LBL-22841 c.2

Lawrence Berkeley Laboratory
UNIVERSITY OF CALIFORNIA

Physics Division

Invited talk presented at the 1987 American Physical Society Annual Meeting, Division of Particles and Fields, Salt Lake City, UT, January 14-17, 1987

PROBING ELECTROWEAK SYMMETRY BREAKING AT THE SSC: A NO-LOSE COROLLARY

M.S. Chanowitz

February 1987

Prepared for the U.S. Department of Energy under Contract DE-AC03-76SF00098

DISCLAIMER

February 11, 1987 LBL-22841

Probing Electroweak Symmetry Breaking at the SSC: A No-Lose Corollary *

Michael S. Chanowitz

Lawrence Berkeley Laboratory
University of California
Berkeley, California 94720, U.S.A.

Invited talk presented at the 1987 Annual Meeting of the Division of Particles and Fields of the American Physical Society, Salt Lake City, January 14-17, 1987. To be published in the proceedings.

*This work was supported by the Director, Office of Energy Research, Office of High Energy and Nuclear Physics, Division of High Energy Physics of the U.S. Department of Energy under Contract DE-AC03-76SF00098.

Probing Electroweak Symmetry Breaking at the SSC: A No-Lose Corollary *

Michael S. Chanowitz

Lawrence Berkeley Laboratory
University of California
Berkeley, California 94720, U.S.A.

Abstract

Low energy theorems are derived for scattering of longitudinally polarized W and Z's, providing the basis for an estimate of the observable signal at the SSC if electroweak symmetry breaking is due to new physics at the TeV scale.

1. INTRODUCTION

This is not a general review of SSC physics but focuses on a particular issue—electroweak symmetry breaking—which is central to the motivation for constructing the SSC. There are other excellent reasons for the SSC: the possibility of completely unanticipated discoveries that could be more important than any of the things we are able to imagine and the theorists' wish list of possible new physics such as additional gauge bosons, further generations of matter, supersymmetry, *etc.*... The energy and luminosity of the SSC will make it a tremendous instrument for these searches, which are in most cases easier than the rather difficult physics of electroweak symmetry breaking I will discuss. But while these other topics are exciting possibilities, electroweak symmetry breaking is a certainty: we know that the W and Z bosons have masses and that the photon is massless, but we do not know why. This is the outstanding open question that must be answered to complete the highly successful unified theory of the weak and electromagnetic interactions.[1] I will argue that with the proposed energy, $\sqrt{s}$ =40 TeV, and luminosity, $\mathcal{L} = 10^{33}$ cm.$^{-2}$ s.$^{-1}$, the SSC is certain to see the manifestations of the new physics responsible for electroweak symmetry breaking.[2] The SSC is strategically placed to see this new physics: I would not be prepared to make the above statement for a machine with one tenth the luminosity, $\sqrt{s}$, $\mathcal{L}$ = 40, 10^{32} or for one with half the energy, $\sqrt{s}$, $\mathcal{L}$ = 20, 10^{33}.

*This work was supported by the Director, Office of Energy Research, Office of High Energy and Nuclear Physics, Division of High Energy Physics of the U.S. Department of Energy under Contract DE-AC03-76SF00098.

Within the general framework of spontaneous symmetry breaking—the only known way of constructing a sensible broken gauge theory—we believe that the W and Z masses must result from new, unknown particles that interact by a new unknown force. We know neither the mass scale M_{SB} of the new particles ("SB" for symmetry breaking) nor the strength λ_{SB} of the new force. With the SSC we will be able to determine whether the new force is a strong interaction

$$\frac{\lambda_{SB}}{8\pi} \cong 0(1) \tag{1.1}$$

or a weak interaction

$$\frac{\lambda_{SB}}{8\pi} \ll 1. \tag{1.2}$$

In the first case I will argue that the new physics lies above 1 TeV. In this case it is also likely that the new spectrum begins below 2 TeV and will be directly observable. In the second, weak coupling case the new particles are much lighter than the 1 TeV scale, *i.e.*, a few hundred GeV or below, and they will be copiously produced at the SSC. These statements taken together are what I call the "No-Lose Corollary".

The basic physical point is that the longitudinal polarization modes of the W and Z, denoted W_L and Z_L, are actually degrees of freedom that originate, by the Higgs mechanism, in the symmetry breaking sector. They are essentially the "pions" of the symmetry breaking sector, and like the pions of hadron physics they obey low energy theorems characteristic of the scattering of Goldstone bosons.[3] If there are no other light (compared to 1 TeV) particles in the symmetry breaking sector, the low energy scattering amplitudes depend only on the known parameters G_F and $\rho = (M_W/M_Z\cos\theta_W)^2$ and not at all on the unknown physics of the symmetry breaking sector, denoted by its lagrangian $\mathcal{L}_{SB}$.[4] For example, one of the low energy amplitudes is

$$\mathcal{M}(W_L^+W_L^- \to Z_LZ_L) = \sqrt{2}G_F\frac{s}{\rho}. \tag{1.3}$$

The low energy theorem provides the correlation between the mass scale M_{SB} and the interaction strength λ_{SB}. Unitarity requires that the amplitude cannot be proportional to s for arbitrarily large s, and the most likely scenario, discussed below, is that the growth in s is cut off at the mass scale M_{SB}. This observation allows us to correlate the strong coupling regime, eq. (1.1), with the mass domain $M_{SB} \gtrsim 1$ TeV (*c.f.* section 3).

For the SSC the crucial observation is that strong W_L, Z_L scattering is observable at the SSC by virtue of increased yields of gauge boson pairs produced with the WW fusion mechanism[5] shown in figure 1. At the design energy and luminosity these extra gauge boson pairs will be observable above the background sources of gauge boson pairs that are present whether $\mathcal{L}_{SB}$ is a strong or weak coupling theory. The conclusion is the

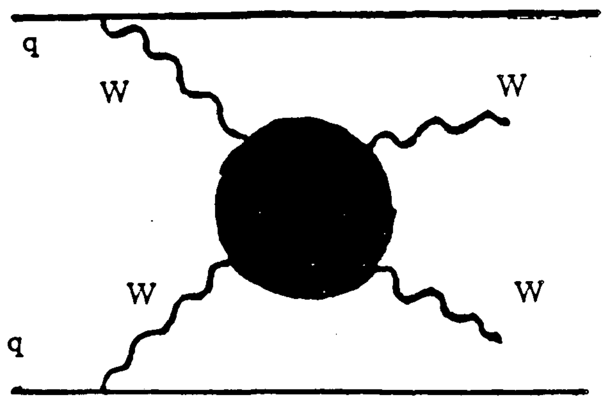

Figure 1: Production of W pairs by WW fusion.

No-Lose Corollary:

Either there are light ($\ll$ 1 TeV) particles from $\mathcal{L}_{SB}$ that can be produced and studied directly

and/or

Excess WW, WZ, ZZ production is observable, signaling strongly-coupled $\mathcal{L}_{SB}$ with $M_{SB} \gtrsim$ 1 TeV.

For the strong coupling case, if as in hadron physics resonances occur when the partial wave amplitudes are 0(1), then probably $M_{SB} \lesssim$ 2 TeV and the low-lying spectrum of $\mathcal{L}_{SB}$ is (just) visible.

In the weak coupling case there is one known exception that should be mentioned: if $\mathcal{L}_{SB}$ is given by the minimal Higgs model and if the mass happens to lie in the interval $2m_t < m_H < 2M_W$, then although 10^6 Higgs would be produced in an SSC year, it is not now known how to detect them in their $H \rightarrow \bar{t}t$ decay mode.[6] For now the only known way of discovering the Higgs boson in this mass range is to build a $\sqrt{s} = 300$ GeV e^+e^- collider with a luminosity of $\sim 10^{32}$ cm.$^{-2}$ s.$^{-1}$. Even in this rather unlikely scenario, the SSC would contribute to our understanding of symmetry breaking by verifying the absence of the excess gauge boson pairs associated with new strong interactions. (Strong coupling models might have light scalars that would approximately mimic a light Higgs boson.) In general the absence of additional gauge boson pairs from WW fusion would be our cue for a redoubled search of the sub-TeV mass scale.

The remainder of the talk is organized as follows: Section 2 reviews relevant aspects of spontaneous symmetry breaking and sketches a proof of the low energy theorems using current algebra methods borrowed from hadron physics. Section 3 is a discussion of unitarity and a conservative strong interaction model that is inspired by the low energy theorems.[7] In Section 4 I will discuss the experimental signals for a strongly interacting symmetry breaking sector that could be seen at the SSC. This includes a careful discussion for the 1 TeV Higgs case and some preliminary estimates for the conservative model for which more careful estimates are in preparation. A brief conclusion is presented in Section 5.

2. SPONTANEOUS SYMMETRY BREAKING AND LOW ENERGY THEOREMS

In order to implement spontaneous symmetry breaking, the lagrangian of the symmetry breaking sector, $\mathcal{L}_{SB}$, must possess a global symmetry group G—analogous to the flavor symmetry of QCD—which breaks by asymmetry of the vacuum to a smaller group H,

$$G \rightarrow H. \tag{2.1}$$

Gauge invariance requires that G include the electroweak $SU(2)_L \times U(1)_Y$ and that H include the unbroken electromagnetic U(1). For each broken generator of G there is a massless goldstone boson in the spectrum of $\mathcal{L}_{SB}$. Three of these couple to the weak currents and are denoted $w^\pm$, z. Others, if any, are denoted by $\{\phi_i\}$. Including the electroweak gauge interactions, the goldstone triplet $w^\pm$, z become longitudinal gauge boson modes $W_L^\pm$, Z_L, and the $\{\phi_i\}$ acquire small masses $0(gM_{SB})$, becoming "pseudo-goldstone" bosons.

As an example, for two flavor QCD with massless quarks the global symmetry G is $SU(2)_L \times SU(2)_R$. After spontaneous symmetry breaking the surviving invariance group is $H = SU(2)_{L+R}$ which is just the isospin group. There are three broken generators, corresponding to the axial generators $SU(2)_{L-R}$, so that three massless goldstone bosons emerge, $\pi^\pm$ and π^0. If there were no other symmetry breaking physics, $\mathcal{L}_{SB}$, $\pi^\pm$ and π^0 would indeed become longitudinal modes of $W^\pm$ and Z, which would however have masses of order 40 MeV rather than $\sim$ 100 GeV.

The statement that the longitudinal modes $W_L^\pm$, Z_L are identified with the goldstone bosons $w^\pm$, z is given a precise meaning by the equivalence theorem, proved to all orders in ref. (7):

$$M\left(W_L(p_1)W_L(p_2)\ldots\right)_U = M\left(w(p_1)w(p_2)\ldots\right)_R + O\left(\frac{M_W}{E_i}\right). \tag{2.2}$$

In eq. (2.2) the left side is an S-matrix element involving longitudinal modes in the U or unitary gauge while the right side is the corresponding goldstone boson amplitude in an R or renormalizeable gauge. As indicated in eq. (2.2), the equivalence holds at energies large compared to the W and Z masses. We can use the equivalence theorem to translate statements about goldstone boson scattering amplitudes into statements about scattering of longitudinally polarized W's and Z's.

As an immediate application, consider the case[7] in which the global symmetry G includes $SU(2)_L \times SU(2)_R$ and H includes an $SU(2)_{L+R}$. For such theories $\rho = 1$ up to electroweak corrections and we may immediately apply the pion low energy theorems derived from current algebra for just this case. For instance, just as for pions we have

$$M(\pi^+\pi^- \rightarrow \pi^0\pi^0) = \frac{s}{F_\pi^2} \qquad s \ll 1 \text{ GeV}^2 \tag{2.3}$$

for $\mathcal{L}_{SB}$ with no other particles that are light compared to M_{SB} we would have

$$M(w^+w^- \rightarrow zz) = \frac{s}{v^2} \qquad s \ll M_{SB}^2 \tag{2.4}$$

where $v = 0.25$ TeV is the familiar vacuum expectation value, $M_W = \frac{1}{2}gv$. With the equivalence theorem this becomes a statement about the scattering of W_L and Z_L in an intermediate energy domain:

$$M(W_L^+W_L^- \rightarrow Z_LZ_L) = \frac{s}{v^2} \qquad M_W^2 \ll s \ll M_{SB}^2. \tag{2.5}$$

Eq. (2.5) and eq. (1.3) are equal up to small corrections, $0(M_W^2/M_{SB}^2)$, for $\rho = 1$.

The assumptions used above, $G \supset SU(2)_L \times SU(2)_R$ and $H \supset SU(2)_{L+R}$, are sufficient to guarantee $\rho = 1$ to all orders in λ_{SB} but they are not known to be necessary conditions. We are therefore motivated to derive the low energy theorems for all candidate groups G and H and for all values of ρ. This has been done by three different methods:[4] a perturbative power counting analysis, the method of effective lagrangians, and current algebra. I will sketch the current algebra derivation below. Along with the low energy theorems for general values of ρ, the derivation establishes a kind of converse to the result quoted above: we find that if $\rho = 1$ then the goldstone boson sector consisting of $w^\pm$, z possesses an effective $SU(2)_{L+R}$ symmetry ("custodial" SU(2)) in the low energy domain $s \ll M_{SB}^2$.

Briefly the derivation is as follows. The global symmetry G must be at least as large as the gauge group, $G \supset SU(2)_L \times U(1)_Y$, so in particular we have the $SU(2)_L$ charge algebra

$$[L_a, L_b] = i\varepsilon_{abc}L_c \tag{2.6}$$

where the corresponding local currents L_a^μ can generally be expanded in terms of the goldstone triplet $w^\pm$, z as

$$L_a^\mu = \frac{1}{2}r_a\varepsilon_{abc}w_b\partial^\mu w_c - \frac{1}{2}f_a\partial^\mu w_a + \ldots \tag{2.7}$$

with terms involving heavy fields omitted. Since $H \supset U(1)_{EM}$ we have $f_1 = f_2$ and $r_1 = r_2$. The f_a are analogues of the PCAC constant and determine the gauge boson masses,

$$M_W = \frac{1}{2} g f_1, \tag{2.8}$$

$$\rho = (f_1/f_3)^2. \tag{2.9}$$

It is straightforward to show that the $SU(2)_L$ algebra requires

$$r_1 = \frac{1}{\sqrt{\rho}}, \tag{2.10}$$

$$r_3 = 2 - \frac{1}{\rho}, \tag{2.11}$$

so that the parameters r_a and f_a in eq. (2.7) are completely determined in terms of G_F and ρ. In particular, $\rho = 1$ implies $f_1 = f_2 = f_3$ and $r_1 = r_2 = r_3 = 1$ which means that the goldstone boson contributions to L_a^μ are the difference of $SU(2)$ vector and axial vector currents. The existence of this vector $SU(2)$ triplet of currents establishes the converse alluded to above.

The rest of the derivation is much like the usual current algebra derivation[3] except that we do not assume an $SU(2)_{L+R}$ isospin invariance. Consequently pole terms which are forbidden by G-parity in the pion case are not forbidden here. Assuming that $w^\pm$, z saturate these pole terms we find goldstone boson low energy theorems such as

$$M(w^+w^- \to zz) = \frac{s}{f_1^2}\frac{1}{\rho} \qquad s \ll M_{SB}^2 \tag{2.12}$$

which using (2.8) reduces to (2.5) for the case $\rho = 1$. By the equivalence theorem we have then

$$M(W_L^+W_L^- \to Z_L Z_L) = \frac{s}{v^2}\frac{1}{\rho} \qquad M_W^2 \ll s \ll M_{SB}^2 \tag{2.13}$$

with $v = f_1 \cong 2M_W/g$. The other two independent amplitudes are

$$M(W_L^+W_L^- \to W_L^+W_L^-) = -\frac{u}{v^2}\left(4 - \frac{3}{\rho}\right), \tag{2.14}$$

$$M(Z_L Z_L \to Z_L Z_L) = 0, \tag{2.15}$$

and by crossing we have also

$$M(W_L^\pm Z_L \to W_L^\pm Z_L) = \frac{t}{v^2}\frac{1}{\rho}, \tag{2.16}$$

$$M(W_L^+W_L^+ \to W_L^+W_L^+) = M(W_L^-W_L^- \to W_L^-W_L^-) = -\frac{s}{v^2}\left(4 - \frac{3}{\rho}\right). \tag{2.17}$$

Like (2.13), eqs. (2.14 – 2.17) are valid in the intermediate domain $M_W^2 \ll E_i^2 \ll M_{SB}^2$.

3. UNITARITY AND A CONSERVATIVE MODEL

The linear growth in s, t, u of the amplitudes (2.13 – 2.17) cannot continue indefinitely or unitarity would be violated. For instance the $W_L W_L \to Z_L Z_L$ amplitude (2.13) is pure s-wave. If we adopt the low energy amplitude (2.13) as a model of the absolute value of the scattering amplitude, then the $J = 0$ partial wave amplitude is

$$\left|a_0(W_L^+ W_L^- \to Z_L Z_L)\right| = \frac{s}{16\pi v^2} \tag{3.1}$$

where here and hereafter we set $\rho = 1$. Unitarity requires $|a_0| \leq 1$ so we see that the growth of a_0 must be cut off at a scale Λ with

$$\Lambda \leq 4\sqrt{\pi}v = 1.8 \text{ TeV}. \tag{3.2}$$

At the cutoff $\sqrt{s} = \Lambda$ the order of magnitude of the amplitude is

$$|a_0(\Lambda)| = \frac{\Lambda^2}{16\pi v^2}. \tag{3.3}$$

For $\Lambda \lesssim v \cong \frac{1}{4}$ TeV we have $|a_0(\Lambda)| \ll 1$ indicating a weakly interacting theory for the symmetry breaking dynamics $\mathcal{L}_{SB}$, while for $\Lambda \gtrsim 1$ TeV we have $|a_0(\Lambda)| \cong 0(1)$, the hallmark of a strong interaction theory. Though there is one counterexample mentioned below, the most likely dynamics is that the cutoff scale Λ is of the order of M_{SB}, the mass scale of the new quanta. Then for $\Lambda \cong 0(M_{SB})$ eq. (3.3) establishes the relationship mentioned in the introduction between the mass scale of the new quanta and the strength of the new interactions: weak coupling for $M_{SB} \ll 1$ TeV and strong coupling for $M_{SB} \gtrsim 0(1)$ TeV.

A weak coupling example is provided by the standard Higgs model[1] with a light Higgs boson, $m_H \ll 1$ TeV, which can be treated perturbatively. Then $a_0(W_L W_L \to Z_L Z_L)$ is given in tree approximation by (where I neglect M_W^2/s)

$$a_0 = \frac{-s}{16\pi v^2} \frac{m_H^2}{s - m_H^2}. \tag{3.4}$$

For $s \ll m_H^2$ this agrees with the low energy theorem (3.2) while for $s \gg m_H^2$ it saturates at the constant value $m_H^2/16\pi v^2$. Comparing with (3.3) we see that m_H indeed provides the scale for Λ.

A strong interaction example is provided by hadron physics. For the $J = I = 0$ partial wave, the low energy theorem[3] gives

$$a_{00}(\pi\pi \to \pi\pi) = \frac{s}{16\pi F_\pi^2} \tag{3.5}$$

with $F_\pi = 92$ MeV. Eq. (3.5) saturates unitarity at $4\sqrt{\pi}F_\pi = 650$ MeV which is indeed the order of the hadron mass scale. The a_{11} and a_{02} amplitudes saturate at 1100 and 1600 MeV.

The two generic possibilities are illustrated in fig. (2). For weak coupling the partial wave amplitudes saturate at values small compared to 1 giving rise to narrow resonances at masses well below 1 TeV. For strong coupling they saturate the unitarity limit with broad resonances in the TeV range.

There is one known model,[8] whose real physical significance is not clear, which has behavior different from figure (2), namely the 0(2N) Higgs model solved to leading order for $N \to \infty$ and then evaluated at $N = 2$ which corresponds to the standard model. (This is only a little worse than the large N_{color} limit for QCD which approximates $3 \cong \infty$; I admit to uneasiness with both approximations.) In that model, which is a strong coupling model, the low energy theorems are of course valid and there is indeed a slow (logarithmic) saturation of partial wave unitarity at the TeV scale, but there are no discernible resonances in the TeV region.

The 0(2N) model or the possibility (which cannot be definitively excluded by the heuristic estimates given above) that resonance structure might be deferred to 2 TeV or above both motivate a conservative model for strong interactions that Mary Gaillard and I have considered.[7] In this model we represent the absolute

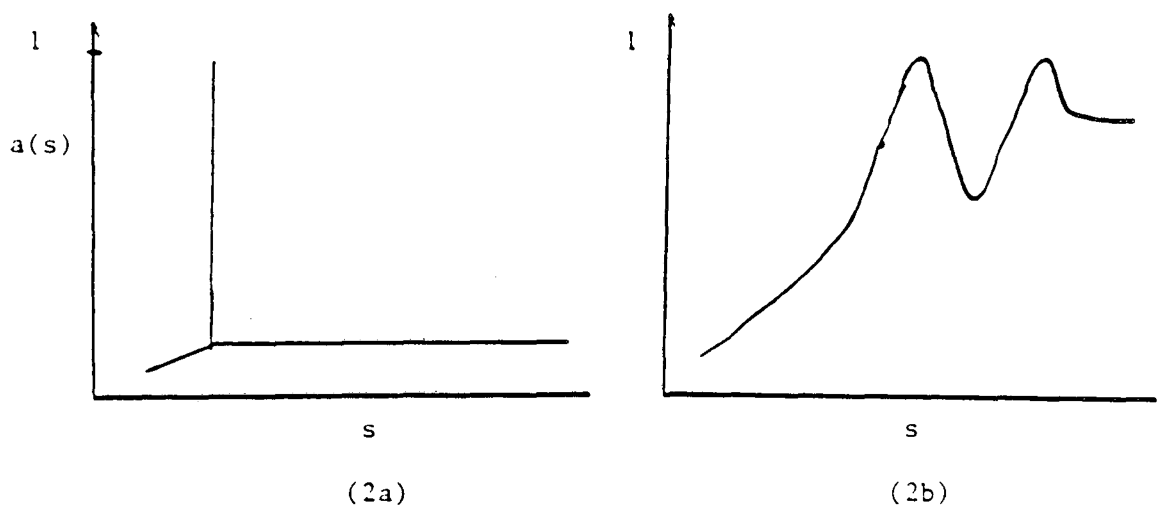

Figure 2. Typical behavior of partial wave amplitudes. Fig. (2a) corresponds to weak coupling — narrow resonance(s) much lighter than 1 TeV and saturation well below the unitarity limit — while fig. (2b) represents strong coupling — broad resonances at the TeV scale and saturation at the order of the unitarity limit.

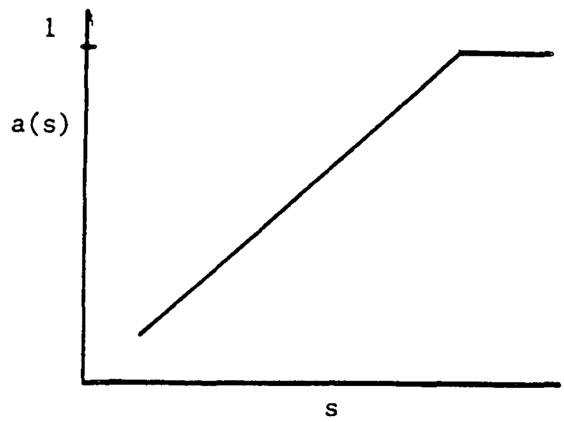

Figure 3. The conservative strong interaction model discussed in the text.

values of the partial wave amplitudes by the low energy theorems up to the energy at which unitarity is saturated and set them equal to one for higher energies, as shown in figure (3). The model is conservative in three respects:

- it neglects possible (or, I would say, likely) resonance structure, underestimating the yield for the 1 TeV Higgs boson by $\sim 50\%$.
- it neglects higher partial waves which surely begin to contribute as the lowest waves saturate.
- it correctly represents the order of magnitude seen in $\pi\pi$ data.

We discuss the experimental implications in the next section.

4. EXPERIMENTAL SIGNALS FOR STRONG INTERACTION MODELS

We consider what might actually be observed at the SSC if $\mathcal{L}_{SB}$ is a strongly interacting theory. The generic strong interaction signal is longitudinally polarized W, Z pairs produced by WW fusion, fig. (1). It will help to measure the polarization of the gauge bosons[9] (if statistics is sufficient) but I will *not* assume polarization information in the results given here. Then the irreducible background is from $\bar{q}q \rightarrow WW$, WZ, ZZ. Since $\mathcal{M}(\bar{q}q \rightarrow WW) = 0(g^2)$ while $\mathcal{M}(qq \rightarrow qqWW) = 0(g^2\lambda_{SB})$, we expect a discernible signal if and only if $\mathcal{L}_{SB}$ is strongly interacting, $\lambda_{SB} = 0(1)$. The signal may occur in W^+W^- and ZZ, as for the standard Higgs boson, but also more generally in $W^{\pm}Z$ and even W^+W^+ and W^-W^- ("I" = 1 and 2 channels).

In order to see the signal over the $\bar{q}q$ background, it is essential to cut on the W, Z rapidity, generally $|y_W| < 1.5$, and the diboson mass, say $M_{WW} > 1$ TeV. Yields from ref. (7) are shown in Table 1. The first number denotes the $\bar{q}q$

	20 TeV	40 TeV
ZZ	150, 90, 250	370, 470, 1100
W^+W^-	660, 120, 500	1600, 630, 2200
$W^\pm Z$	290, 120	670, 670
$W^+W^+ + W^-W^-$	0, 200	0, 1200

Table 1: yields in events per 10^4pb^{-1} for 20 and 40 TeV pp colliders, taken from ref. (7). Cuts are $M_{VV} > 1.0$ TeV and $|y_V| < 1.5$ except for $W^+W^+ + W^-W^-$ for which the rapidity cut is relaxed to $|y_V| < 4.0$. In each entry the first number is the $\bar{q}q$ annihilation background, the second is the conservative strong interaction model, and the third is the 1 TeV Higgs boson.

background, the second is the conservative model, and the third where present is for the 1 TeV Higgs boson. The units are events per $10^4 pb^{-1}$, *i.e.*, 10^7 sec. of running at $\mathcal{L} = 10^{33}$ cm.$^{-2}$ s.$^{-1}$. For like-charged W pairs the rapidity cut is relaxed to $|y_W| < 4$ since there is no $\bar{q}q$ background, the leading background from single gluon exchange being perhaps ~ 5 times smaller than the $\bar{q}q$ backgrounds in the other channels.[10]

Notice the uncharacteristically sharp dependence on the machine energy, 20 TeV versus 40 TeV. This is simply because the signal is at the edge of phase space. Not only the signal but also the signal:background ratio suffer at lower energy. The greater sensitivity of the signal than the background probably reflects the four body phase space of the signal, $qq \to qqWW$, compared to the two body phase space of the background $\bar{q}q \to WW$. To compensate for lower energy, luminosity would need to be increased beyond what would be needed just to equal the signal of a higher energy machine.

Table 1 is chiefly a theoretical exercise since it does not necessarily correspond to experimentally implementable signals. We need to consider how the gauge bosons decay and are detected. This has been done with some care for the 1 TeV Higgs but not yet for the conservative model.

The cleanest channel, $ZZ \to e^+e^-/\mu^+\mu^- + e^+e^-/\mu^+\mu^-$, has a small branching ratio, $B = 3.6 \cdot 10^{-3}$. For the 1 TeV Higgs and the cuts of Table 1 this gives a signal of 4 events over a background of 1 for $\sqrt{s} = 40$ TeV. A more promising leptonic channel is[7,11] $ZZ \to e^+e^-/\mu^+\mu^- + \bar{\nu}\nu$ with $B \sim .02$, analogous to observing $W \to e/\mu + \nu$. The signal is defined by 1) $Z \to e^+e^-/\mu^+\mu^-$ at large p_T with central rapidity, 2) large missing p_T, and 3) no hot jet activity in order to veto the background from $W+$ jet. The latter especially must be studied with a Monte

Carlo but is likely to be both clean and efficient. For the 1 TeV Higgs, Cahn and I found[11] that cuts of $|y_{e^+e^-}| < 1.5$ and $p_T(e^+e^-) > .45$ TeV give a signal of 43 events over a $\bar{q}q$ background of 7 for 40 TeV (16 standard deviations) compared to 7 over 3 for 20 TeV. The $\bar{q}q$ backgrounds are under adequate theoretical control and will in any case be measured at the SSC.

A second channel that has been well studied[12,10] is $H \to WW \to e\bar{\nu}/\mu\bar{\nu} + u\bar{d}/c\bar{s}$. Though the QCD background from Wjj is formidable, Gunion and Soldate have identified cuts which allow the signal to emerge. For the 1 TeV Higgs my extrapolation of their result for $m_H = 0.8$ TeV gives a signal of ~ 300 events over a background of ~ 500 at $\sqrt{s} = 40$ TeV (or $\sim 13\ \sigma$).

For the conservative strong interaction model defined in Section 4, comparable calculations are not yet completed. Except for $ZZ \to e^+e^-/\mu^+\mu^- + e^+e^-/\mu^+\mu^-$ (with a signal of 2 events over a background of 1), I can now offer only some crude guesses. It is however clear that the conservative model will be even more sensitive to machine energy than the 1 TeV Higgs since the signal of the former is smeared over larger values of M_{WW}.

For $ZZ \to e^+e^-/\mu^+\mu^- + \bar{\nu}\nu$ and $WW \to e\bar{\nu}/\mu\bar{\nu} + u\bar{d}/c\bar{s}$ I expect the yields to be roughly half of those quoted above for the 1 TeV Higgs boson. Signals comparable to $ZZ \to e^+e^-/\mu^+\mu^- + \bar{\nu}\nu$ are also expected in $WZ \to \ell\bar{\nu} + e^+e^-/\mu^+\mu^-$ and in $W^+W^+ + W^-W^- \to \mu\mu\nu\nu$. If the electron charge can also be measured the signal for the like-charged W channel is multiplied by a factor 4. The channel $WW \to \ell\bar{\nu} + \bar{\ell}\nu$ may also provide a useable signal without determining the lepton charges.

5. CONCLUSION

With the SSC design parameters, $\sqrt{s} = 40$ TeV and $\mathcal{L} = 10^{33}\mathrm{cm.}^{-2}\mathrm{s.}^{-1}$, we are assured of the capability to see the signal of a strongly interacting symmetry breaking sector. The SSC is strategically positioned in that it is close to being the minimal machine about which this statement can be made. If we do see signs of TeV scale, strong interaction, symmetry breaking physics, then (hard though it may be to imagine now) we are certainly going to think about a next generation facility that will allow more detailed studies.

Because the SSC has the capability to see the strong interaction signal, we will learn something even by not seeing the signal. We then learn that the physics of symmetry breaking involves quanta that are below 1 TeV and not above. If we have not already found them we will know where we need to look. Large numbers of the new quanta would be produced at the SSC, so that detailed studies are likely to be possible in most physics scenarios.

Light quanta—*e.g.*, pseudo goldstone bosons—can occur in TeV-scale strong interaction models. Thus without detailed studies of their properties, the discovery of light Higgs candidates may not immediately and unambiguously tell us whether $\mathcal{L}_{SB}$ involves TeV scale strong interactions. Therefore even if light scalars are discovered, say at SLC or LEP, we will need to look for the strong interaction signal at the SSC in order to clearly understand the nature of the symmetry breaking physics.

Finally it is important to keep in mind, as discussed in the introduction, the physics I have not discussed. It is potentially very rich and the SSC probes it with great facility. I have concentrated just on what must be there—electroweak symmetry breaking—and have argued that at the design energy and luminosity the SSC is sure to illuminate the underlying physics.

References

1. S. Weinberg, Phys. Rev. Lett. 19 (1967) 1264; A. Salam in "Elementary Particle Physics" (ed. N. Svartholm, Almqvist and Wiksells, Stockholm, 1968) p. 367.

2. M. Chanowitz, LBL-21973, to be published in the Proceedings of the XXIII International Conference on High Energy Physics, Berkeley, 1986.

3. S. Weinberg, Phys. Rev. Lett. 17 (1966) 616.

4. M. Chanowitz, H. Georgi and M. Golden, Phys. Rev. Lett. 57 (1986) 2344 and manuscript in preparation.

5. R. Cahn and S. Dawson, Phys. Lett. 136B (1984) 196; E 138B (1984) 464. See also D. Jones and S. Petcov, Phys. Lett. 84B (1979) 440; and Z. Hioki, S. Midorikawa, and H. Nishiura, Prog. Theor. Phys. 69 (1983) 1484. For the effective W approximation see also M. Chanowitz and M. Gaillard, Phys. Lett. 136B (1984) 196; S. Dawson, Nucl. Phys. B29 (1985) 42; G. Kane, W. Repko, and W. Rolnick, Phys. Lett. 148B (1984) 367.

6. B. Cox and F. Gilman, Proc. Snowmass 1984 (ed. R. Donaldson) p. 87.

7. M. Chanowitz and M. Gaillard, Nucl. Phys. B261 (1985) 379.

8. M. Einhorn, Nucl. Phys. B246 (1984) 75; R. Casalbuoni, D. Dominici, and R. Gatto, Phys. Lett. 147B (1984) 419.

9. M. Duncan, G. Kane, and W. Repko, Nucl. Phys. B272 (1986) 517.

10. J. Gunion *et al.*, Theoretical Report of W/Z/Higgs SSC Working Group, UCB-86-39, to be published in Proc. Snowmass 1986.

11. M. Chanowitz and R. Cahn, Phys. Rev. Lett. 56 (1986) 1327.

12. J. Gunion and M. Soldate, Phys. Rev. D34 (1986) 826.