ^i.

So far we have considered only massless bound states modes. However, it is quite natural to assume that also some massive boundstates are formed, whose masses are of order of the confining scale Λ_2. In calculating the quantum corrections we shall use a cutoff Λ_{cut} to remove all altraviolet divergencies, since the non-linear sigma model is not renormalizable in 4 dimensions /F.5/. Usually, the ultraviolet cutoff Λ_{cut} is taken to be the same as the confining scale Λ_2 and hence the massive modes may not contribute to the quantum corrections. However, in the case of more than four families /F.6/, the running-coupling constant of the strong gauge field V_μ is slowly varying, and one may naturally generate a hierarchic situation /9/ so that really $\Lambda_2 \ll \Lambda_{cut}$. If this is the case, it may be reasonable to take into account the contributions due to massive modes in calculating the quantum corrections for that field. For example, let us consider the $2 \times 4n$ massive Dirac fermion boundstates, $\Psi^a_k = \psi^a_L \phi'_k$ and $\Psi'^a_k = \psi^a_L \phi'^*_k$ /F.7/, with masses of order $M \sim \Lambda_2 \sim 1$ TeV. An integration over these massive fermions Ψ^a_k and Ψ^a_k generates an effective kinetic term for A_μ^i at a low-energy scale which is below the heavy mass M /F8/. Rescaling the gauge field to obtain a standard kinetic energy term identifies /4/ the coupling constant $\bar{\alpha}_2$ of the hidden gauge field A_μ^i. In the leading order in a 1/M expansion /10/ one finds

$$\bar{\alpha}_2 \simeq \frac{3\pi}{8n \log \Lambda_{cut}/M} . \qquad (9)$$

In order to understand the observed value $\bar{\alpha}_2 \sim 1/25$, one must assume Λ_{cut}/M 100 for e.g. five families $n = 5$ /F.9/. It seems rather unlikely to us that such a large hierarchy $\Lambda_{cut} \sim 100$ M really emerges dynamically, from these simple considerations. We need, therefore, other sources for the effective kinetic term which may be, perhaps, related to more involved preon dynamics.

In conclusion, we have shown that the low-energy Lagrangian of the extended Abbott-Farhi model possesses a hidden local SU(2) symmetry which is identifiable with the weak gauge group in the standard model. Because of the presence of this hidden gauge invariance, some important features of weak bosons (e.g. $\rho = m_W / m_Z \cos\theta_W = 1$) are automatically understood, provided that the kinetic term of A_μ^i is indeed produced by the quantum corrections. Although it proves difficult to explain the observed small coupling constant $\bar{\alpha}_2 \sim 1/25$ by the quantum effects considered, we believe that our preon model illustrates nicely how one may solve dynamically one of the basic problems in composite models for weak intermediate bosons - that of the near gauge nature of this particle.

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Footnotes

/F.1/ It seems more likely to have $v_1 = v_2$, where the SP(4) remains unbroken. However, if there is a significant mass difference between the two scalars $\phi^{(i)}$; $m_1 < m_2 \ll \Lambda_2$, the usual simple argument that identifies the vacuum-expectation values of two scalars is no longer applicable. Therefore, one expects that the NG bosons χ's acquire their masses from such an explicit breaking term ($m_1 \neq m_2$). We thank M. Lüscher for raising the dynamical possibility that $v_1 \neq v_2$, necessarily.

/F.2/ A similar situation also happens in QCD. In the broken phase of the chiral $SU(2)_L \times SU(2)_R$ symmetry, the constituent quark q' is defined by $q' \equiv \xi_{L(R)}\, q_{L(R)}$, where $\xi_L^{\dagger} \xi_R = \exp(i\, 2\pi/f_\pi)$ /8/. Here q' transforms as a doublet of the hidden local $SU(2)_H$ but as a singlet of the global $SU(2)_L \times SU(2)_R$, since $\xi_{L(R)}$ transforms as ($\underline{2}$, $\underline{2}$) under the $SU(2)_H \times SU(2)_{L(R)}$. Thus, the mass term, $m\bar{q}'q'$, is clearly invariant under the chiral symmetry.

/F.3/ In the unbroken phase of SP(4) the simplest fermion boundstate $F_{Li}^a \equiv \psi_L^a \Omega_i$ ($i = 1 \sim 4$) does not satisfy the 't Hooft consistency condition, since F_{Li}^a transforms as (4n, 4) under the SU(4n) x SP(4). This may, perhaps, imply that either SU(4n) or SP(4) is likely broken and can be adduced as a justification for the formation of the $\langle \phi^{(i)*} \phi^{(i)} \rangle$ condensates.

/F.4/ The composite quarks and leptons F_{Lk}^a will get masses through Yukawa couplings $\psi_L^a\, \bar{q}_R (\bar{l}_R)\, \phi^{(i)}$. These interactions also generate the masses of NG boson χ^i, since they violate not only the SU(4n), but also SP(4) invariance. However, these interactions never break the hidden local symmetry, since all preons ψ_L^a and $\phi^{(i)}$ are singlets of the hidden group.

/F.5/ It is natural to have the ultraviolet cutoff at the momentum $\sim \Lambda_{cut}$, because our effective theory is valid only at distances larger than $1/\Lambda_{cut}$.

/F.6/ For the five families n = 5, the running coupling constant of $SU(2)_L$ strong interactions is slowly varying, since $\beta(g)$ is very small, ($b = 22 - 4n - 1 = 1$). In this case one may have a dynamical situation in which the scale Λ_{cut} is much greater than the confining scale Λ_2 /9/. Thus it might not be unreasonable to think of having an effective field theory in the intermediate region between Λ_2 and Λ_{cut}, where the "constituent" scalar preons ϕ_k' also interact with the NG bosons. The preon-loop diagrams also contribute to the generation of the kinetic term. However, the calculation of such effects is far beyond the scope of this short note, since we cannot neglect the non-perturbative effects of the strong $SU(2)_L$ interactions.

/F.7/ Ψ_k^a or $\Psi_k'^a$ may be thought as excited-states of the 't Hooft massless composite fermions F_{Lk}^a.

/F.8/ The Lagrangian of the massive fermions is given by

$$\mathcal{L} = \bar{\Psi}^a \not{D} \Psi^a + \bar{\Psi}'^a \not{D} \Psi'^a .$$

The loop diagrams of Ψ^a_k and Ψ'^a_k will give rise to a kinetic term for A^i_μ of the form: $\mathcal{L}_{eff} = -\frac{1}{4} \kappa^{-1} F^i_{\mu\nu} F^{i\mu\nu}$. In the leading order in an 1/M expansion one finds $\kappa = 12\pi^2/8n \log \Lambda_{cut}/M$ which leads to Eq. (9).

/F.9/ In this case, a careful analysis is necessary, since the QCD coupling constant increases in the region between the heavy mass M and Λ_{cut}.

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