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New Phenomena Below the Mega TeV Scale*

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Abstract

After some musings about a Mega TeV scale, I contrast the difference expected between theories where the $SU(2)\times U(1)$ group is broken dynamically and theories where this breakdown is induced by the vacuum expectation value of an elementary Higgs field. Although high precision electroweak data, whose impact I review, do not favor the dynamical option they also do not rule it out, making it imperative to explore the Higgs-dynamical symmetry breaking dichotomy at higher energies. Real differences between these options emerge in W_LW_L scattering, with this process showing both resonance formation and multiparticle production if $SU(2)\times(1)$ is broken dynamically. Expectations for multi W_L production in the Eloisatron energy range are discussed, as well as some of the difficulties and challenges one faces in trying to generate fermion masses dynamically.

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1. Musings on a Mega TeV scale

Thinking about physics below a Mega TeV (10^9 GeV) is a nice, liberating experience. For one thing, viewed from this scale, electroweak symmetry breaking is truly a "low energy" phenomenon, whose manifestations should by then all be apparent. Of course, it is possible that a scale around a Mega TeV has physical significance per se. For instance, $M = 10^9$ GeV is numerically the geometric average of the Planck Scale and the dynamical scale of QCD,

$$M = \sqrt{M_P \Lambda_{QCD}} \tag{1}$$

As such it can serve as an ideal intermediate scale.[1] It is also a "see saw" scale for neutrino masses, since it gives a mass for ν_τ which renders this particle a plausible dark matter candidate[2]:

$$m_{\nu_\tau} \simeq \frac{m_\tau^2}{M} \sim 0(eV) \tag{2}$$

For the purposes of this lecture, however, the principal virtue of M is really that it is a scale which is sufficiently far above the electroweak symmetry breaking scale $v \simeq 250$ GeV, so that the statement $M >> v$ truly applies. As we shall see, being able to consider electroweak symmetry breaking effects far above threshold has its rewards. Furthermore, M is sufficiently large that it is possible (likely ?) that, if the electroweak symmetry breaking is dynamical, the scale Λ_m associated with the dynamics of fermion mass generation is also below M. Unfortunately, the fermion mass issue in models of dynamical symmetry breaking is still poorly understood,[3] so I will concentrate for the most part in this lecture on effects connected purely with the weak boson sector. Nevertheless, even with this restriction, we will encounter quite a rich and exciting physics.

2. What do we know: electroweak symmetry breaking and the equivalence theorem

The very successful comparison of data from LEP with the standard model predictions,[4] provides strong evidence that electroweak interactions are described by a spontaneously broken $SU(2) \times U(1)$ gauge theory.[5] In addition, the fact that the measured ρ parameter[4]

$$\rho = 0.9968 \pm 0.0048 \tag{3}$$

is so close to unity argues that the theory possesses some (approximate) custodial global symmetry to guarantee this property.[6] These results notwithstanding, however, one still has no knowledge of what the detailed nature of the spontaneous symmetry breaking is. Thus one can still contemplate the alternatives that the order parameter for the $SU(2) \times U(1) \rightarrow U(1)_{em}$ breakdown is either:

i) the vacuum expectation value of an elementary scalar field: $< \Phi > \neq 0$

or

ii) the condensate of fermion-antifermion pairs of some underlying theory: $< \bar{T}_L T_R > \neq 0$

The physics associated with these two options is quite distinct. The first option, the elementary Higgs option,[7] is perfectly compatible with weak coupling. That is, one can imagine the formation of a vacuum expectation value for the Higgs field Φ, even though the interaction strength λ in the Higgs potential

$$V = \lambda(\Phi^\dagger \Phi - \frac{v^2}{2})^2 \tag{4}$$

is small, $\lambda << 1$. In the second option, called here Technicolor[8] for short, however, one cannot avoid strong coupling in the underlying theory, since only in this regime will the $< \bar{T}_L T_R >$ condensate form. The strength of this condensate is measured by the dynamical scale Λ_T at which the coupling constant $\alpha_T(q^2)$ becomes of order unity:

$$< \bar{T}_L T_R > \simeq \Lambda_T^3 \tag{5}$$

If this condensate is to be associated with electroweak symmetry breaking, then Λ_T must be of order of the electroweak symmetry breaking parameter v:

$$\Lambda_T \sim v = (\sqrt{2} G_F)^{-1/2} \simeq 250 \ GeV \tag{6}$$

So, in this case, we expect that effects associated with the strong coupling underlying theory be visible in this energy range.

In both the Higgs case and in Technicolor, the physics of the spontaneous breakdown will be reflected in the amplitudes for longitudinal gauge boson scattering - $W_L W_L$ scattering. The amplitudes for processes involving longitudinal gauge bosons, including $W_L W_L$ scattering itself, are governed by an important "equivalence theorem", proven long ago by Cornwall, Levin and Tiktopoulos.[9] This theorem states that, in the limit in which the energies of all longitudinal gauge bosons are much bigger than their masses ($E_i >> M_W$), one can obtain the longitudinal gauge boson amplitude directly from an amplitude in which each W_L is replaced by its corresponding Goldstone boson, ensuing from the spontaneous breakdown. That is, schematically, if one denotes the Goldstone boson corresponding to W_L by w one has, in the high energy limit,

$$A(W_{L_1}, W_{L_2}, \cdots, W_{L_n}) \ = \ A(w_1, w_2, \cdots, w_n) \tag{7}$$

For the elementary Higgs case, one can write the Higgs field Φ as

$$\Phi = \frac{1}{\sqrt{2}} \begin{pmatrix} v + H + i w_3 \\ w_1 + i w_2 \end{pmatrix} \tag{8}$$

where H is the physical Higgs field. Then the potential (4) becomes simply

$$V = \frac{\lambda}{4}(\vec{w}^2 + H^2)^2 + \lambda v H(\vec{w}^2 + H^2) + \lambda v^2 H^2 \tag{9}$$

One sees that, if λ is small, one will have weak interactions among the Goldstone fields and with the Higgs field H. Hence, from the equivalence theorem one expects also weak interactions among the longitudinal gauge bosons W_L. If, on the other hand, the spontaneous breakdown of the electroweak theory occurs dynamically, through the formation of condensates in some strongly interacting underlying theory, then one would expect that the interactions among the Goldstone bosons in the underlying theory will also be strong. In turn this will lead to strong interactions among the W_L. It is this phenomena which we want to explore in some detail here.

Before doing so, however, it is important to note a feature of the Higgs potential which needs to be incorporated in all underlying dynamically symmetry breaking theories. The potential (4), with Φ a complex $SU(2)$ doublet, has a larger symmetry that $SU(2) \times U(1)$, namely 0(4). When Φ acquires a vacuum expectation value $< \Phi >$, $0(4) \rightarrow 0(3) \sim SU(2)$. It is this custodial $SU(2)$ symmetry[6] which, neglecting fermion mass effects, guarantees that $\rho = 1$. For the dynamical symmetry breaking theory one wants to retain this desirable feature. This is most easily done by requiring that the Technicolor theory possess a global $SU(2)_L \times SU(2)_R$ symmetry, under which T_L and T_R transform as $(2,1)$ and $(1,2)$, respectively. The condensate $< \bar{T}_L T_R > \neq 0$ then breaks this symmetry to the diagonal subgroup, $SU(2)_L \times SU(2)_R \rightarrow SU(2)_{L+R}$, and one has again a custodial $SU(2)$ symmetry to guarantee that $\rho = 1$.

3. Signals of dynamical symmetry breaking near threshold

As I mentioned above, the tests of the electroweak theory at $\sqrt{s} \simeq 100\ GeV$ are not terrifically sensitive to the symmetry breaking sector. One hopes to learn indirectly about this sector through its effects on radiative corrections. However, for the Higgs case one knows that these effects are small as a result of the, so called, Veltman's screening theorem.[10] This theorem basically states that, at $0(\alpha)$, radiative corrections depend on the Higgs mass only logarithmically. Although there are also quadratic dependences on the Higgs mass, these only enter at $0(\alpha^2)$, so that

$$\text{Rad. Corr} \sim \frac{\alpha}{\pi}\, \ell n\ \mathrm{M_H/M_Z} + (\frac{\alpha}{\pi})^2 \frac{\mathrm{M_H^2}}{\mathrm{M_Z^2}} \tag{10}$$

Because of Eq. (10), one does not expect significant problems with the usual perturbative treatment of the Higgs sector until the Higgs mass is sufficiently big, so that the first term above is of order of the second:

$$M_H^2 \sim \frac{\pi}{\alpha} M_Z^2 \tag{11}$$

This corresponds to the regime where the strength λ of the Higgs potential is much greater than unity ($\lambda >> 1$), corresponding to strong coupling for the elementary Higgs theory. If λ is small, however, the radiative corrections are largely insensitive to the value of the Higgs mass.

There is a roughly analogous situation in the case of dynamical symmetry breaking. It is rather easy to show that, at low energies, there is weak scattering among the Goldstone bosons, arising from the fact that because they are Goldstone bosons they couple to themselves and to other fields derivatively. So, for instance, the amplitude for $w^+w^- \to zz$ scattering is given by

$$\mathcal{A}(w^+w^- \to zz) \;=\; \frac{s}{v^2} \tag{12}$$

where $\sqrt{s}$ is the w^+w^- CM energy. This amplitude is weak as long as $\sqrt{s} < v$. Furthermore, even if it grows with energy, in fact it does not hit the elastic unitarity limit until $\sqrt{s} = 4\sqrt{\pi}\, v \;\simeq\; 1.7\; TeV$.[11] Nevertheless, in general there are differences in the contributions to the radiative corrections to the electroweak theory depending on whether the symmetry breakdown is caused by an elementary Higgs field acquiring a vacuum expectation value or whether this breakdown is more dynamical.

To appreciate these differences, it is useful to focus on the gauge propagators of the electroweak theory. Let us denote the self energy functions for these propagators as $\Pi_{AB}(q^2)$, where the indices AB span the 4 possibilities of the $SU(2) \;\times\; U(1)$ theory: ($AB = WW; ZZ; \gamma\gamma; \gamma Z$). Since one is interested in the effects of the physics of the symmetry breaking sector for the relatively low q^2 values represented by the W and Z masses, it suffices to retain just the first two terms in a q^2 expansion of $\Pi_{AB}(q^2)$. Thus, let us approximate[12]

$$\Pi_{AB}(q^2) \;\simeq\; \Pi_{AB}(0) + q^2\Pi'_{AB}(0) \;\; . \tag{13}$$

Because of photonic gauge invariance, one can check that $\Pi_{\gamma\gamma}(0) = \Pi_{\gamma Z}(0) = 0$, so Eq. (13) contain 6 independent parameters. Three of these parameters can be expressed in terms of the 3 independent parameters of the $SU(2)\times U(1)$ theory, the coupling constants g_2 and g_1, and the Fermi scale v. In fact, it is more convenient instead of using the set $\{g_1, g_2, v\}$ to replace these parameters by three equivalent physical constants which are well measured experimentally: α, M_Z and G_F.[13] The remaining 3 constants in Eq. (13) have been chosen variously in the literature, and I find it useful here to follow the conventions of Peskin and Takeuchi[14] who have denoted these additional parameters simply as S, T and U.

The parameter S, which is defined as[15]

$$S = 16\pi\{\Pi'_{33}(0) - \Pi'_{3\gamma}(0)\} \tag{14}$$

is particularly interesting because it is rather small for a large range of Higgs masses, but is quite sensitive to possible chirality breaking contributions in a

theory where the symmetry breaking is done dynamically. It is useful to pick a reference Higgs mass (and a reference top mass, although that will not be relevant for our discussion here) and try to predict where S will be relative to this standard. For the case of pure Higgs breaking, of course, this is trivial. Keeping leading contributions only, one has

$$\Delta S \simeq \frac{1}{12\pi} \ell n \; M_H^2/(M_H^2)_{\rm ref} \tag{15}$$

Picking $(M_H)_{\rm ref} = 100\ GeV$, one sees that $\Delta S \leq 0.1$ even for $M_H = 1000\ GeV$.

For any dynamical symmetry breaking model one can, in principle, try to compute the expected S value relative to the $(M_H)_{\rm ref}$ standard. This is not straightforward because it requires making various assumptions on the underlying strong dynamics. Nevertheless, Peskin and Takeuchi[14] made some progress in this direction for a class of theories in which a vectorial $SU(2)$ is conserved by the underlying theory. In view of the dynamical need for such an $SU(2)$ to obtain $\rho = 1$, these theories are certainly interesting candidates. In these theories only the vector-vector (VV) and axial-axial (AA) components enter in the vacuum polarization Π_{33}, while Π_{3_γ} only contains a VV piece. Thus S can be rewritten entirely in terms of the vacuum polarization functions which enter in the current correlation functions for the $SU(2)_V$ and $SU(2)_A$ currents in the underlying theory. A simple calculation then gives

$$S = -4\pi\{\Pi'_{VV}(0) - \Pi'_{AA}(0)\} \tag{16}$$

This nice formula for S allows one to attempt estimating this quantity in theories of dynamical symmetry breakdown by "scaling up" what happens in QCD. Two flavor QCD with $m_u = m_d = 0$ has a global $SU(2)_V \times SU(2)_A$ symmetry which is broken spontaneously to $SU(2)_V$ by the presence of condensates $< \bar{u}u > = < \bar{d}d > \neq 0$. In QCD, the vacuum polarization difference entering in Eq. (16) can be rewritten in terms of an integral difference between the vector and axial vector spectral functions of the theory

$$\Pi'(0)_{VV} - \Pi'_{AA}(0) = -\frac{1}{4\pi}\int_0^\infty \frac{ds}{s}[v(s) - a(s)] \tag{17}$$

This integral can be estimated by making use of the two Weinberg spectral function sum rules[16] obeyed by $v(s)$ and $a(s)$

$$\int_0^\infty ds[v(s) - a(s)] = 4\pi f_\pi^2 \tag{18a}$$

$$\int_0^\infty s ds[v(s) - a(s)] = 0 \quad , \tag{18b}$$

where f_π is the pion decay constant, along with the assumption that $v(s)$ and $a(s)$ are saturated by the ρ and A_1 poles. Saturating Eqs. (18) with the ρ and A_1, poles, respectively,

$$v(s) = g_\rho^2\delta(s - m_\rho^2) \;\; ; \;\; a(s) = g_{A_1}^2\delta(s - m_{A_1}^2) \tag{19}$$

allows one to determine the strengths g_ρ^2 and $g_{A_1}^2$ in terms of m_ρ^2, $m_{A_1}^2$ and f_π^2, and gives for the integral that enters in Eq. (17) the formula

$$\int_0^\infty \frac{ds}{s}[v(s) - a(s)] = \frac{4\pi f_\pi^2}{m_\rho^2}[1 + \frac{m_\rho^2}{m_{A_1}^2}] \tag{20}$$

The original technicolor theory of Susskind and Weinberg[8] was precisely just a "scaled" up version of 2 flavor QCD where f_π is replaced by v, the Fermi scale, and the Goldstone pions are replaced by technipions which, effectively, get absorbed to become the longitudinal components of the weak gauged bosons. Many other dynamical symmetry breaking theories have the same structure, except that perhaps the technicolor group is $SU(N)_{TC}$ not simply $SU(3)$ and there may be a variety of species of technicolor doublets, not just simply T_u and T_d to mimic the u and d quarks of QCD. For all of these theories, one should be able to estimates S by just "scaling up" the formula (20) with some simple modifications. Large N arguments suggest that[17]

$$\frac{f_\pi^2}{m_\rho^2} \to \frac{3}{N_T}\frac{v^2}{m_{\rho_T}^2} \quad ; \quad \frac{m_\rho^2}{m_{A_1}^2} \to \frac{m_{\rho_T}^2}{m_{A_T}^2} \quad , \tag{21}$$

with m_{ρ_T} and m_{A_T} the analogons of m_ρ and m_{A_1}. Furthermore, if there are n technicolor doublets, each of these doublets should give an equivalent contribution. Hence, one arrives at the simple formula

$$S \simeq \frac{N_T}{3}\, n \int_0^\infty \frac{ds}{s}[v(s) - a(s)]\Big|_{QCD} \simeq 0.25[\frac{N_T}{3}n] \tag{22}$$

The above rough estimate of S agrees qualitatively with the more careful evaluation of the QCD integral of Eq. (17) done by Peskin and Takeuchi.[14] It is also in general accord with the work of Cahn and Suzuki,[18] although these authors argue that the errors in evaluating the QCD integral could actually be consistent with a much smaller value for S than that given in Eq. (22). At any rate, two things regarding S are apparent from this exercise of "scaling up" QCD:

i) $S > 0$

ii) In technicolor models containing a full complement of techniquarks, to account for both color and lepton number, $n = 4$ and then $S \sim 0(1)$.

It is possible to extract values for S (and T, U) allowed by experiment by comparing observable parameters, like the mass of the W boson or the leptonic width of the Z boson, to the theoretical predictions of this generalized $SU(2) \times U(1)$ analysis. The value of any experimental observable $0_{\rm exp}$ can be written as

$$0_{\rm exp} = 0_{SM}(M_H^{\rm ref} \ , m_t^{\rm ref}) + \alpha_0 S + \beta_0\, T + \gamma_0 U \ , \tag{23}$$

where $0_{SM}(M_H^{\rm ref}\ ,\ m_t^{\rm ref})$ is the standard model prediction, assuming a single Higgs doublet does the breaking, with some reference choice for the values of the Higgs and top masses. The coefficients α_0, β_0 and γ_0 depend on the observable in question, but are fixed once $M_H^{\rm ref}$ and $m_t^{\rm ref}$ are chosen. Table 1, taken from Peskin and Takeuchi,[14] gives the dependence of various observable on S, T and U for the choice of $M_H^{\rm ref} = 1\ TeV$ and $m_t^{\rm ref} = 150\ GeV$.

Table 1: Dependence of Observables on S, T and U

$$
\begin{aligned}
M_W/M_Z &= 0.8787 - 3.15\times 10^{-3}S + 4.86\times 10^{-3}T + 3.70\times 10^{-3}U \\
\Gamma_Z(GeV) &= 2.484 - 9.58\times 10^{-3}S + 2.615\times 10^{-2}T \\
\Gamma_\ell(MeV) &= 83.5 - 1.91\times 10^{-1}S + 7.83\times 10^{-1}T \\
\Gamma_u(MeV) &= 296.2 - 1.92S + 3.67T \\
\Gamma_d(MeV) &= 382.3 - 1.72S + 4.20T \\
\Gamma_b(MeV) &= 377.9 - 1.72S + 4.20T \\
\Gamma_{\rm had}(GeV) &= 1.7348 - 9.00\times 10^{-3}S + 1.993\times 10^{-2}T \\
R_Z = \frac{\Gamma_{\rm had}}{\Gamma_\ell} &= 20.78 - 5.99\times 10^{-2}S + 4.24\times 10^{-2}T \\
\sin^2\theta_w^{\rm eff}(M_Z) &= 0.2337 + 3.59\times 10^{-3}S - 2.54\times 10^{-3}T \\
A_{LR} &= 0.1297 - 2.82\times 10^{-2}S + 2.00\times 10^{-2}T \\
A_{FB}^b &= 0.0848 - 1.97\times 10^{-2}S + 1.40\times 10^{-2}T \\
A_{FB}^\ell &= 0.0126 - 6.72\times 10^{-3} + 4.76\times 10^{-3}T \\
g_L^2 &= 0.3001 - 2.67\times 10^{-3}S + 6.53\times 10^{-3}T \\
g_R^2 &= 0.0302 + 9.17\times 10^{-4}S - 1.94\times 10^{-4}T \\
R_\nu &= 0.3126 - 2.32\times 10^{-2}S + 6.46\times 10^{-2}T \\
R_{\bar\nu} &= 0.3824 - 2.77\times 10^{-3}S + 6.03\times 10^{-3}T \\
Q_W(Cs) &= -73.31 - 0.79S - 0.01T
\end{aligned}
$$

In Fig. 1, taken again from Peskin and Takeuchi,[14] I show the allowed range (both at 68% C.L. and 90 % C.L.) for S and T from a combined analysis of all extant data. This analysis clearly favors **negative** S values, with $S \simeq -0(1)$. Similar results to these have been obtained also by Altarelli, Barbieri and Jadach[19] and by Langacker and Luo.[20] Also shown in Fig. 1 are the expectations for S and T in the standard model, as one lets M_H and m_t vary from their reference values of 1000 GeV and 150 GeV, respectively. Note that T is very sensitive to changes in m_t, but quite insensitive to M_H changes. S is largely insensitive to changes in both m_t and M_H. Note also that the expected values for S for a full generation ($n = 4$) of techniquarks of an $SU(4)_{TC}$ theory is about 3σ away from the best fit.

Although the above result is disappointing if one believes in dynamical symmetry breaking, it is not totally disastrous. The result for S, one is comparing

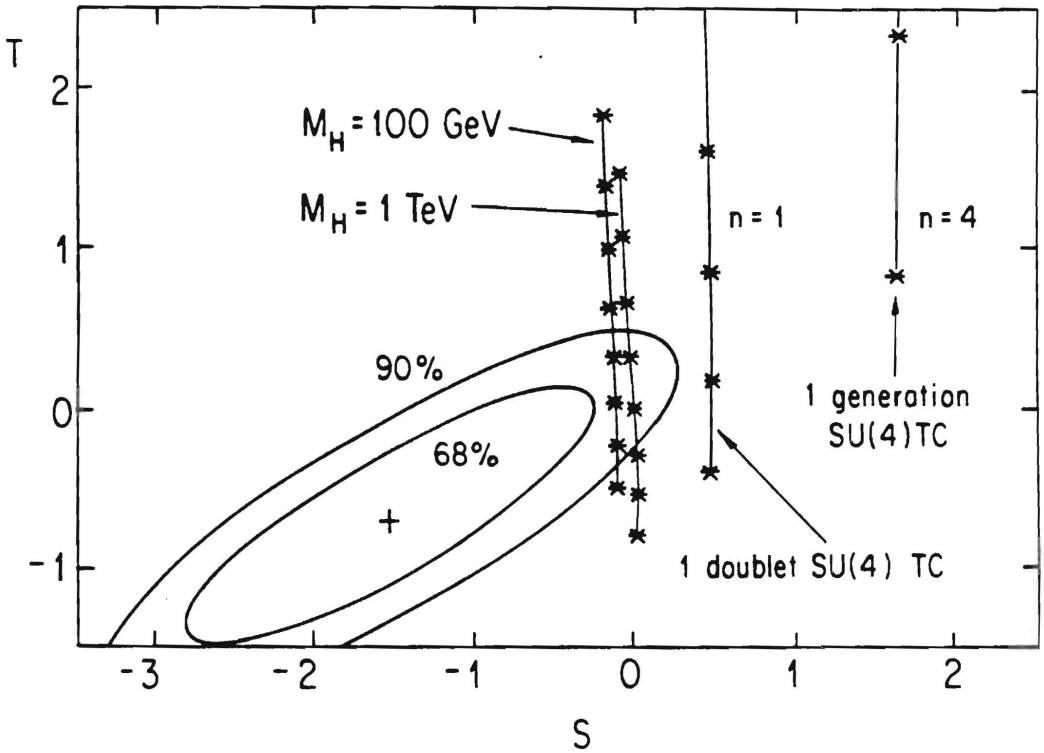

Figure 1: 68% C.L. and 90% C.L. curves in the S and T plane determined by the global fit to all data by Peskin and Takeuchi.[14] Also shown are the predictions for the standard model and two dynamical symmetry breaking models, as a function of m_t and M_H. The values of m_t are shown by the crosses which, starting from the bottom, increase by 20 GeV from $m_t = 90$ GeV.

data to, was obtained by "scaling up" QCD and there is no guarantee that this is necessarily applicable for the "correct" theory of dynamical symmetry breaking. Furthermore, even in QCD, one knows that the Weinberg sum rules are rather slowly convergent,[21] so a saturation of the integral in Eq. (17) solely by low energy states may not be totally accurate. This lack of convergence may, in fact, be even more of a problem in the more sophisticated versions of Technicolor, so called Walking Technicolor, which are needed to address the issue of fermion masses. I will briefly discuss some aspects of Walking Technicolor theories at the end of this lecture, but I note here only that precisely because of the slow running of $\alpha_{WTC}(q^2)$ with q^2 one can argue[22] that it is likely that $S_{WTC} < S_{TC}$. In fact, in the model analysis of Hsu and Sundrum,[22] values of $S_{WTC} \sim (\frac{1}{4} - \frac{1}{10})S_{TC}$ were obtained

4. Signals for dynamical symmetry breaking far above threshold

As the above discussion demonstrates, it is difficult to learn much about the symmetry breaking sector even from the very precise data obtained at LEP, mostly because LEP is really a low energy machine! Things will be considerably amelio-

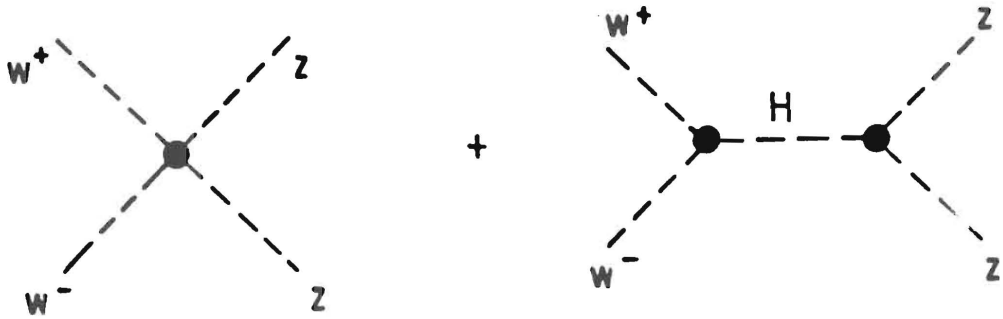

Figure 2: Feynman graphs contributing to the process $w^+w^- \to zz$ in the purely Higgs theory

rated at the LHC and the SSC, because one can begin to probe TeV subenergies in the $q\bar{q}$ subsystem. However, even here, despite the stupendous CM energies of the LHC and the SSC one is not very far above threshold for probing W_LW_L scattering, which is so crucial to ascertain whether the symmetry breaking sector is that of a weakly coupled or a strongly coupled theory. Only at even higher energies, those of the Eloisatron and a fantomatic Mega TeV machine, will the distinction become really manifest.

In this section I want to examine W_LW_L scattering far above threshold, to learn what are the characteristics which distinguish between having a weakly coupled or a strongly coupled symmetry breaking dynamics. By the equivalence theorem,[9] for these purposes it suffices to study the dynamics of the associated Goldstone bosons. As we discussed earlier, at threshold the behaviour of the scattering amplitudes for Goldstone bosons is solely fixed by the coset space of the breakdown (e.g. $SU(2)_L \times SU_R(2)/SU(2)_{L+R}$) and not by the specific dynamics of the theory. For instance for the process $w^+w^- \to zz$, in the pure Higgs theory, the evaluation of the Feynman graphs of Fig. 2 gives the scattering amplitude

$$\mathcal{A}(w^+w^- \to zz) = -2\lambda\left\{1 + \frac{2\lambda v^2}{s - 2\lambda v^2}\right\} \tag{24}$$

This amplitude agrees at threshold, $s << 2\lambda v^2$, with that of Eq. (12) obtained by considering only the dynamics of the nonlinearly coupled Goldstone bosons

$$\mathcal{A}(w^+w^- \to zz) \to \frac{s}{v^2} \tag{25}$$

So, obviously, nothing can be learned about the nature of the mechanism for symmetry breakdown precisely at this point.

Above the threshold region, however, there are real distinctions between the scattering amplitude for $w^+w^- \to zz$ predicted by the Higgs theory and that

precisely the same pattern to emerge in a dynamically broken $SU(2) \times U(1)$ theory. The underlying theory which governs the interactions of the technipions, which eventually become the longitudinal gauge bosons, should manifest both **resonance formation** and **multiparticle production**, with the former process being more prevalent at lower CM energies in the scattering of technipions, and the latter processes dominating at higher CM energies.

Hadronic machines like the SSC and LHC are primarily suited to probe the hard scattering of quarks and gluons. One can, of course, also probe the scattering of virtual weak gauge bosons in these machines, but the relevant accessible CM subenergy is much degraded since the $W's$ only carry a fraction of the energy of the quarks, which themselves only carry a fraction of the energy of the initial protons. As a result the invariant subenergy effectively available for $W_L W_L$ scattering at the SSC and LHC is small, and even searching for resonance formation in the $W_L W_L$ scattering process is challenging. This is particularly so because resonance detection in the $W - W$ channel is hampered by a copious background of non resonant W pairs produced by quark annihilation, $q\bar{q} \to WW$. Furthermore, experimentally one must isolate the W signal through their decay into leptons, which effectively reduces much further the number of events one can study. Nevertheless, there is broad agreement that, although difficult, a resonant signal in the $W_L W_L$ channel with resonance mass in the range $2 - 4$ TeV is observable.[23] For instance, a $\rho_{TC}^{\pm}$ with mass of 1.8 TeV, decaying into $W^{\pm}Z^0$, for a total integrated luminosity of $10 fb^{-1}$ gives a signal of 30 trilepton events $\ell^{\pm}\ell^{+}\ell^{-}$, from the further decays of the $W^{\pm}Z^0$, compared to a background of about 9 events, so reduced by appropriate cuts. For a $\rho_{TC}^{\pm}$ of mass of 4 TeV, the signal is down to 13 events, but the background stays the same, so this effect is just marginally detectable.[23]

At the LHC and SSC the available subenergy for $W_L W_L$ scattering is too small to hope to see the opening up of inelastic channels, to be expected if there is strong dynamics in the symmetry breaking sector. However, if one had a machine of sufficient energy, like the Eloisatron, it may be that the signal provided by this second aspect of strong dynamics could be simpler to detect than the presence of resonances in the $W_L\ W_L$ channel. In collaboration with Duncan Morris, and Rogerio Rosenfeld, I have recently investigated this issue at UCLA.[24] What we imagined is that, above a certain threshold energy $\sqrt{s_0}$, instead of having purely resonant dominated, elastic $W_L\ W_L$ scattering, $W_L\ W_L$ scattering would give rise predominantly to multiparticle production: $W_L\ W_L \to nW_L$ $(n > 2)$.

If our experience with hadronic scattering is any guide, $\sqrt{s_0}$, for $W_L W_L$ scattering will not be, comparatively, a terrifically high energy. Typically in strong interaction physics, multiparticle production begins to dominate at around $\sqrt{s_0} \simeq 2\ GeV$, which for a Technicolor theory, with the scaling implied by Eq. (21), gives a threshold $\sqrt{s_0} \simeq 4 - 5$ TeV. Furthermore, at $\sqrt{s_0} \simeq 2\ GeV$ in hadron physics the average multiplicity of produced states, which are dominantly pions, is already quite high. Typically $< n >_{\text{charged}} \simeq 4$ and $< n >_{\text{total}} \simeq 6$. So, analogously, we are

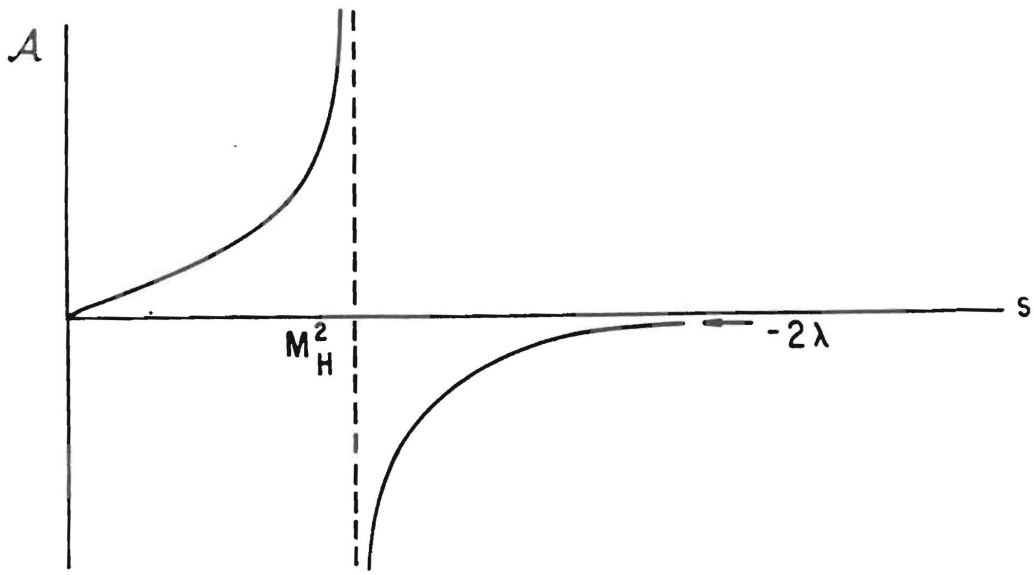

Figure 3: Behaviour of the amplitude $\mathcal{A}(w^+w^- \to zz)$ as a function of s for the Higgs model. The amplitude displays resonance behaviour as $s \to M_H^2$ and goes to an asymptotic value $\mathcal{A} \to -2\lambda$.

predicted by some model of dynamical symmetry breakdown. The amplitude (24) for the Higgs case, when plotted as a function of s, displays only one remarkable feature, a resonance pole at the Higgs mass $M_H^2 = 2\lambda v^2$, as shown in Fig. 3. Of course, even in this case, there is not a real pole in $\mathcal{A}$ because as $s \to M_H^2$ the result (24) ceases to be valid since one must include the effect of higher order in λ graphs which build up the Higgs width. For the case of dynamical symmetry breakdown, on the other hand, in complete analogy to what happens in the known strongly interacting theory which describes $\pi - \pi$ interactions (QCD), one expects that $\mathcal{A}$ should contain resonances in other partial waves (like the f_2 in $\pi^+\pi^- \to \pi^0\pi^0$), besides the $J = 0$ Higgs boson:

$$\mathcal{A} = 32\pi \sum_J (2J+1) P_J(\cos\Theta) a_J(s) \quad . \tag{26}$$

Secondly, again in analogy to $\pi - \pi$ scattering, one expects the opening up of inelastic channels, so that the partial waves $a_J(s)$ should not have unit strength

$$a_J(s) \;=\; \frac{1}{2i}\{\eta_J(s) e^{2i\delta_J(s)} - 1\} \;\; ; \;\; \eta_J(s) \;<\; 1 \tag{27}$$

In $\pi - \pi$ scattering, as in other strongly interacting scattering processes, resonance formation is the dominant process at lower energies. Eventually, however, when the CM energy is high enough the principal feature of the scattering process is the copious opening up of inelastic channels. Indeed, at high energy, the dominant feature of hadronic interactions is multiparticle production, which is mostly that of pions, since these particles are the lightest hadrons. One should expect

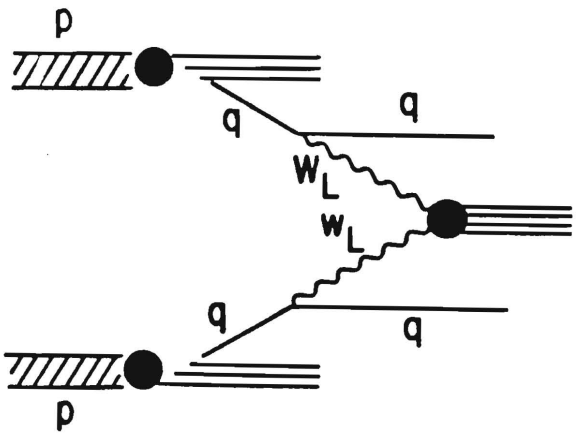

Figure 4: W_L W_L scattering arising from quark bremsstrahlung

led to expect that if we can access energies of the order of $4-5$ TeV in the $W_L W_L$ CM system, the average multiplicity of produced $W's$ should also be given by similar numbers. The existence of copious multi W production at these energies is a clear signal that one has a strongly interacting symmetry breaking sector, for otherwise a process like $W_L W_L \rightarrow 6W_L$ would be hopelessly suppressed by the weakness of the electroweak coupling constant.

Morris, Rosenfeld and I[24] computed the cross section for $W_L W_L$ scattering due to the quark bremsstrahlung process of Fig. 4 in the effective W approximation.[25] In this approximation, the contribution of $W_L W_L$ scattering processes in a pp scattering process, of total center of mass energy $\sqrt{s}$, is given by

$$\sigma_{pp}^{W_L}(s) \;=\; \int_{\tau_{\min}} \mathcal{L}_{W_L W_L}(\tau)\; \sigma_{W_L W_L}(\hat{s} = \tau s) \tag{28}$$

Here $\sqrt{\hat{s}}$ is the CM subenergy of the $W_L W_L$ system and $\mathcal{L}_{W_L W_L}(\tau)$ is the effective W_L W_L luminosity, for producing a $W_L W_L$ system of a given invariant mass $\hat{s} = \tau s$.

The $W_L W_L$ luminosity function is readily calculable as the square of the convolution of the probability of finding a W_L in a given quark $P_{W_L q}$, times the quark distribution function in the proton

$$\mathcal{L}_{W_L W_L} = (f_q \bigotimes P_{W_L q})^2 \tag{29}$$

The result of this calculation is shown in Fig. 5, plotted as a function of τ. As expected, this is a sharply falling function of τ, as τ increases towards unity. For a fixed value of $\sqrt{s}$, it is always most favorable to fuse together two W_L $'s$, with the smallest possible invariant mass. Wanting to access high values of $\sqrt{\hat{s}}$, to see the opening up of multi W_L production, forces τ to larger values with a consequent decrease in the effective $\mathcal{L}_{W_L W_L}$ luminosity. For a putative threshold of $\sqrt{s_0} =$ 5 TeV, the minimum value of τ reached at the LHC ($\sqrt{s} = 16$ TeV)

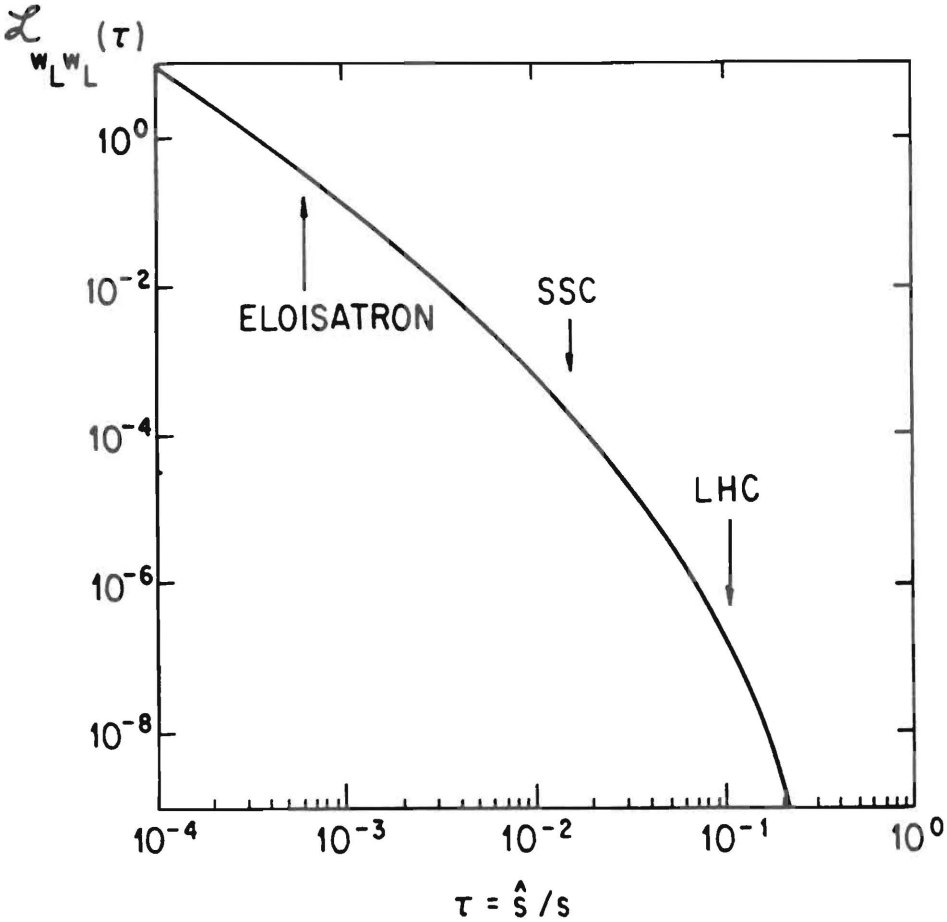

Figure 5: Effective $W_L W_L$ luminosity, plotted as a function of $\tau = \hat{s}/s$, from.[24] Also indicated are the values of $\tau_{\min}$ (for $\sqrt{\hat{s}_0} = 5\ TeV$) reachable at the LHC, SSC and Eloisatron.

SSC ($\sqrt{s} = 40\ TeV$) and Eloisatron ($\sqrt{s} = 200\ TeV$) are, respectively, $\tau_{\min} = 9.77 \times 10^{-2}$, 1.56×10^{-2}, 6.25×10^{-4}. These minimum values, which correspond to the maximal $W_L W_L$ luminosity available in each machine for producing $W_L W_L$ interactions with at least $\sqrt{\hat{s}_0} = 5\ TeV$, are also indicated in Fig. 5.

Because we are only interested in the signal for multiparticle production in $W_L\ W_L$ scattering, it is perfectly sensible to entirely neglect any $W_L\ W_L$ scattering below $\sqrt{\hat{s}} = \sqrt{\hat{s}_0}$. Above this threshold, on the other hand, we assume that the $W_L\ W_L$ scattering is dominated by multiparticle production with a constant total cross section. That is, the simple model adopted[24] is that

$$[\sigma_{W_L W_L}(\hat{s})]^{\text{Multi}} = \theta(\hat{s} - \hat{s}_0)\sigma_0 \tag{30}$$

Scaling from $\pi - \pi$ scattering, using Eq. (21), suggest a value $\sigma_0 \simeq 1\ nb$ for the $W_L\ W_L$ multiparticle cross section. Using the model of Eq. (30) in Eq. (28) gives for the component of the pp scattering cross section which involves multi W_L production the formula

$$[\sigma_{pp}^{W_L}(s)]^{\text{Multi}} = \sigma_0 \int_{\tau_{\min}}^{1} d\tau \mathcal{L}_{W_L\ W_L}(\tau) \simeq \frac{2\tau_{\min}}{3}\sigma_0 \mathcal{L}_{W_L\ W_L}(\tau_{\min}), \tag{31}$$

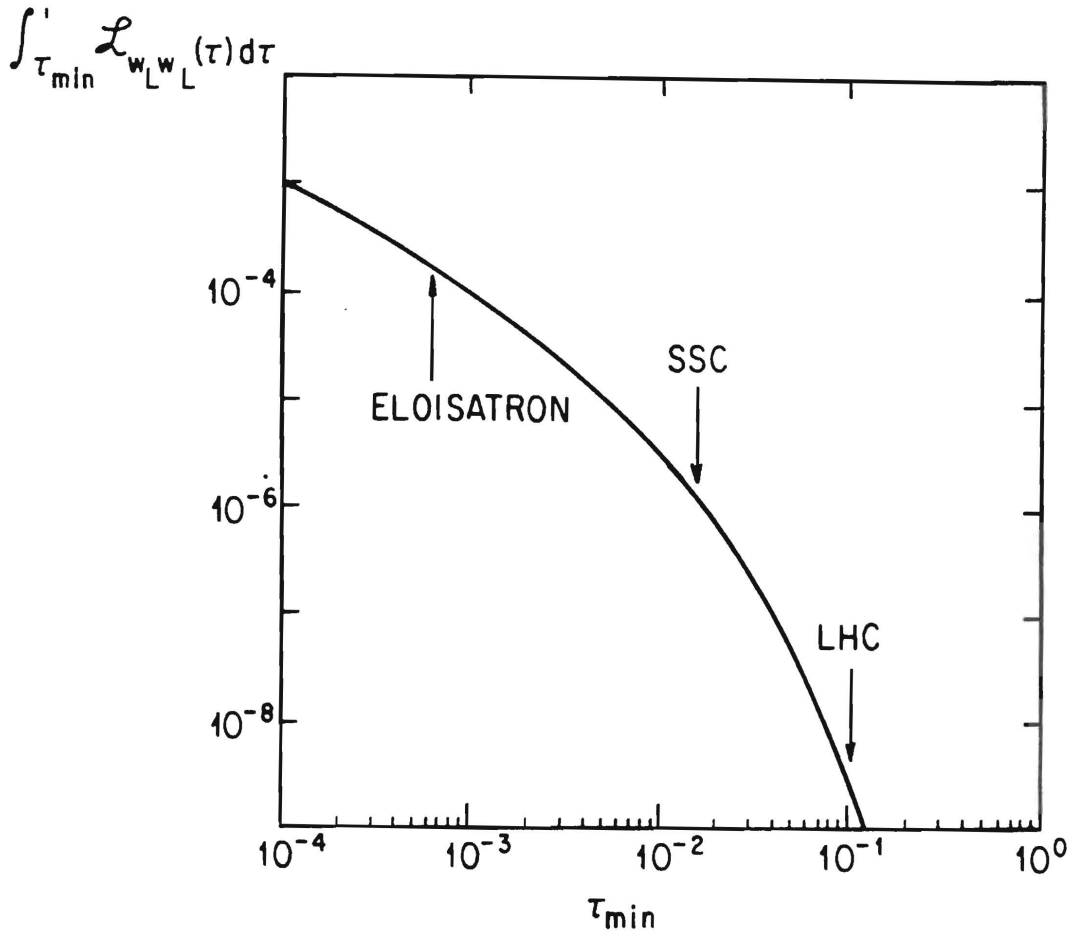

Figure 6: Plot of $\int_{\tau_{min}}^{1} \mathcal{L}_{W_L\ W_L}(\tau)d\tau$ as a function of $\tau_{\min}$. The arrows indicate the value of $\tau_{\min}$ for the Eloisatron, SSC and LHC, assuming $\sqrt{\hat{s}_0} = 5\ TeV$.

where the last line uses the fact that the luminosity function decreases approximately as $\tau^{-5/2}$ for small τ. For the same amount of integrated luminosity, it is clear from Eq. (31) that to see this phenomena one really needs the energy of the Eloisatron, with the number of events scaling roughly as the CM energy of the machine cubed. This is demonstrated explicitly in Fig. 6 where the integral of $\int_{\tau_{\min}}^{1} \mathcal{L}_{W_L W_L}(\tau)$ is given as a function of $\tau_{\min}$. Assuming a given threshold value for multi W_L production, the value of this integral for the Eloisatron, SSC and LHC can be read off from the corresponding $\tau_{\min} = \hat{s}_0/s$ in the figure. The arrows in Fig. 6 for these accelerators assume as a threshold $\sqrt{\hat{s}_0} = 5\ TeV$.

The assumptions made concerning multi W_L production, Eq. (30), are rather conservative, in that both σ_0 could be larger and the threshold energy $\sqrt{\hat{s}_0}$ could be smaller. Nevertheless, even in this case, assuming an equally conservative yearly integrated luminosity of $\int \mathcal{L}dt = 10^{40} cm^{-2} = 10 fb^{-1}$ one gets about 2000 multi W events per year.[24] These events will mostly contain $< n >_{\text{total}} \simeq 6\ W_L$, if the W_L multiparticle distributions are rather narrowly peaked around $< n >$ as are the pion distributions. Fig. 7, taken from the paper of Morris, Rosenfeld and I,[24] shows a guesstimate of what the distribution of W_L events as a function of n might look like. Although the tails are not large, there are of 0(100) events expected with 8 W_L and a handful of events containing 12 W_L 's. Just for fun, one can also try

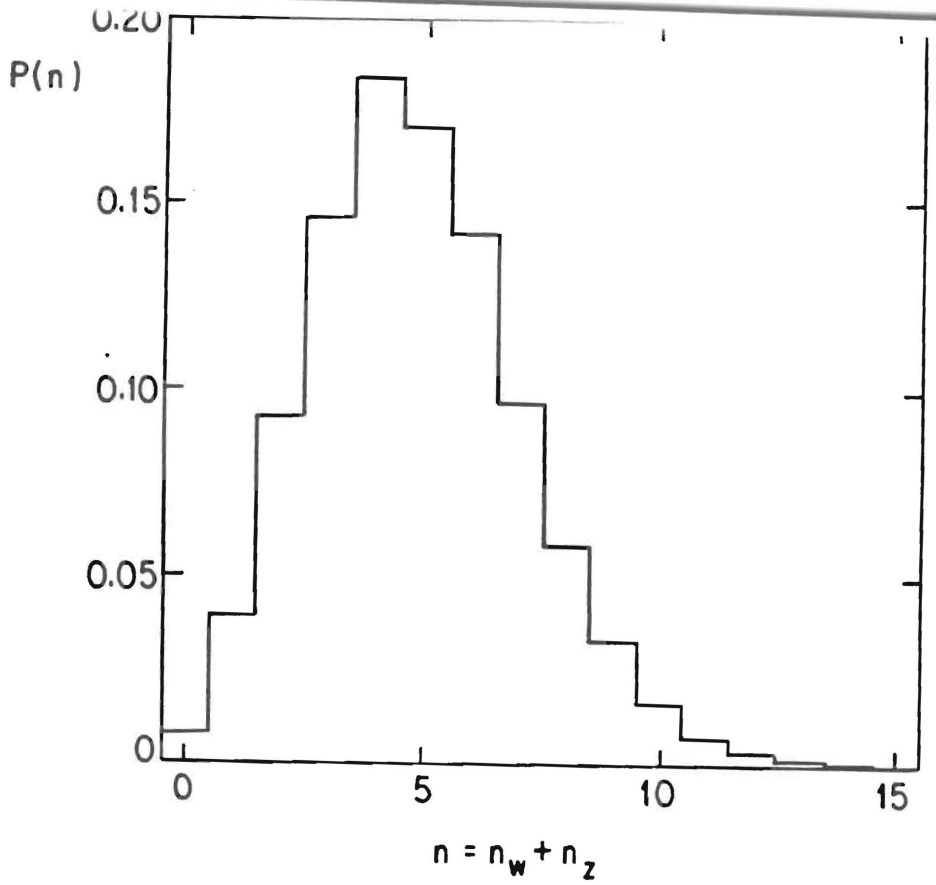

Figure 7: Guesstimates for the distribution of multi W_L events at the Eloisatron. From.[24]

to estimate the corresponding number of events at a Mega TeV machine (a good name for it might be Megatron, or even better, Megaton!). This number should be in the neighborhood of 10^{12} multi W_L events, still mostly containing 6 or so W_L - the threshold value of $< n >_{\text{total}}$ - but now events with large numbers of W_L (20 or above) in the tails are in the thousands!

Although multi W signals are characteristic of a theory having strong W_L W_L interactions, their observation (based solely on the number of events expected) only appears to begin to be possible at the Eloisatron. Even then, it is not at all clear that one can really dig out this signal from the background. Because the number of events is small it appears to be necessary to use the multi jet nature of the events, coming from hadronic decays of the $W's$, as a signature combined with the presence of multileptons. Multiple top production is then a natural background, but fortunately, one does not expect an overwhelming number of events containing 3 or more $t\bar{t}$ pairs.[24] Clearly before one can make any more concrete statements about the observability of this signal, a rather thorough phenomenological analysis will be necessary. Some aspects of this analysis are contained in ref. (24). There we show that events containing 2 high p_T leptons coming from W decay and ten or more jets, or one leptonically reconstructed Z with one or more high p_T lepton and eight or more jets, have a combined cross section of a few fb. These events, furthermore, are a factor of 3-5 above background.

5. The Messy problem of fermion mass

Fermion masses are the unresolved crux in models of spontaneous symmetry breakdown. Because this issue is so central, and so difficult to resolve in a

satisfactory fashion, it is important to contrast the "solutions" which have been put forward in these models with the way fermion masses are generated in the standard model.[5] Because of the presence of non trivial Yukawa couplings between the Higgs doublet field and the quark and lepton fields, in the standard model the Higgs doublet vacuum expectation value $< \Phi >= \frac{1}{\sqrt{2}}v$ serves to give mass to all fermions. Schematically, the fermion mass matrix is given by

$$M_f = \Gamma_f < \Phi > \quad , \tag{32}$$

where the Γ_f are the Yukawa couplings. Because the Yukawa couplings are arbitrary, although the masses are proportional to the Fermi scale v, the masses themselves are not predictable in the model.

The above negative feature of the standard model is offset by two very nice properties. Because the Yukawa couplings Γ_f have no symmetry reason to be diagonal, the fermion mass matrix M_f will, in general, be non diagonal. The diagonalization of M_f, necessary to go to a physical basis, then introduces naturally into the theory a mixing matrix for the quarks - the Cabibbo Kobayashi Maskawa[26] matrix V_{CKM}. Furthermore, if only one Higgs doublet contributes to M_f, as in Eq. (32), then diagonalizing M_f necessarily also diagonalizes the Yukawa couplings Γ_f. Thus, in the standard model, Higgs exchange never generates any flavor changing neutral currents (FCNC), in complete accord with experiment where one knows that these processes are very suppressed.[27]

Fermion masses cannot be generated by Technicolor models[8] without introducing further structure. Although technifermion condensates $< \bar{T}T >$ can serve to break down $SU(2) \times U(1)$ - thus playing an analogous role to the Higgs vacuum expectation value $< \Phi >$ - to generate mass for the fermions it is necessary to be able to connect the technifermions to the quarks and leptons. That is, one must be able to introduce something analogous to the Yukawa couplings Γ_f of the standard model.

The connection between ordinary fermions and technifermions can be achieved in a variety of ways. The most popular approach, called Extended Technicolor (ETC), is to introduce further gauge interactions under which ordinary fermions and technifermions transform irreducibly.[28] Another idea is that there are some residual interactions between these states because both the quarks and leptons and the technifermions are bound states of some still more elementary fermions.[29] Here, for concreteness, I will concentrate only on the ETC suggestion describing briefly both the idea, some of its problems and some of their putative solutions in terms of Walking Technicolor models.[30]

In ETC models[28] one introduces a further non Abelian gauge theory (ETC) besides Technicolor. Both quarks and leptons and the technifermions are presumed to transform according to some irreducible ETC representation. Furthermore, the ETC gauge theory is assumed to be broken spontaneously at a scale Λ_m. If we let F denote generically the fermions ETC multiplets, then the exchange of an ETC

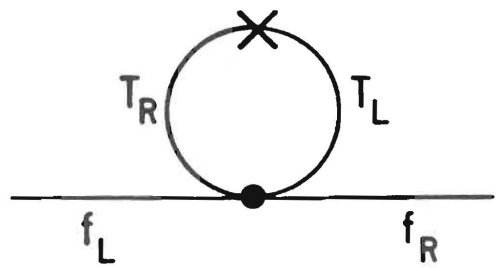

Figure 8: Mass generation in ETC models

gauge boson will give rise to effective interactions among these fermions of the form

$$\mathcal{L}_{\rm ETC} \sim \frac{1}{\Lambda_m^2}\bar{F}_L\lambda_a\gamma_\mu F_L\bar{F}_R\lambda_a\gamma^\mu F_R \tag{33}$$

As a result of the formation of technicolor condensates $< \bar{T}_L T_R > \sim \Lambda_{TC}^3 \simeq v^3$ the interaction in Eq. (33) will give rise to mass terms for the ordinary fermions. This process is shown schematically in Fig. 8. and the magnitude of the dynamically generated mass is clearly of the order

$$m_f \sim \frac{< \bar{T}_L T_R >}{\Lambda_m^2} \sim \frac{v^3}{\Lambda_M^2} \tag{34}$$

Unfortunately, the interaction (33) will generate not only masses for the fermions f, but once fermion mixing is taken into account, also FCNC among these fermions. To avoid too large FCNC one needs to make the ETC dynamical scale rather large,[31] typically $\Lambda_m \sim 10^2 - 10^3\ TeV$. However, if Λ_m is so large then it is very difficult to generate via Eq. (34) large masses for the fermions, since this equation gives masses in the MeV not GeV range!

The apparent incompatibility of having sizable fermions masses and small FCNC interactions is not the only problem of ETC models. In general, it is extremely difficult to construct any realistic models. Furthermore, even semirealistic models - like the Farhi Susskind model[32] - are plagued by the presence of very light pseudo Goldstone bosons arising from the formation of the Technicolor condensates.[33] All of these problems have, rather naturally, discouraged people from thinking about dynamical symmetry breaking theories. However, as Holdom and others have pointed out,[30] many of the above problems may in fact not arise in more realistic theories. They are only problems for theories which, in their dynamics, resemble rather closely QCD. As we saw already, when we considered the S parameter, such theories run into other experimental difficulties purely related to the weak interactions.

In QCD the expectation value of u - quarks in the vacuum and the pion decay constant, f_π are closely related, since they both are determined by the dynamical scale of QCD, Λ_{QCD}:

$$< \bar{u}u > \sim \Lambda^3_{QCD} \; ; \; f_\pi \sim \Lambda_{QCD} \tag{35}$$

If this also obtains in the underlying Technicolor theory then, most likely, not only will one run into trouble experimentally with the S parameter, but also any attempts to build realistic ETC models will fail - for the reasons discussed above. It is therefore important to try to understand under what conditions the QCD relation (35) may be obviated by the dynamics of the underlying theory one is using to trigger the spontaneous breakdown of the electroweak $SU(2) \times U(1)$ theory.

For these purposes it is important to note that f_π and $< \bar{u}u >$ are, in principle, quite distinct parameters. The pion decay constant typifies the strength of the coupling of the Goldstone pions to the broken axial current J^μ_A:

$$i f_\pi k^\mu = < \pi(k)|J^\mu_A|0 > \quad . \tag{36}$$

$< \bar{u}u >$, on the other hand, is the order parameter for the breaking of the chiral symmetry associated with the currents J^μ_A in QCD. Both f_π and $< \bar{u}u >$, however, can be related to the quark self energy $\Sigma(p)$, defined as usual as the part of the full inverse quark propagator, $S^{-1}(p)$, which is not proportional to $\gamma \cdot p$

$$S^{-1}(p) = \gamma \cdot p A(p^2) + \Sigma(p) \tag{37}$$

Graphically, the relation of $< \bar{u}u >$ and f^2_π to $\Sigma(p)$ can be easily inferred from Fig. 9. Retaining only the leading momentum behaviour one sees that

$$< \bar{u}u > \quad \sim \int d^4p \, \frac{\Sigma(p)}{p^2} \quad \sim \int dp^2 \Sigma(p) \tag{38a}$$

$$f^2_\pi \quad \sim \int d^4p \, \frac{\Sigma^2(p)}{(p^2)^2} \quad \sim \int \frac{dp^2}{p^2} \Sigma^2(p) \tag{38b}$$

The presence of a nontrivial quark self energy $\Sigma(p)$ is a signal of the breaking of the chiral symmetry present in the absence of quark masses. Thus $\Sigma(0)$ serves as a chiral breaking parameter and naturally one expects in QCD that $\Sigma(0) \sim \Lambda_{QCD}$. One can show in QCD[34] further that, up to logarithmic factors, for large momentum

$$\Sigma(p) \sim \Sigma^3(0)/p^2 \quad . \tag{39}$$

Because of the different weighting of the p^2 integrals associated with $< \bar{u}u >$ and f^2_π, one sees that f^2_π probes the region of the quark self energy near $p^2 = 0$, while

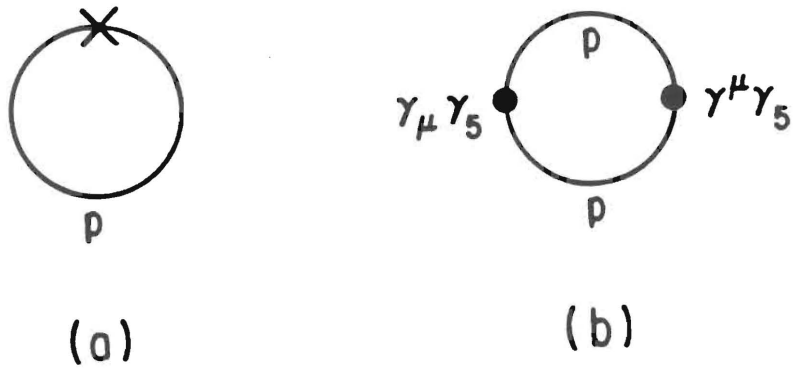

Figure 9: Graphs relating $< \bar{u}u >$ (a) and f_π^2 (b) to the quark self energy $\sum(p)$.

$< \bar{u}u >$ feels more the large momentum structure of $\Sigma(p)$. Using Eq. (37), and (38) one deduces immediately that, up to logarithmic factors,

$$< \bar{u}u > \quad \sim \int dp^2 \sum(p) \quad \sim \Sigma^3(0) \tag{40a}$$

$$f_\pi^2 \quad \sim \int \frac{dp^2}{p^2} (\Sigma(p))^2 \quad \sim \Sigma^2(0) \tag{40b}$$

The above natural relationship between $< \bar{u}u >$ and f_π need not apply in theories where the self energy of the underlying fermions in the important region for mass generation falls much slower with p^2 than does the quark self energy in QCD., This is the basic idea behind Walking Technicolor theories.[30] Although purely on dimensional grounds $\Sigma(p) \sim \Sigma^3(0)/p^2$ at large p^2, the region of p^2 important for fermion mass generation due to the ETC interactions (33) is $p^2 \sim \Lambda_m^2$. It may well be that for a realistic Technicolor theory at these large, but not yet asymptotic, values of p^2, the techniquark self energy is much larger than its asymptotic value. Although the ordinary fermion masses are still given by Eq. (34), the second approximate equality in this formula no longer necessarily applies. Rather, more properly, one has

$$m_f \sim \frac{< \bar{T}T >}{\Lambda_m^2} \sim \frac{1}{\Lambda_m^2} \int^{\Lambda_m^2} dp^2 \Sigma(p) \sim \Sigma(\Lambda_m) \tag{41}$$

$\Sigma(\Lambda_m)$ is not known a priori without a dynamical calculation. Of particular interest is its relation with the technipion decay constant v, which is the scale that is related directly to the gauge boson masses. This scale, as is f_π in QCD, is connected [cf Eq. (40)] to the technifermion self energy at zero momentum transfer

$$v \simeq \Sigma(0) \tag{42}$$

Using (42), it is clear that, if at scales of $0(\Lambda_m)$ the technifermion self energy has

already approached its asymptotic value (39), then

$$\Sigma(\Lambda_m) \sim \frac{\Sigma^3(0)}{\Lambda_m^2} \sim \frac{v^3}{\Lambda_m^2} \tag{43}$$

and one recovers the estimate made previously in Eq. (34). The assumption one makes in Walking Technicolor theories,[30] however, is that the dynamics is such that $\Sigma(\Lambda_m)$ is still much above this value. That is

$$[\Sigma(\Lambda_m)]_{WTC} >> \frac{v^3}{\Lambda_m^2} \tag{44}$$

To obtain a behaviour like (44) it is necessary that all physical quantities evolve slowly with momentum - hence the moniker walking for these theories. This can be achieved if the underlying theory is very nearly not asymptotically free, or, perhaps, a fixed point theory.[30] I do not have the time, or the relevant expertise, to enter into the details of the dynamics of these theories here. However, I note that (44) neatly resolves one of the difficulties of ETC models. Namely, one can have Λ_m be sufficiently large to avoid trouble with FCNC and yet get fermion masses which are sufficiently large. Furthermore, since the fermion masses are given by $\Sigma(\Lambda_m)$ and this self energy can be different for different technifermions[35] one can begin building in this way some of the observed hierarchies in the masses of the quarks and leptons.

6. Concluding Remarks

Although Walking Technicolor models resolve some of the dynamical conundrums of ETC models, they are not a panacea and many issues are still unclear. For instance, despite intense activity in the last few years no real full models are yet in sight.[36] Furthermore, even in semirealistic models the dynamics and group structure introduced are in general quite complex.[37] In addition, many problems continue to be difficult to tackle. For example, although it is possible to build a hierarchy between m_t and m_b, it is difficult to do so without distabilizing by the dynamics the result that $\rho \simeq 1$.

These comments not withstanding, it is important not to let the difficult issue of fermion masses discourage one from contemplating some of the possible physical consequences of theories where the symmetry breaking of $SU(2) \times U(1)$ is done more dynamically. After all, it could well be that fermion mass generation really requires a new idea. Much in the same way in which quark models of hadrons were dynamically untenable before the discovery of asymptotic freedom, it could well be that to understand fermion masses we will need to break away from our present conventional mode of thinking.

Because fermion masses most probably arise due to new physics at a scale $\Lambda_m >> v$, it is important to try to explore the consequences of dynamical symmetry breaking in the Eloisatron energy range, which lies between Λ_m and v. Because

of the energy degradation that occurs when one focuses on the W_L constituents of the colliding protons [cf Fig. 4], for the telltale $W_L W_L$ scattering processes high energy machines like the LHC and SSC are only threshold machines for probing dynamical symmetry breaking. However, it is quite likely that new and interesting phenomena may begin to open up at Eloisatron energies (and would be spectacular at $M = 10^6\ TeV$!) As we discussed, beside the presence of characteristic resonances in the $W_L W_L$ channel, one should expect to see the first signs that the $W_L W_L$ total cross section remains approximately constant through the opening of inelastic channels. The discovery of multi W_L events ($n_{W_L} > 6 - 8$) at the Eloisatron would be a spectacular confirmation of the existance of a new strongly interacting sector in the physics of fundamental particles.

7. Acknowledgement

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