## Title
Probing Electroweak Symmetry Breaking at Multi-TeV Colliders



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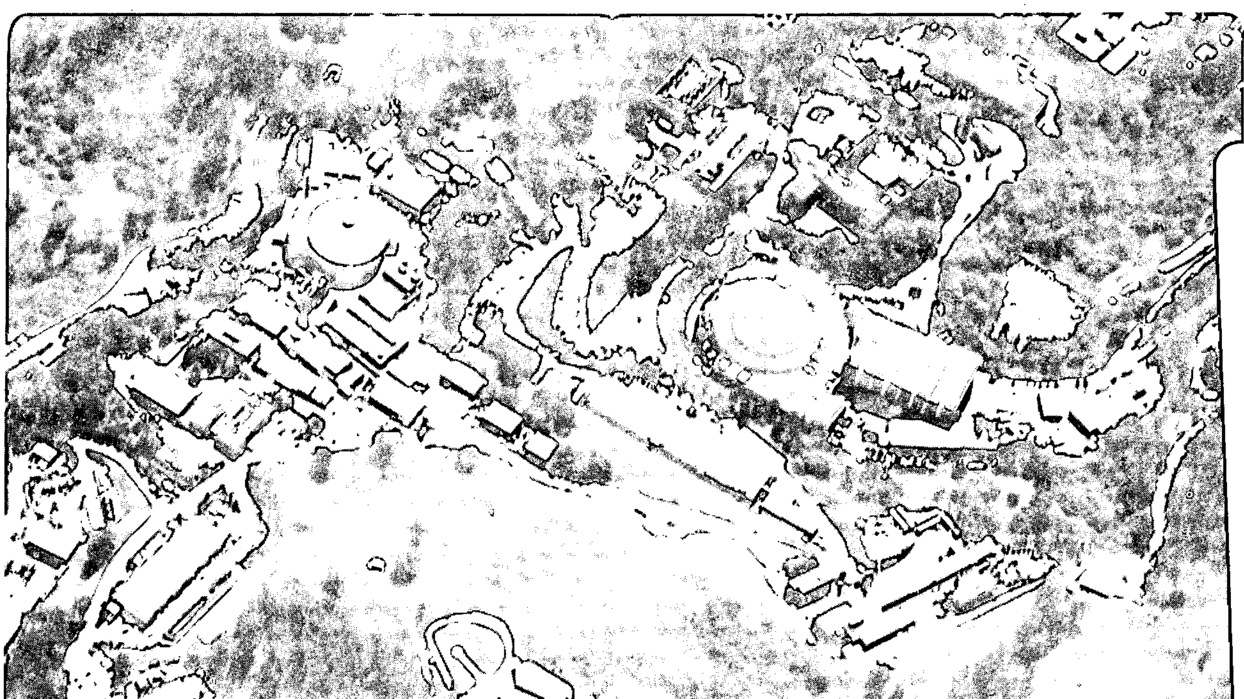



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# Probing Electroweak Symmetry Breaking at Multi–TeV Colliders *

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**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 if electroweak symmetry breaking is due to new physics at the TeV scale. A $pp$ collider with $\mathcal{L}, \sqrt{s} = 40$ TeV, $10^{33} cm.^{-2} s^{-1}$ is just sufficient to observe the signal while $pp$ colliders with 40, $10^{32}$ or 20, $10^{33}$ are not. A collider that is sensitive to the TeV–scale signal provides valuable information about symmetry breaking whether the masses of the associated new particles are below, within, or above the 1–2 TeV region.

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## 1. INTRODUCTION

This is not a general review of TeV collider physics but focuses on a particular issue—electroweak symmetry breaking—which is central to the motivation for constructing multi–TeV scale *pp* colliders. There are other excellent reasons for these colliders: 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.*... This physics is 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, $\sqrt{s}$, $\mathcal{L} = 20$, $10^{33}$.

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 $\lambda_{SB}$ 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 within or near the 1–2 TeV region 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 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

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

The low energy theorem provides the correlation between the mass scale $M_{SB}$ and the interaction strength $\lambda_{SB}$. 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

*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 $O(1)$, then probably $M_{SB} \lesssim$ 0(2) TeV and the low-lying spectrum of $\mathcal{L}_{SB}$ is (just) visible. However, the 1–2 TeV strong coupling signal would be observable even if $M_{SB} >> 2$ TeV and the new particles were too heavy to produce.

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] (See however the possibility discussed by G. Kane in these proceedings.) 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.

I want to make a few comments to put the above statements in perspective. TeV scale symmetry breaking physics puts the greatest demands on energy and luminosity. Most other physics that has been contemplated is less demanding and less sensitive to the difference between 20 and 40 TeV. However, symmetry breaking physics has a special claim on our attention, since while other possible new physics (e.g., SUSY, $Z', \cdots$) may or may not occur in nature, we know for sure that the electroweak symmetry is broken. The no–lose corollary says that a facility able to "see" the 1–2 TeV signal of strongly–coupled symmetry breaking gives us valuable information about the mass scale $M_{SB}$ of the symmetry breaking physics whether the new particles occur below, within, or above the 1–2 TeV region. As will become clear from the results presented below, a *pp* collider with $\sqrt{s} = 40$ TeV and $\mathcal{L} = 10^{33} cm.^{-2} s.^{-1}$ is *just* able to observe the 1–2 TeV strong–interaction signal but is not sufficient to study it in detail. In this sense it is an optimal probe, since we would not want to consider a more ambitious facility without first knowing that there is TeV–scale physics to study.

Though the problem of observing the strong interaction symmetry breaking signal at TeV $e^+e^-$ colliders has not been as extensively studied, several authors have had a first look at the signal and/or backgrounds.[7] The conclusion from Bento and Llewellyn–Smith is that the signal of the conservative model discussed below would not be visible at an $e^+e^-$ collider with $\sqrt{s}, \mathcal{L} = 2$ TeV, $10^{33} cm^{-2} s.^{-1}$ but requires a collider in the range $\sqrt{s}, \mathcal{L} \sim 3 - 5$ TeV, $(1 - 2\frac{1}{2})10^{33} cm.^{2} s.^{-1}$ Incidentally, though we are accustomed to $e^+e^-$ colliders being cleaner than hadron colliders, it is amusing in this instance that the dominant background to WW fusion at an $e^+e^-$ collider, from the usual two photon process, $\gamma\gamma \rightarrow WW$, is relatively unimportant at *pp* colliders (because of the quark and electron charges that contribute to the fourth power and also because the smaller electron mass enhances the factor $[\ln E/m_e]^4$). On the other hand the dominant background at the *pp* colliders, from $\bar{q}q \rightarrow WW$ as discussed below, has no counterpart in $e^+e^-$ collisions.

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.[4] Section 3 is a discussion of unitarity and a conservative strong interaction model that is inspired by the low energy theorems.[8] In Section 4 I will discuss the experimental signals for a strongly interacting symmetry breaking sector that could be seen at the SSC. Work in this area continues and is far from complete. 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 $H$,

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

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 $\pi^0$. If there were no other symmetry breaking physics, $\mathcal{L}_{SB}$, $\pi^{\pm}$ and $\pi^0$ would indeed become longitudinal modes of $W^{\pm}$ and $Z$, which would however have masses of order 40 MeV rather than $\sim$ 100 GeV. (This would have made a rather different world than the one we live in!)

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. (8):

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

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[8] 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[3] derived from current algebra for just this case. For instance, just as for pions we have

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

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

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

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:

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

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 $\lambda_{SB}$ 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 $\rho$. The problem we face is equivalent to that of obtaining the pion–pion scattering low energy theorems in the absence of isospin symmetry. We derive the low energy theorems by three different methods:[4] a perturbative power counting analysis, nonlinear chiral lagrangians, and current algebra. I will sketch the current algebra derivation below. Along with the low energy theorems for general values of $\rho$, 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 $L_a^\mu$ can generally be expanded in terms of the goldstone triplet $w^\pm$, $z$ as

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

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,

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

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

Corrections are suppressed by inverse powers of order $M_{SB}$ or, because of quantum corrections, by inverse powers of $4\pi f_a$.

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 $\rho$. 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 $L_a^\mu$ 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

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

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

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

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

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 $M_W^2 \ll E_i^2 \ll M_{SB}^2, (4\pi v)^2$.

## 3. UNITARITY AND A CONSERVATIVE MODEL

It would be difficult to test the low energy theorems directly at a $pp$ collider because of the predominance of the $\bar{q}q \to WW$ background for $s << 1$ TeV$^2$. The only hope is to use the polarization information since the signal consists predominantly of longitudinally polarized gauge boson pairs while the background is predominantly transversely polarized. This seems unlikely as a first generation experiment, though it has not yet been examined carefully.

In any case, at present a more interesting application is to use the low energy theorems to estimate the generic signal we should expect if $\mathcal{L}_{SB}$ is a strongly coupled sector at the TeV scale or above. The problem we face is like the one that physicists of the 1930's would have faced if they knew nothing of nuclei, baryons or other hadrons, but had discovered the pion, measured the PCAC constant $F_\pi$, and recognized (!?) the pion as an almost Goldstone boson. They would have then been able to derive the pion–pion low energy theorems, such as eq. (2.3), and the problem would be to use this information to reconstruct the scale of hadron physics. Though it would take a skilled writer of science fiction to make this a plausible plot line for the 1930's, it is precisely the situation we are in today if $\mathcal{L}_{SB}$ is strongly coupled: our pions are the longitudinal modes of $W$ and $Z$, our "PCAC" constant is $v = 0.25$ TeV, and we have the low energy theorems eqs. (2.13–2.17).

The central ingredient in our considerations is unitarity. 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 $W_LW_L \to Z_LZ_L$ 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 $\Lambda$ 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 $\Lambda$ 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 $\Lambda$.

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,[9] 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.[8] In this model we represent the absolute 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[10] (if statistics is sufficient) but I will *not* assume polarization information in the results given here. Then the irreducible background is from $\bar{q}q \to WW$, $WZ$, $ZZ$. Since $\mathcal{M}(\bar{q}q \to WW) = 0(g^2)$ while $\mathcal{M}(qq \to qqWW) = 0(g^2\lambda_{SB})$, 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).

To compute the expected yields we must convolute the luminosity distribution functions to find $q$ and $\bar{q}$ beams in the incident protons with the luminosity distribution to find longitudinally polarized gauge bosons in the incident $q$'s or $\bar{q}$'s, convoluted finally with the $2 \to 2$ scattering cross section of the longitudinally polarized gauge bosons:

$$\sigma(pp \to Z_L Z_L + \cdots) = \int_\tau \left.\frac{\partial \mathcal{L}}{\partial \tau}\right|_{qq/pp} \cdot \int_x \sum_{V_L} \left.\frac{\partial \mathcal{L}}{\partial x}\right|_{V_L V_L/qq} \cdot \sigma(V_L V_L \to Z_L Z_L) \quad (4.1)$$

where $\tau = s_{qq}/s_{pp}$ and $x = s_{ZZ}/s_{qq}$. In the effective $W$ approximation the luminosity function for $W_L^+ W_L^-$ pairs is [11]

$$\left.\frac{\partial \mathcal{L}}{\partial x}\right|_{W_L W_L/qq} = \frac{\alpha^2}{16\pi sin^4\theta_W}\frac{1}{x}\left[(1+x)\ln\frac{1}{x} + 2x - 2\right] \quad (4.2)$$

The $x$ dependence of this function, tabulated in table 1, shows the strong dependence on the available phase space that will be familiar to practitioners of two photon physics at $e^+e^-$ colliders. We see that doubling the energy of the incident quarks from $\sqrt{s_{qq}} = 2\sqrt{s_{WW}}$ to $\sqrt{s_{qq}} = 4\sqrt{s_{WW}}$ increases the $W_L W_L$ luminosity by a factor 20! This accounts for the great sensitivity of the signals discussed below to the total collider energy.

Yields computed with eqs. (4.1) and (4.2) are shown in figures (4–6), taken from reference 8. Figures (4) and (5) show the $ZZ$ signals for the conservative model of section 3 and for the 1 TeV Higgs boson respectively, assuming integrated luminosity of $10^4 pb.^{-1}(10^{33}cm.^{-2}s.^{-1}$ for $10^7 s.)$ at $pp$ colliders with $\sqrt{s} = 10, 20, 30, 40$ TeV. A rapidity cut $|y_Z| < 1.5$ has been imposed in the figures to reduce the $\bar{q}q \rightarrow ZZ$ background which is strongly forward while the signal is relatively isotropic. Figure (6) shows the two signals for the 40 TeV collider as increments to the $\bar{q}q \rightarrow ZZ$ background, again with $|y_Z| < 1.5$. From figure (6) we see the necessity to detect gauge boson pairs at large invariant mass, $m_{VV} \gtrsim O(1)$ TeV, where the signal can emerge from the rapidly falling background.

Yields are presented in Table 2 for the rapidity cut $|y_V| < 1.5$ and invariant mass cut $m_{VV} > 1.0$ TeV. Results are shown for 20 and 40 TeV colliders with $10^4 pb.^{-1}$ integrated luminosity. In each box of the table the first number denotes the $\bar{q}q$ background, the second is the conservative model, and the third, where present, is the 1 TeV standard model Higgs boson. The values are taken from reference (8) except that an error in the Higgs boson yield has been corrected (reducing those yields by $\sim 20\%$ relative to reference (8)). For like–charged $WW$ pairs the rapidity cut in relaxed to $|y_W| < 4$ since there is no $\bar{q}q$ background, the leading background from single gluon exchange being perhaps $\sim 5$ times smaller than the $\bar{q}q$ backgrounds in the other channels.[12]

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 reflects the four body phase space of the signal, $qq \rightarrow qqWW$, compared to the two body phase space of the background $\bar{q}q \rightarrow WW$. 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 2 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 more completely for the 1 TeV Higgs than for the conservative model, though in both cases much work remains to be done.

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 with $|y_Z| < 1.5$ and $m_{ZZ} > 0.9$ TeV this gives a signal of 4 events over a background of 1 for $\sqrt{s} = 40$ TeV. A more promising leptonic channel is[8,13] $ZZ \to e^+e^-/\mu^+\mu^- + \bar{\nu}\nu$ with $B \sim .026$, analogous to observing $W \to e/\mu + \nu$ at the SPS collider. 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 find[13] that cuts of $|y_{e^+e^-}| < 1.5$ and $p_T(e^+e^-) > .45$ TeV give a signal of 27 events over a $\bar{q}q$ background of 8 for 40 TeV (10 standard deviations) compared to 6 over 3 for 20 TeV. (These valves correct an error found in ref. 13 by M. Golden.) The $\bar{q}q$ backgrounds are under adequate theoretical control and will in any case be measured at the SSC.

The $ZZ$ yields from the conservative, resonance–free strong interaction model of section 3 are smaller by about a factor 2 than for the 1 TeV Higgs boson. In the channel $ZZ \to e^+e^-/\mu^+\mu^- + e^+e^-/\mu^+\mu^-$ with $|y_Z| < 1.5$ and $m_{ZZ} > 1.0$ TeV we find a signal/background of only 2/1 events for $\sqrt{s} = 40$ TeV and only 0.4/0.7 for $\sqrt{s} = 20$ TeV. For the channel $ZZ \to e^+e^-/\mu^+\mu^- + \bar{\nu}\nu$ with $|y_Z| < 1.5$ and $p_T > 0.45$ TeV for the observed $Z$ we find signal/background of 15/8 events at 40 TeV (for $5\sigma$ statistical significance) compared to 3/3 events at 20 TeV.

A straightforward leptonic channel is $W^\pm Z \to e\nu/\mu\nu + e^+e^-/\mu^+\mu^-$ with branching ratio 0.011. Though there is no WZ signal in the standard Higgs model, this channel may more generally exhibit important strong interaction effects. Defining the signal by $|y_{W,Z}| < 1.5$ and $m_{WZ} > 1.0$ TeV, the yield at $\sqrt{s} = 40$ TeV from the conservative model is $7\frac{1}{2}$ events over a $\bar{q}q \to WZ$ background of 3. If $W \to \tau\nu$ is also included, signal and background are increased by 50%. Dramatic resonance effects might arise in this channel. For instance, in the $N = 4$ technicolor model, the techni–rho meson is predicted at 1.8 TeV, and its charged states will decay predominantly to $W_L^{\pm} Z_L$. It is produced both by $\bar{q}q$ annihilation and $WZ$ fusion. For the 40 TeV collider we find[8] 12 spectacular events over a $\bar{q}q \to WZ$ background of only 1 event in the leptonic channel with only $e$'s and $\mu$'s, increased to 18 over 1.5 if $W \to \tau\nu$ is included.

The like–charged $WW$ signal of the conservative model will also give rise to striking leptonic events of the form $W^+W^+/W^-W^- \to \ell^+\ell^+/\ell^-\ell^- + \nu\nu$, with no background from $\bar{q}q$ annihilation. For $\ell = e$ or $\mu$ the decay branching ratio is $B = .028$ so that corresponding to the cuts $|y_W| < 4$ and $m_{WW} > 1.0$ TeV there would be 33 events. Of course these cuts are not experimentally implementable in this decay channel. An experimentally meaningful set of cuts remains to be defined, but I expect cuts can be found that will lead to a signal of a few tens of events. Notice that for this channel we must be able to determine the sign of the charge for muons and electrons. If we could only measure the muon charge the yields would drop by a factor 4.

The leptonic decay channals discussed above are experimentally clean but suffer from small branching ratios. There are larger yields to be had if we can detect the hadronic decays of $W$ and $Z$, however we then encounter formidable QCD background. For instance, the "mixed" channel $H \to WW \to u\bar{d}/c\bar{s} + e\bar{\nu}/\mu\bar{\nu}$ has a QCD background from $pp \to Wjj$ where the dijet fakes a $W \to \bar{q}q$ decay. Assuming a (perhaps optimistic) 5% resolution on the $jj$ mass, the QCD background is two orders of magnitude larger than the signal.[14] Since the QCD background is dominated by gluon jets, it would be very helpful if we could learn to distinguish quark and gluon jets.

So far two approaches have been taken to the problem of winnowing this mixed decay signal from the the enormous QCD background. One approach[15] is to apply $p_T$ cuts to the observed jets, using the tendency of the longitudinally polarized $W$'s to decay into jets that are transverse to the $W$ line of flight, unlike both the QCD dijets and the transversely polarized $W$'s which both tend to give jets along the $W$ line of flight. Extrapolating the results of ref. (15) for $m_H = 0.8$ TeV to $m_H = 1.0$ TeV suggests a possible signal of $\sim 400$ events over a background of $\sim 400$. This would be a formidable result if it can actually be accomplished in the laboratory, since it corresponds to reducing the background by almost three orders of magnitude while diminishing the signal by less than one order of magnitude.

A second approach to the mixed modes is to borrow a trick from two photon physics at $e^+e^-$ colliders where detecting a final state $e^\pm$ in the forward direction is a powerful way to isolate a clean sample of two photon events. The analogous idea[16] is to tag the forward jets that occur in $WW$ fusion, $qq \to WW + q + q$, with transverse momentum of 50–100 GeV because of the $W$ mass. Of course this approach also has its own QCD backgrounds, from processes with a dijet faking a $W$ and one or two forward jets faking the tagged quark or quarks. Estimating (but not yet calculating) the QCD background, the results of ref. (16) suggest that for the 1 TeV Higgs boson a signal of a few hundred events over a comparable background is probably attainable.

All this work, for both leptonic and mixed decay modes, is still at a preliminary state. Much remains to be done to refine the cuts that define the signals at the parton level and to determine how well they can be implemented by using Monte Carlo methods with detector simulations.

## 5. CONCLUSION

With the SSC design parameters, $\sqrt{s} = 40$ TeV and $\mathcal{L} = 10^{33}$cm.$^{-2}$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.

Absence of the TeV–scale gauge boson pair signal would indicate that $\mathcal{L}_{SB}$ is weakly interacting with quanta well below the TeV scale. In that case the quanta of $\mathcal{L}_{SB}$ are copiously produced and in most physics scenarios it would be possible to study them in some detail.

As shown by the conservative strong interaction model, which assumes no resonances, we may expect the strong interaction signal to be visible even in the (unlikely) event that $M_{SB} >> 1$ TeV so that the strongly interacting quanta of $\mathcal{L}_{SB}$ are too heavy to produce at a 40 TeV collider. In fact, in that case the low energy theorems would enjoy the maximum possible range of validity. But a more likely prospect is for resonances to occur with the onset of strong scattering as happens in hadron physics. In that case we would observe resonances in the gauge boson pair cross section at a mass scale $M_{SB} \lesssim 0(2 \text{ TeV})$.

*Epilogue: A Better World for our Children*

The subprocess cross section for $W_L W_L$ scattering obtained by extrapolating the low energy theorems into the TeV energy range where they saturate unitarity is

$$\hat{\sigma}_{W_L W_L}|_{1-2\text{ TeV}} \sim \frac{1}{2v^2} \sim O(1nb.) \tag{4.3}$$

which at the SSC translates into an observable cross sections of order a few tenths of a picobarn. If $\mathcal{L}_{SB}$ is strongly coupled then at high energies we may expect geometrical total cross sections with the size scale set by the $W$ boson Compton wave length, just as the pion Compton wave length fixes the scale of hadron cross sections. Then we would have

$$\hat{\sigma}_{W_L W_L}\big|_{>>1\ \mathrm{TeV}} \sim \pi R^2 \sim \frac{\pi}{M_W^2} \sim O(100nb), \tag{4.4}$$

two orders of magnitude larger than eq. (4.3). At low energy, eq. (4.3), the cross section is dominated by just the lowest partial waves, whereas at high energy many partial waves contribute, up to $\ell = pR$. If the SSC probe of the TeV scale reveals strongly coupled structure from the symmetry breaking sector, then our children or grandchildren will use the SSC tunnel and high temperature superconducting magnets for a multi-hundred TeV accelerator that will be for the physics of symmetry breaking what the PS and AGS have been for the physics of hadrons.

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I wish to thank Mitchell Golden, Robert Cahn, and Mary Gaillard for helpful discussions. 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.

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## Tables

| $\frac{1}{\sqrt{x}} = \frac{m_{qq}}{m_{WW}}$ | 1 | 2 | 3 | 4 | 5 | 10 |
|---|---|---|---|---|---|---|
| $\frac{16\pi \sin^4\theta_W}{\alpha^2} \left.\frac{\partial\mathcal{L}}{\partial x}\right\vert_{W_L W_L/qq}$ | 0 | 0.9 | 6 | 17 | 36 | 270 |

Table 1: The luminosity distribution function in the effective $W$ approximation for longitudinally polarized $W$ pairs as a function of the ratio of the intial state quark pair center of mass energy to the final state $WW$ pair center of mass energy, $1/\sqrt{x} = m_{qq}/m_{WW}$.

| | 20 TeV | 40 TeV |
|---|---|---|
| $ZZ$ | 150, 90, 200 | 370, 470, 890 |
| $W^+W^-$ | 660, 120, 410 | 1600, 630, 1790 |
| $W^\pm Z$ | 290, 120 | 670, 670 |
| $W^+W^+ + W^-W^-$ | 0, 200 | 0, 1200 |

Table 2: yields in events per $10^4\mathrm{pb}^{-1}$ for 20 and 40 TeV pp colliders, taken from ref. (8) (with correction discussed in the text). 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 $\overline{q}q$ annihilation background, the second is the conservative strong interaction model, and the third is the 1 TeV Higgs boson.

## Figures

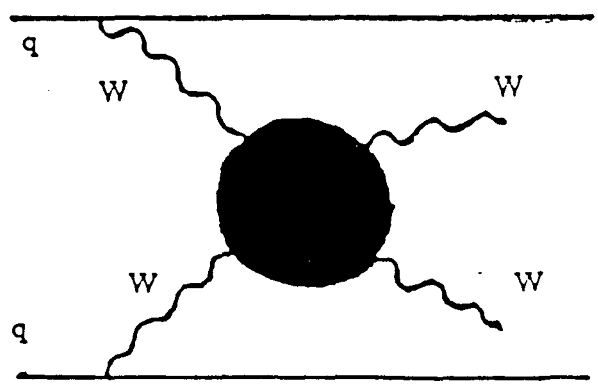
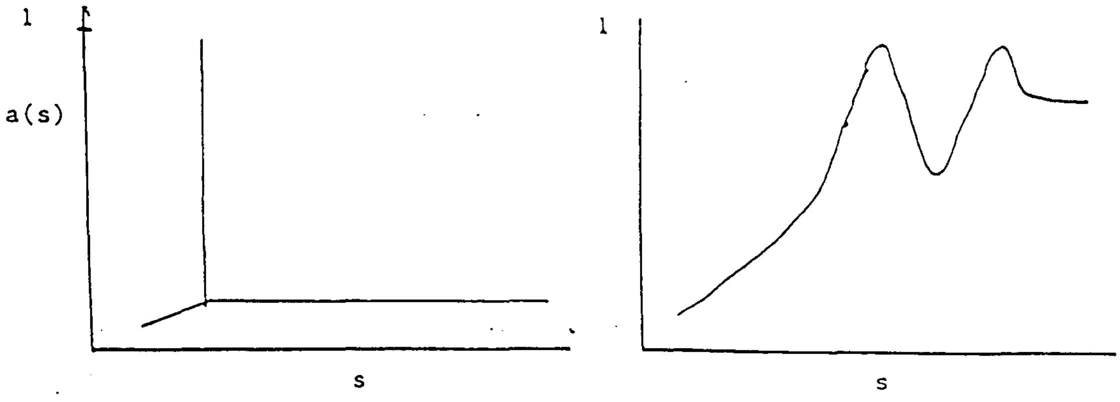



Figure 1: Production of $W$ pairs by $WW$ fusion.




(2a) (2b)

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.

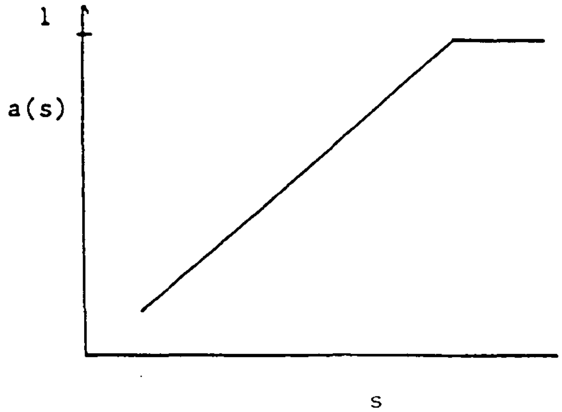



Figure 3. The conservative strong interaction model for the partial wave amplitudes discussed in the text.

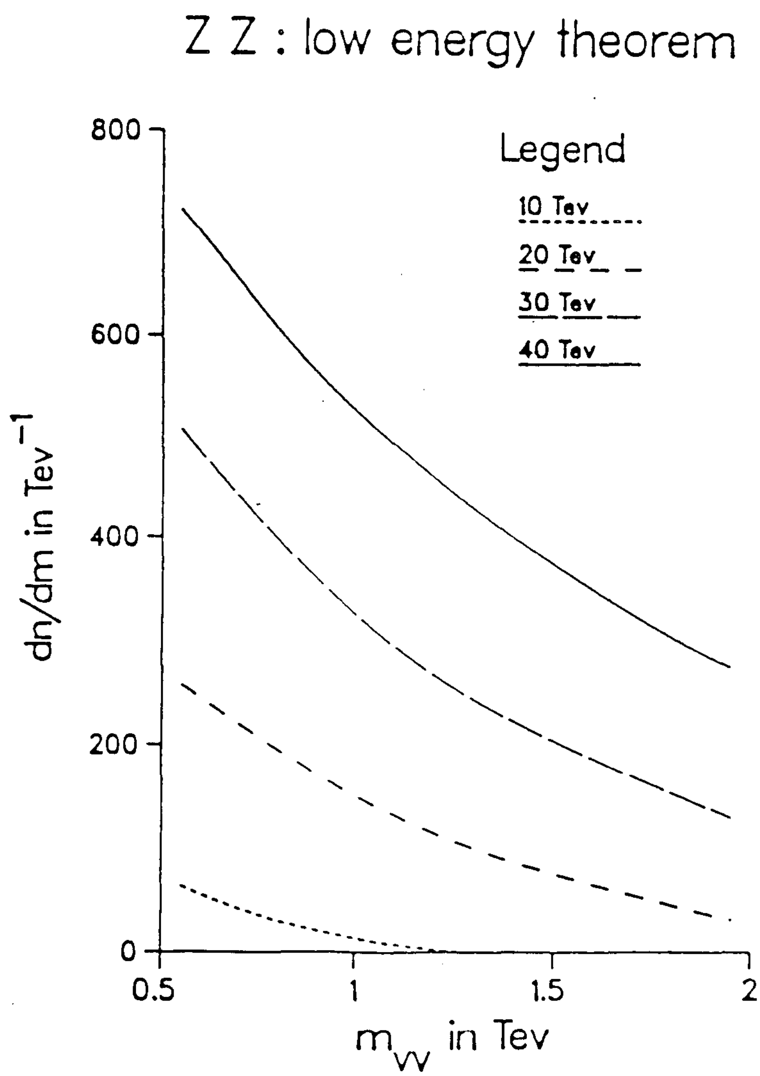



Figure 4. Mass distribution of $Z$ pairs computed from the conservative strong interaction model discussed in the text, with a rapidity cut $|y_Z| < 1.5$. Results are for $10^{-4} pb.^{-1}$ and $pp$ colliders of energy $\sqrt{s} = 10, 20, 30, 40$ TeV.

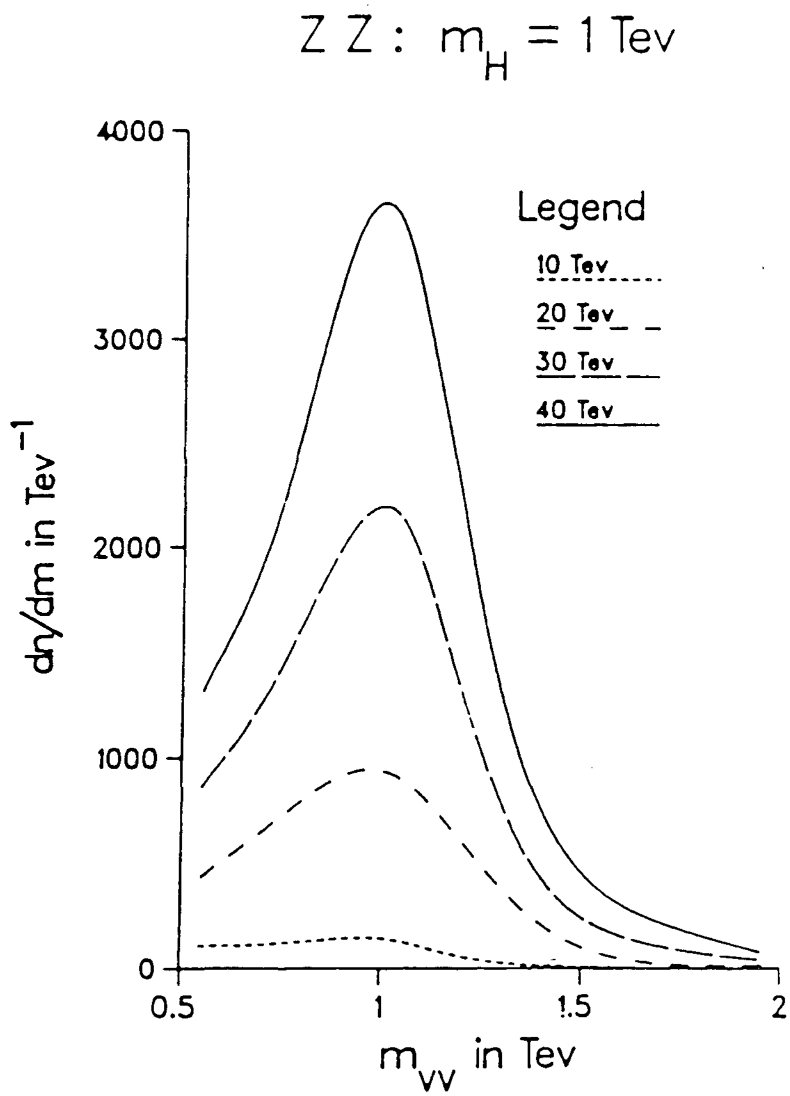



Figure 5. Same as figure 4 for 1 TeV standard model Higgs boson.

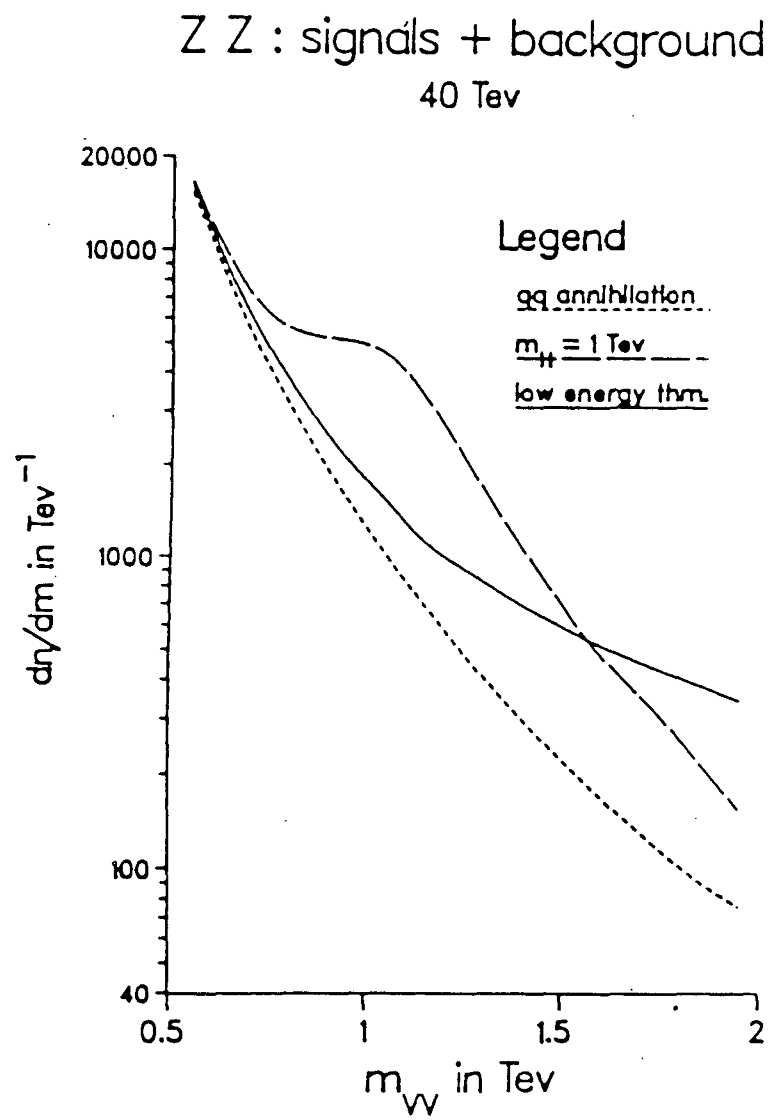



Figure 6. The signals for the consevative strong interaction model and the 1 TeV Higgs boson superimposed inrementally on the background from $\bar{q}q \rightarrow ZZ$, for $10^4 pb^{-1}$ at $\sqrt{s}$ =40 TeV and requiring $|y_Z| < 1.5$.

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