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# ON COMPOSITENESS IN THE WEAK BOSON SYSTEM *)

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**Abstract:**

We discuss some aspects of compositenesa related to the hypothesis that the weak boson system represents a mass protected bound state. We present a framework which possesses a non-perturbative lattice regularization and where no No-go theorems valid for QCD like theories deny this possibility a priori. Above the weak bosons one more "light" isotriplet vector boson with some peculiar properties, presumably easy to distinguish experimentally, is predicted. Combination with a model of quark-lepton compositeness allows for heuristic arguments suggesting that three generations might be prefered over other numbers of replica. The arguments can in principle be sharpened in a non-perturbative formulation of the model.

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The experimental absence of nearby resonances and scalar bound states in the weak boson system leads us to the tentative concept of a mass protected bound state, that is if $m_w$ is the weak boson mass and $\Lambda_w$ the scale of compositeness one should have

$$m_w << \Lambda_w \quad ; \tag{1}$$

in the isotriplet spin one channel.- The breaking of the, now global, $SU_2^{weak}$ is accomplished by the Hung-Sakurai mixing mechanism [(1)]. The Higgs [(2)] is absent from the theory, $m_w$ is due to the bound state dynamics.
To maintain (1) against radiative corrections the low energy ($E << \Lambda_w$) theory of the weak boson must be renormalizable. Since the weak boson is a massive and (here globally) nonabelian particle this implies the vanishing of all dimension 4 self-couplings[(3)]. Among other things this leads to a definite relation between the tree level vertices of $Z_o \longrightarrow W^+W^-$ in the present model and the Standard Model[(2)] [S.M.]. (Roughly a suppression by a factor $-tg^2\,\theta_w$ occurs relative to the S.M.[(4)] ($\theta_w$ is the Weinberg angle[(2)]).

Fig. 1 gives an example of this relation and shows that at LEP 200 one can unambiguously decide whether the weak boson is a gauge particle[(2)] or a "light" bound state obeying (1).

To fix a theoretical framework we first constrain the allowed preon fields by the requirements:

(1) The bound state theory must not have Goldstone bosons triggered by confinement.

(2) The theory must not have scalar bound states lighter than the weak bosons.

This excludes Dirac fermions and scalars as preons. We therefore try Weyl fermions in $SU_2$ representations. More precisely we consider a field $F^A_{\alpha,i}(x)$ where $\underline{\alpha}$ refers to the (undotted) spin index, $\underline{i}$ to the fundamental representation of an $SU_2$ hypercolor gauge group and $\underline{A}$ denotes the fundamental representation of the global $SU_2^{weak}$ .

Now the only Lorentz and gauge invariant condensate that can form, $\langle 0|F^T(x)\ i\sigma_2\ i\tau_2\ iT_2\ \ F(x)|0\rangle$ (where $i\sigma_2$, $i\tau_2$, $i\cdot T_2$ are the matrix $\begin{pmatrix} 0 & 1 \\ -1 & 0 \end{pmatrix}$ in spin, hypercolor and $SU_2^{weak}$ space respectively), is also $SU_2^{weak}$ invariant. Hence no analogue of the pion (scalar isotriplet) exists. Isosinglet states become heavy[5]. Indeed one can argue that the spin one isotriplet state may well be the lightest one contrary to QCD[6].

We note that the strongly coupled Weyl fermion theory has a gauge invariant lattice regularization quite similar to the Wilson lattice[7] except that the species doublers are suppressed by Majorana masses instead of Dirac masses. Because of the required brevity of this report we do not give the details here[4].

To get a feeling for the dynamics let us compare with QCD and ask: can QCD exist in a phase with $m_\rho = 0$ and $m_\pi \neq 0$ ? We expect the answer to be negative[6] but it is interesting to see how it happens in the lattice strong coupling limit. If we adjust in the effective action $m_\rho = 0$ we end up with $m_\pi^2 < 0$ . A pertinent shift, of the pion field $\vec{\pi} \longrightarrow \langle\vec{\pi}\rangle + \delta\,\vec{\pi}$ ($\vec{\pi} = \bar{\Phi}\gamma_5\vec{\tau}\,\Phi$) would introduce a spontaneous breaking of isovector rotations, which is forbidden in QCD[8].

In the Weyl fermion theory however there is no scalar particle whose mass is protected by the chiral structure of the vacuum so as to produce an inconsistency if one adjusts for $m_W = 0$ . However then a spontaneous breaking of Lorentz invariance in the isotriplet channel is required (at least in the continuum limit; this satisfies the theorem of ref. (9), as well as an "inverse Goldstone theorem"[4]):

$$\langle 0|F^\dagger(x)\ \sigma^U\ \ T^A\ F(x)\ |0\rangle \quad \neq 0\ ; \tag{2}$$

Note that there is no analogue of the theorem of ref. (8) that would forbid (2).- One can give a heuristic argument[4] to the extent that a small variation of the effective action constructed from the lattice (namely of the Wilson type adjustable mass parameter[4],[7] $M \longrightarrow M + \delta M$, $\delta M \ll \Lambda_W$) can lead to a phase with

$m_\omega \neq 0$ and $\langle 0 | F^\dagger(x)\, \sigma^\mu\, T^A\, F(x) | 0 \rangle = 0$ . If this transition is smooth one expects roughly

$$m_\omega \sim \delta M << \Lambda_\omega ; \qquad (3)$$

The nearby 'Goldstone' phase protects $m_\omega$ still in the Lorentz invariant phase.

The lattice formulation allows for a nonperturbative check of this proposal which is also of interest per se, independent of the applied notion of compositeness.

If the theory behaves according to this proposal then before varying $M$ away from the point where $m_\omega = 0$ the two condensates $\langle 0 | F^\dagger(x)\, \sigma^\mu\, T^A\, F(x) | 0 \rangle$ and $\langle 0 | F^T(x)\, i\sigma_2\, i\tau_2\, i \cdot T_2\, F(x) | 0 \rangle$ can coexist (without introduction of a tachyon, as we have seen). In this case also the condensate

$$\langle 0 | B^A_{\alpha\beta}(x) | 0 \rangle = \langle 0 | F^T(x)\, \sigma^T_\alpha\, \sigma^*_\beta T^A F(x) | 0 \rangle \neq 0 ; \qquad (4)$$

will occur ( $\sigma^0 = \begin{pmatrix} 1 & 0 \\ 0 & 1 \end{pmatrix}$ , $T$ means transposed). The same arguments used with (2) also apply to (4) and we conclude that there is another "light" particle created by the antisymmetric and self-dual field $B^A_{\alpha\beta}(x)$ after $M \longrightarrow M + \delta M$ was performed as described above (since it is "baryon like" we expect $m_\omega < m_B << \Lambda_\omega$ ). However the selfdual field $B^A_{\alpha\beta}(x)$ cannot be quantized in Minkowski space, instead it can be shown that the field $D^A_\beta(x) = \partial^\alpha B^A_{\alpha\beta}(x)$ propagates the spin one isotriplet modes with mass $m_B$ in Minkowski space (see ref. (4) for a discussion of this point). Since $\partial^\alpha D^A_\alpha(x) = 0$ any (chiral) current coupling to $D^A_\beta(x)$ must obey $\partial^\alpha J^A_\alpha(x) = 0$ . No such dim = 3 current exists. Hence $J^A_\alpha(x)$ must be a dim > 3 current which is made conserved by construction (and does not violate electric charge conservation when coupled to $D^A_\alpha(x)$; we choose $Q_{up} = \frac{1}{2}$; $Q_{down} = -\frac{1}{2}$ for the $F(x)$ field)

$$J^A_\beta(x) \sim \frac{1}{\Lambda_Q^2}\, \partial^\alpha\, [\bar{Q}_L(x)(\,\partial_\alpha\sigma_\beta - \sigma_\alpha\partial_\beta) T^A Q_L(x)] ; \qquad (5)$$

( $Q_L(x)$ denotes a lefthanded quark or lepton field and $\Lambda_Q$ the compositeness scale for quarks and leptons).

Further since chirality is not a good quantum number $D^A_\beta(x)$ embodies two independent fields (its Re- and Jm-part). Therefore if produced in $e^+e^-$ collisions via $\gamma^*$ the cross section will look as for a double charged particle, although the ionization per track corresponds to unit charge. Because of (5) coupling to quarks and leptons is highly suppressed ($\Lambda_Q \gtrsim$ 100 TeV presumably). The same holds for coupling to weak bosons because the particle related to $D^A_\beta(x)$ is again mass protected[4]. Therefore it will be unusually long lived. Its discovery would be as important for the present theory as the Higgs is for the S.M..

Finally we wish to relate the model to one with quark and lepton compositeness. Introducing an $SU_2(\Lambda_Q)$ hypercolor different from the earlier $SU_2(\Lambda_w)$ where in brackets we have denoted the scale where the gauge force becomes strong ($\Lambda_Q >> \Lambda_w$ allows to understand why weak interactions are weak[4], [10]) we write the composite lefthanded up-down quarks and leptons collectively as

$$U_L^{c,N_F} = L_a^{c,N_F} \cdot \emptyset_a^* \; ; \quad D_L^{c,N_F} = L_a^{c,N_F} \varepsilon_{ab} \emptyset_b \; ; \qquad (6)$$

where $c = 1,2,3$ is for SU(3) QCD color and $c = 4$ for leptons.

$\emptyset_a$ and $L_a^{c,N_F}$ are scalar and lefthanded fermions respectively in the fundamental representation of $SU_2(\Lambda_Q)$ and $N_F$ denotes family replication. The $F^A_{\alpha,i}(x)$ fields are now composite at scale $\Lambda_Q$

$$F_i^{up} = X_{ia} \emptyset_a^* \; ; \qquad F_i^{down} = X_{ia} \varepsilon_{ab} \emptyset_b \; ; \qquad (7)$$

up-down play the role of the earlier index $\underline{A}$ and $X_{ia}(x)$ is a Weyl fermion that transforms both in the fundamental representation of $SU_2(\Lambda_w)$ and $SU_2(\Lambda_Q)$ .

We sketch some heuristic arguments indicating that $N_F$ could be fixed by requiring that the $SU_2(\Lambda_Q)$ dynamics be confining (in the sense of absence of hypercolor asymptotic states) and chirality preserving, that is $\langle 0 | L^{TM}(x)\, i\sigma_2\, i\tau_2\, L^{M'} | 0 \rangle = 0$ with $M, M' = 1,\ldots,4\,N_F$ .

Asymptotic freedom requires $N_F \leq 4$ (this bound occurs rather narrow, $4\,N_F < 19{,}5$ , and thus depends on the details of the

model's field content). Further for $N_F = 4$ the 2-loop $\beta$-function gives an infrared stable fixed point (IRSFP) at $x = g^2/(4\pi)^2 = 0,1$ and for $N_F = 3$ at $x = \frac{10}{11}$ (it was assumed that only the $L(x)$ - preons remain massless in the infrared). At an IRSFP no dimensional transmutation can occur and the gauge forces do not generate a confinement scale. Clearly the $N_F = 3$ IRSFP we cannot trust and for the rest of this note we assume that it is not there. The $N_F = 4$ IRSFP is deeper in the perturbative region and we take it as an indication that with a better knowledge of $\beta(g^2)$ (no further zeros ?) we could eventually rule out a 4th family. We are then left with $N_F = 1,2$ or 3. Now the formation of chirality breaking condensates can stop if the fermionic screening is strong enough (which is particularily severe for massless particles(12); a related effect (11)
is observed in QCD(13)). A rough estimate leads to the condition $\exp\left(\frac{1}{4N_F}\right) \sim 1$ for fastest screening. So the largest allowed $N_F$, namely $N_F = 3$, has the best chance to keep chirality and to screen asymptotic hypercolor charges. Since a nonperturbative formulation can be set up(11) these heuristic considerations can be sharpened by looking, for example, for a chiral phase transition at $N_F = 3$. We also note that in this model $\theta^{QCD}$ can be rotated away(11). Further the absence of W - families is related to the reqiured absence of Goldstone bosons. This is further discussed in ref.(11).

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Figure Caption:

Figure 1 shows the angular distribution for $e^+e^- \longrightarrow W^+W^-$ at threshold energy for the S.M. and for the case of a 'light' bound state weak boson. In the latter case the effect of a variation of the anomalous magnetic moment <u>k</u> is also plotted.

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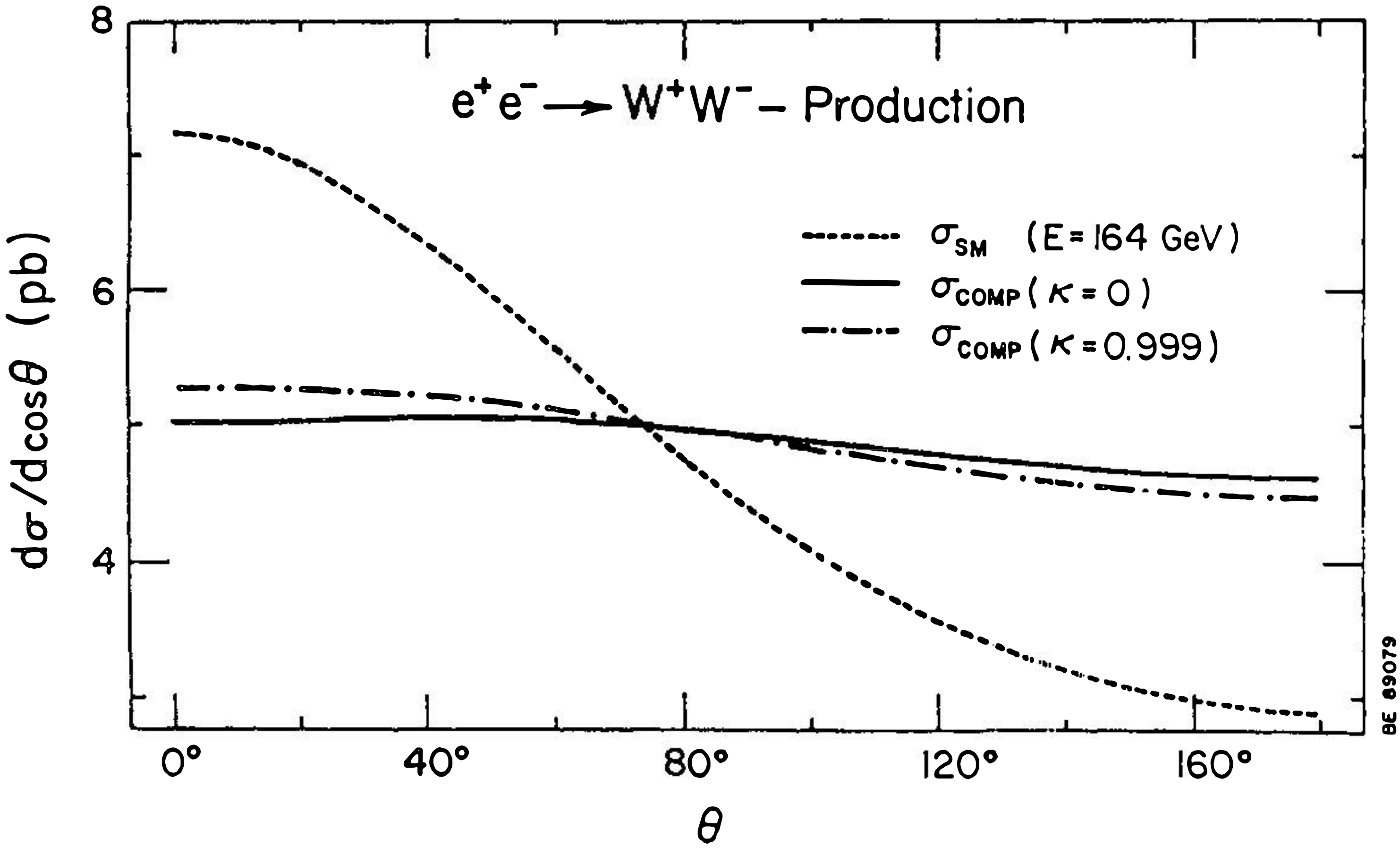
e+e- → W+W- – Production
σSM (E=164 GeV)
σCOMP (κ=0)
σCOMP (κ=0.999)
dσ/dcosθ (pb)
8
6
4
0°
40°
80°
120°
160°
θ
BE 89079