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The No-Higgs Signal: Strong WW Scattering at the LHC [1]

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Abstract

Strong WW scattering at the LHC is discussed as a manifestation of electroweak symmetry breaking in the absence of a light Higgs boson. The general framework of the Higgs mechanism — with or without a Higgs *boson* — is reviewed, and unitarity is shown to fix the scale of strong WW scattering. Strong WW scattering is also shown to be a possible outcome of five-dimensional models, which do not employ the usual Higgs mechanism at the TeV scale. Precision electroweak constraints are briefly discussed. Illustrative LHC signals are reviewed for models with QCD-like dynamics, stressing the complementarity of the $W^{\pm}Z$ and like-charge $W^+W^+ + W^-W^-$ channels.

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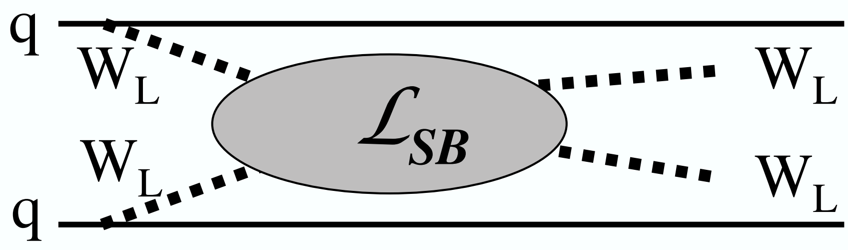

Figure 1: $W_L W_L$ fusion.

1. Introduction

Fifty years of high energy physics have led us to a fundamental question — *What breaks electroweak symmetry?* — that differs from other fundamental questions in one respect: *we know how to find the answer!* The way to the answer is to build and run the LHC. The ability to observe strong WW scattering is an essential part of this prescription. If we do observe it, we learn that elecroweak symmetry breaking is accomplished by strongly coupled quanta above 1 TeV. If we do not observe it (and know we could have it if it were present) then we can conclude that elecroweak symmetry breaking is due to weakly coupled quanta below 1 TeV, which will be Higgs bosons if the Higgs mechanism is valid.

If we do not initially see light Higgs bosons or strong WW scattering, the confirmed absence of strong WW scattering would be a signal to look harder below 1 TeV rather than to expand the search to scales much greater than 1 TeV. This is a key difference with respect to many other searches for new physics, where failure to find the signal at a given energy typically sends us off to search at still higher energies. The ability to observe strong WW scattering confers a "no-lose" capability to determine the mass scale of electroweak symmetry breaking physics.

Even if a light Higgs boson is discovered, it will still be important to measure the WW scattering cross section in the TeV region. If symmetry breaking is due to a light Higgs boson, a central prediction of the Higgs mechanism is that strong WW scattering does *not* occur. As discussed below, strong WW scattering is first-cousin to the famous "bad high energy behavior" of massive vector boson scattering, which it is a principal mission of the Higgs mechanism to remove. If electroweak symmetry breaking is driven by a strong interaction, the cross section for scattering of longitudinally polarized W bosons grows toward the unitarity upper limit, while for symmetry breaking by a weak force it cuts off while it is still small, well below where unitarity would be saturated. In considering the experimental signals at the LHC we should consider both the capability to observe strong WW scattering if it is present and to exclude it if it is not.

The basic idea is that we have already discovered three quanta from the Higgs sector:

the longitudinal spin modes of the $W^{\pm}$ and Z bosons, which in the Higgs mechanism are essentially Higgs sector quanta. By measuring the scattering of the longitudinal modes $W_L W_L \rightarrow W_L W_L$ (where L denotes longitudinal) we are probing Higgs sector interactions,[1] a statement made precise by the 'equivalence theorem.'[2, 3] At the LHC we look for $W_L W_L$ pairs produced by WW fusion, shown in figure 1: the W_L bosons emitted by the colliding quarks are off-shell and must rescatter to emerge on-shell in the final state. Rescattering of electroweak strength is assured by the interactions of the electroweak Lagrangian $\mathcal{L}_{EW}$; these WW pairs are predominantly transversely polarized. If the symmetry breaking sector $\mathcal{L}_{SB}$ is strongly interacting, then the yield is significantly enhanced by the production of longitudinally polarized $W_L W_L$ pairs, indicated in figure 1.

2. The Higgs *mechanism*

The ingredients of the Higgs mechanism are a gauge sector and a symmetry breaking sector,[4]

$$\mathcal{L} = \mathcal{L}_{\text{gauge}} + \mathcal{L}_{\text{SB}} \; . \tag{2.1}$$

$\mathcal{L}_{\text{gauge}}$ has an unbroken *local* $SU(2)_L \times U(1)_Y$ symmetry, with massless, transversely polarized gauge bosons $W^{\pm}$, Z, and γ. $\mathcal{L}_{\text{SB}}$ is the symmetry breaking Lagrangian that describes the dynamics of the symmetry breaking force and the associated quanta. The Higgs mechanism requires $\mathcal{L}_{\text{SB}}$ to have a *global* symmetry group G that breaks spontaneously to a subgroup H, $G \rightarrow H$. We do not know G or H but we do know that $G \supset SU(2)_L \times U(1)_Y$ and $H \supset U(1)_{EM}$, which ensures that the resulting Goldstone bosons include three, $w^{\pm}$ and z, that couple to the three gauge currents corresponding to the three spontanteously broken symmetries of $SU(2)_L \times U(1)_Y$. The Higgs mechanism then ensures that $w^{\pm}$ and z become the longitudinal modes of the gauge bosons $W^{\pm}$and Z which acquire masses, *whether there is a Higgs boson or not.*

The equivalence theorem[2], valid to all orders in gauge and symmetry breaking interactions,[3] codifies the fact that the longitudinal gauge boson modes, $W_L^{\pm}, Z_L$, behave as quanta from $\mathcal{L}_{\text{SB}}$ at high energy, $E \gg m_W$,

$$\mathcal{M}(W_L(p_1), W_L(P_2), \ldots) = \mathcal{M}(w(p_1), w(p_2), \ldots)_R + O\left(\frac{m_W}{E_i}\right), \tag{2.2}$$

where the subscript R denotes a covariant renormalizable gauge choice such as Landau gauge. Equation 2.2 underlies the statement that WW fusion probes $\mathcal{L}_{\text{SB}}$ as indicated in figure 1.

Goldstone boson scattering obeys low energy theorems (LET's), as shown by Weinberg for $\pi\pi$ scattering. The LET's for w, z can be derived without knowing G and H; for instance[5]

$$\mathcal{M}(w^+ w^- \rightarrow zz) = \frac{1}{\rho}\frac{s}{v^2}, \tag{2.3}$$

where v and ρ are the usual vev and rho parameter. Equation (2.3) is valid at low energy, $s \ll M_{\rm SB}^2$, where $M_{\rm SB}$ is the typical mass scale of $\mathcal{L}_{\rm SB}$. Combining the LET (2.3) and the ET (2.2) we obtain the gauge boson LET,

$$\mathcal{M}(W_L^+W_L^- \to Z_L Z_L) = \frac{1}{\rho}\frac{s}{v^2} \tag{2.4}$$

valid in the intermediate energy domain $m_W^2 \ll s \ll M_{\rm SB}^2$.[5]

To understand strong WW scattering it is instructive to consider a U-gauge derivation of the LET which makes no reference to the underlying Goldstone bosons or the ET.[5] In leading order the U-gauge amplitudes involving only gauge sector quanta exhibit the "bad" high energy behavior that would make massive vector boson theories nonrenormalizable. E.g., with just the gauge sector Feynman diagrams we have at high energy

$$\mathcal{M}(W_L^+W_L^- \to Z_L Z_L)_{\text{gauge sector}} = \frac{g^2 s}{4\rho m_W^2} + O(g^2 s^0) \simeq \frac{s}{\rho v^2}, \tag{2.5}$$

which would eventually violate unitarity and render the theory nonrenormalizable. This "bad" behavior is cancelled at the scale $M_{\rm SB}$ by exchange of quanta from $\mathcal{L}_{\rm SB}$, but at low energy, $s \ll M_{\rm SB}^2$, it can be shown that $\mathcal{L}_{\rm SB}$ decouples to all orders. Therefore we again obtain the LET (2.4) for $m_W^2 \ll s \ll M_{\rm SB}^2$. We see that *the LET is precisely the low energy tail of the "bad" UV behavior.*

There are two important conclusions from this discussion:

- In the Higgs mechanism, $W_L W_L$ scattering exhibits the Goldstone boson dynamics of $\mathcal{L}_{\rm SB}$.
- Even if the Higgs mechanism does not occur we see from the U-gauge derivation that the $W_L W_L$ LET is still valid if the physics $\mathcal{L}_{\rm SB}$ that cuts off the amplitude decouples at low energy.

3. Unitarity and the scale of Strong WW scattering

Unitarity implies a rigorous upper bound on the energy at which quanta from $\mathcal{L}_{\rm SB}$ must cut off the growth of $W_L W_L$ scattering. For example, setting $\rho = 1$ (assuming $\mathcal{L}_{\rm SB}$ has a custodial isospin symmetry I), the LET for the $I = J = 0$ partial wave is

$$|a_{00}(W_L W_L)| = \frac{s}{16\pi v^2}. \tag{3.1}$$

Below four-body thresholds unitarity requires $|a_{00}| \leq 1$ and Re $a_{00} \leq 1/2$, which would be violated at 1.8 and 1.2 TeV respectively. These values imply that $M_{\rm SB}$ cannot be too much greater than 1 TeV, say $M_{\rm SB} \leq O(2)$ TeV.

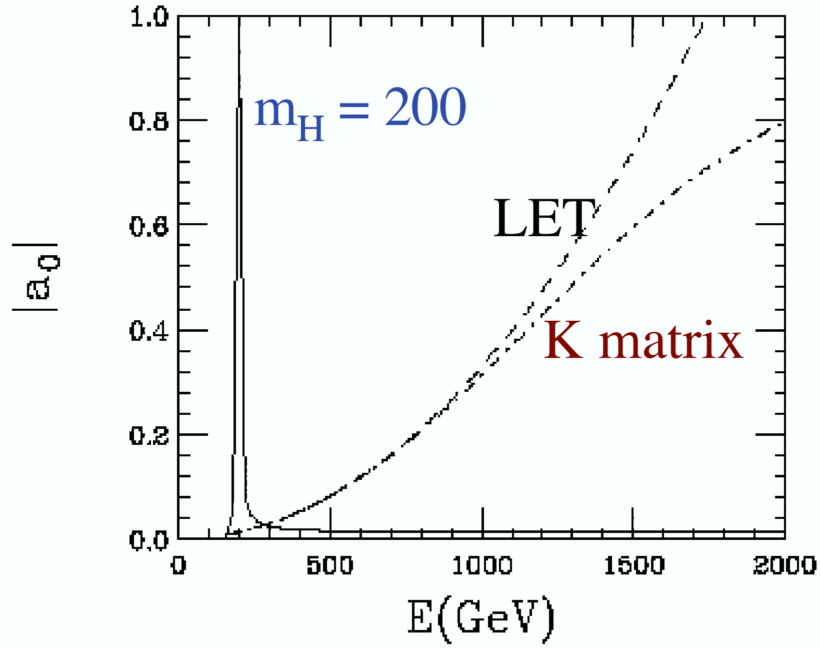

Figure 2: $|a_0(W_L^+W_L^- \to Z_LZ_L)|$ for $m_H = 200$ GeV and two strong scattering models.

The existence of strong WW scattering then depends on the the scale of $M_{\rm SB}$. If $\mathcal{L}_{\rm SB}$ is weakly coupled, then $M_{\rm SB} \ll 1$ TeV and WW scattering never becomes strong. For instance, consider the SM with a light Higgs boson, $M_{\rm SB} = m_H \ll 1$ TeV. Then $\mathcal{M} \simeq s/v^2\ (1 - s/(s - m_H^2))$ where the second term, from H exchange, cancels the first term (2.5) from gauge sector interactions. At high energy the sum is $\simeq m_H^2/v^2 = 2\lambda_H$ and $a_{00} \simeq (m_H/1.8\text{TeV})^2$. On the other hand if $M_{\rm SB}$ is greater than 1 TeV then the partial wave amplitudes rise toward their unitarity limits, $a_{00} \simeq O(1)$, for $E > 1$ TeV, so that the scattering is strong. This is illustrated in figure 2, which shows $|a_0(W_L^+W_L^- \to Z_LZ_L)|$ for two strong scattering models (LET and K-matrix) compared to the SM with $m_H = 200$ GeV.

4. No Higgs mechanism? — EWSB in 5 dimensions

The Higgs mechanism has been an article of faith in high energy physics for about 30 years but it has not been tested experimentally. We will test it at the LHC. The experimental success of the SM implies that $\mathcal{L}_{SU(2)\times U(1)}$ is a good effective theory below the scale of new physics, *even if the Higgs mechanism does not occur in nature.* From the U-gauge derivation we know that the WW low energy theorems are valid in this case, and unitarity would require SOMETHING to cut off $a_0(W_LW_L)$. If $\Lambda_{\rm SOMETHING} \geq O(1)$ TeV then strong WW scattering would occur, just as it would for the Higgs mechanism with $M_{\rm SB} > 1$ TeV.

An example of "SOMETHING" has recently emerged from extra-dimensional theories.

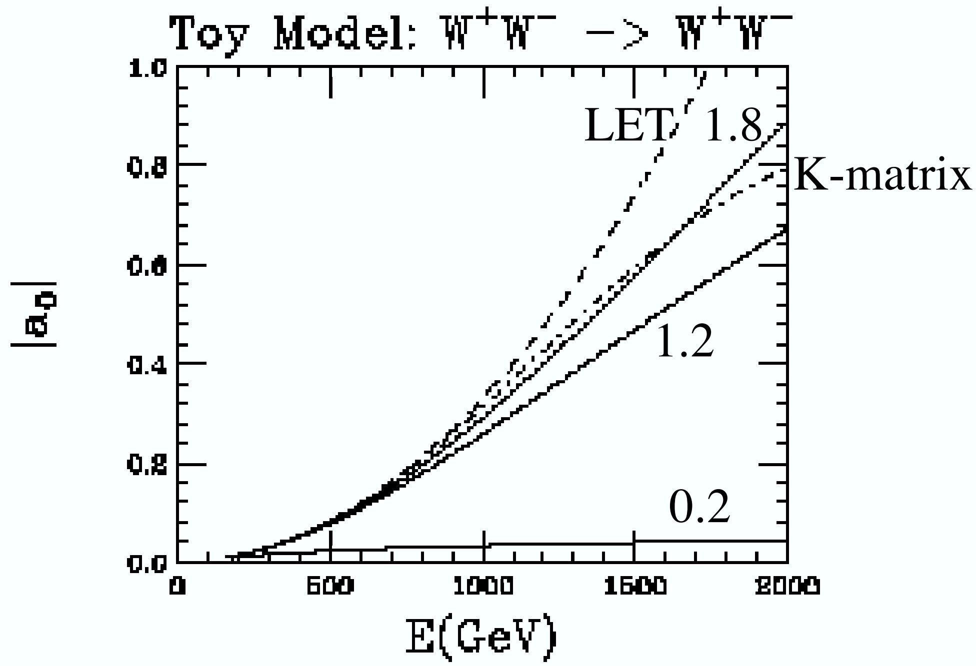

Figure 3: $|a_0(W_L^+W_L^- \to W_L^+W_L^-L)|$ for 5-d model with $m_1 = 0.2, 1.2, 1.8$ TeV.

Models in which $\mathcal{L}_{SU(2)\times U(1)}$ is broken by boundary conditions on a compact fifth dimension can postpone the violation of unitarity to tens of TeV. The cancellation of bad UV behavior at the 1 TeV scale is accomplished by exchange of the Kaluza-Klein excitations of the gauge bosons,[6] W_n, Z_n, γ_n, with masses $M_n \simeq n/R$ where R is the size of the compact dimension.

Yang-Mills theory in five dimensions is nonrenormalizable and can only be an effective theory below a cutoff Λ_5. It is possible for Λ_5 to be an order of magnitude larger than the TeV scale, in which case unitarity in the effective four-dimensional theory is preserved up to Λ_5 by cancellations from exchanges of the Kaluza-Klein gauge bosons.[6] Above Λ_5 the physics underlying the five dimensional theory (strings?) would emerge.

At the LHC W_LW_L scattering could be weak or strong depending on the mass of the KK bosons. For $M_1 \ll 1$ TeV it would be weak while for $M_1 \simeq$ O(TeV) it would be strong. To illustrate this I have considered $W_L^+W_L^- \to W_L^+W_L^-$ scattering in a toy model (actually the Georgi-Glashow model) considered by Csaki *et al.*[7] with symmetry breaking $SU(2) \to U(1)$. The bad high energy behavior of the gauge sector amplitudes is cancelled in this case by the s and t channel exchanges of γ_1, the first Kaluza-Klein excitation of the photon. By varying the radius R of the compactified fifth dimension we can specify the mass M_1 of γ_1. Figure 3 compares the results for $M_1 = 0.2$, 1.2, and 1.8 TeV with the two strong scattering models shown in figure 2. For $M_1 = 0.2$ TeV scattering is weak, as for a light

Higgs boson, but for $M_1 > 1$ TeV there is strong WW scattering as the partial waves grow toward O(1) above 1 TeV.

The strongly coupled versions of these models are preferred experimentally because they are better equiped to hide both from precision electroweak constraints[8] and from direct KK searches. These strongly coupled models resemble technicolor, but the flexibility of the extra-d scenarios give them better prospects to incorporate fermion masses without large flavor changing neutral currents.[9]

Professor Higgs may however have the last laugh: the leading candidate for this class of models, AdS_5, has a dual CFT_4 description, in which symmetry breaking is driven by a strong gauge force, i.e., technicolor in a CFT_4 setting.[10] In this case dynamical symmetry breaking *alla* the Higgs mechanism emerges as the fundamental four dimensional description. And even in the 5-d theory a version of the Higgs mechanism may be at play: the fifth components of the gauge fields act like the Goldstone bosons that become longitudinal gauge boson modes[6] and, extrapolating from a study of brane-induced SUSY breaking, it appears that Wilson integrals looping over the fifth dimension may provide the symmetry breaking condensates.[11]

5. Precision electroweak constraints

The SM fit of the precision electroweak data favors a light Higgs boson with $m_H < 240$ GeV at 95% CL, implying that there would not be strong WW scattering at the LHC. But going beyond the Standard Model, new physics could raise the scale of m_H arbitrarily, even into the realm of dynamical symmetry breaking, as discussed below. Furthermore, the longstanding 3σ discrepancy between A_{LR} and A_{FB}^b, the two most important asymmetry measurements in the SM fit of m_H, raises questions about the reliability of the SM determination of m_H.[12] *LEP cannot definitively determine the scale of electroweak symmetry breaking. LHC can.*

The discrepancy between A_{LR} and A_{FB}^b could be a genuine manifestation of new physics. If it is, the SM fit is moot and we cannot predict m_H from the precision data until the new physics is known. If on the other hand the discrepancy is the result of underestimated systematic uncertainty in the A_{FB}^b measurement (and the two lower precision hadronic asymmetry measurements, A_{FB}^c and Q_{FB} — see [12]), then the SM prediction for the Higgs boson mass is very low, $m_H = 58$ GeV, and is inconsistent at 90% CL with the LEP II lower limit, $m_H > 114$ GeV. In this case new physics would be required to raise the predicted value of m_H into the experimentally allowed region, and again we could not predict m_H from the precision data until the new physics were known. To sustain the SM prediction and upper limit for m_H, the 3σ A_{LR} - A_{FB}^b discrepancy cannot be new physics or a systematic effect but must be a statistical fluctuation. This is surely possible, but so are new physics or underes-

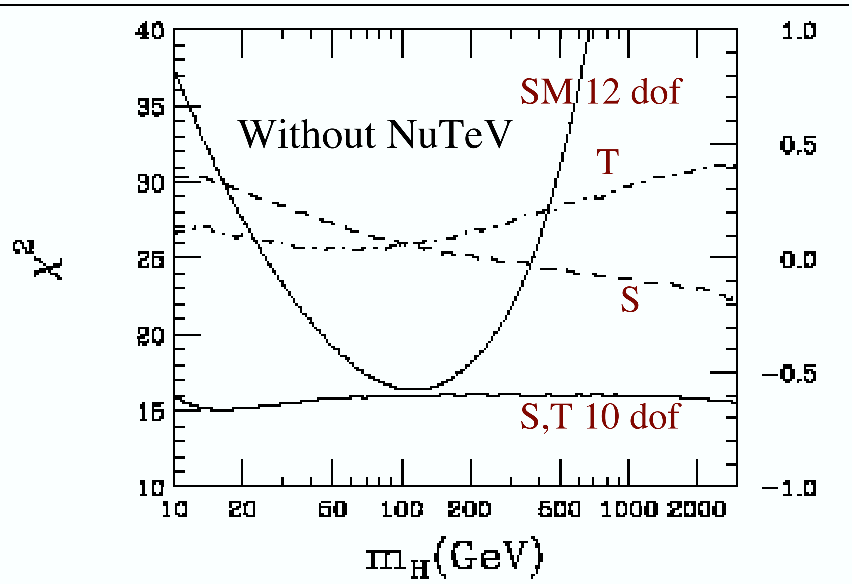

Figure 4: SM and new physics fits to precision data. S, T are read to the right axis.

timated systematic uncertainty. In the latter two cases the precision data does not give us any information about the scale of EWSB and strong WW scattering remains a possibility.

Even if the discrepancy between A_{LR} and A^b_{FB} is a statistical fluctuation so that the SM fit is acceptable at face value, it is possible that the data could also be explained by new physics models with *any* value of m_H. This can be seen explicitly in models in which the new physics is "oblique," i.e., contributes dominantly via corrections to the gauge boson two-point functions. Figure 4, from [12], shows both the SM fit to the precision data and a new physics fit using the oblique parameters S, T[13]. The χ^2 distribution for the S, T fit is almost flat, with no significant preference for any value of m_H between 10 and 3000 GeV. Using the current Summer 2004 data, the confidence level for the oblique fit is 0.17, not very different from the confidence level for the SM fit, which is 0.23 at the χ^2 minimum. (The NuTeV data is not included in these fits; they correspond to "fit C" from [12].)

In figure 4 for each m_H the values of S, T at the χ^2 minimum are shown, plotted against the right vertical axis. To reach the domain of dynamical symmetry breaking and strong WW scattering we require negative S and positive T. Positive T is readily available in new physics models — it corresponds to breaking of custodial isospin — but negative S is harder to come by, although models with $S < 0$ do exist.[14] Recently models of the type discussed above, with electroweak symmetry breaking from boundary conditions on a compact fifth dimension, have been formulated with negative S of the required magnitude.[15] However in a broad class of such models Chivukula *et al.*[16] have obtained the interesting inequality

$$S - 4\cos^2\theta_W \; T = \frac{4}{\alpha}\sin^2\theta_W \cos^2\theta_W \; m_Z^2 \Sigma_n \frac{1}{M_n^2} \geq \frac{1}{3} \tag{5.1}$$

where the sum is over the Kaluza-Klein excitations of the Z boson with mass M_n. The lower limit follows because M_1 cannot be arbitrarily large in order to unitarize WW scattering, as described in the preceding section. It then appears that negative S implies negative T in these models, contrary to what is needed in figure 4. However the values of S, T in these calculations are not equivalent to the experimental S, T, since the latter are normalized with respect to the SM with a reference value for m_H while there is no Higgs boson at all in these 5-d models.[10] It is an open problem to extract oblique parameters from the 5-d models that can be compared directly with experiment.

6. Signals at LHC: complementarity in a QCD'ish example

Strong WW scattering is among the most challenging physics goals of the LHC, requiring the full energy and luminosity. Theorist-level simulations[17, 18, 19] indicate that signals will be observable at the $\geq 5\sigma$ level with $\geq 150 fb^{-1}$ data samples. Studies with realistic detector simulations have been done by ATLAS[20] and CMS[21], but it is fair to say that simulation studies are still in early days.

As a generic example I will briefly review the results of a study[17, 18] of strong WW scattering signals in $SU(N_{TC})$ technicolor, using a scaled version of a chiral effective Lagrangian model of π and ρ interactions[22], which is also the basis of the BESS model[23] of strong EWSB. The partial wave amplitudes obtained from the effective Lagrangian are unitarized by the K-matrix method, which is equivalent to the conventional Breit-Wigner resonance parameterization with the constant imaginary part of the denominator, $m\Gamma$, replaced by $\sqrt{s}\Gamma(\sqrt{s})$.

In [17] the model was shown to give a surprisingly good description of the $I = J = 1$ and $I = 2$, $J = 0$ $\pi\pi$ partial waves. The amplitudes are determined with no free parameters from the known values of F_π, m_ρ, and Γ_ρ. The results are shown in figure 5. While the quality of the fit at 1.2 GeV is probably fortuitous, we can take seriously the qualitative agreement of the model with the data. The ρ meson exchange obviously enhances the signal in the a_{11} partial wave. Less obviously ρ exchange *suppresses* the a_{20} partial wave, causing the slope to begin to flatten out at $\simeq$ 800 MeV, in accord with the data shown in figure 5b. It is then useful to use the model to explore the consequences of varying the ρ mass and width. In particular we find that this enhancement/suppression from ρ exchange implies that LHC signals from the a_{11} and a_{20} partial waves are *complementary*: as m_ρ is decreased the signal at the LHC from a_{11} is enhanced while the signal from a_{20} is suppressed. Conversely for large $m\rho$, a_{20} is enhanced while a_{11} is suppressed.

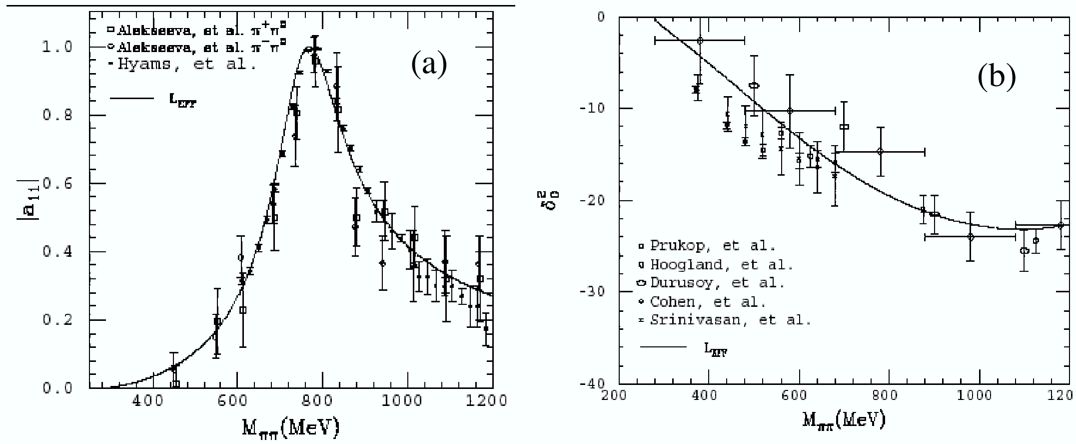

Figure 5: $\pi\pi$ partial waves from chiral Lagrangian.[17].

The ρ mass and width for $SU(N_{TC})$ are obtained from the hadronic m_ρ, Γ_ρ by the scale factor $v/F_\pi \simeq 2700$ modified by factors of $N_{TC}/3$ using the large N scaling laws for $SU(N)$ gauge theories. For $N_{TC} = 2$ and 4 we have $m_\rho = 1.78$ TeV and 2.52 TeV respectively. To represent the possibility that the resonances of $\mathcal{L}_{\mathrm{SB}}$ may be heavier than the naively anticipated 1 - 3 TeV region a 4 TeV ρ meson is also considered, with its width determined from the hadronic $f_{\rho\pi\pi}$ coupling.

The complementarity of the a_{11} and a_{20} channels is evident in figure 6. As m_ρ increases, the amplitudes approach the nonresonant K-matrix model amplitude for $|a_{11}|$ from above and $|a_{20}|$ from below, since chiral invariant ρ exchange enhances the former and suppresses the latter. At the LHC the 4 Tev ρ signal is indistinguishable from the signal of the nonresonant K-matrix model. The fact that the ρ resonance amplitude approaches the nonresonant K-matrix amplitude for large ρ mass is a very general feature, independent of the specific properties of vector meson exchange. It explains the sense in which smooth unitarization models, such as the linear and K-matrix models, are conservative: they represent the "fail-safe" nonresonant scattering signals that are anticipated if the resonances are unexpectedly heavy. This is the most general meaning of complementarity. A more specific meaning, special to vector meson exchange as constrained by chiral symmetry, is the inverse relationship of the $I = 1$ and $I = 2$ channels discussed here.

The experimental signals and backgrounds for the model were computed in [17, 18]. The final states are $W^\pm Z$, which includes the direct channel $\rho^\pm$ resonance, and $W^+W^+ + W^-W^-$, which is pure $I = 2$. The irreducible backgrounds are the SM $\mathrm{O}(\alpha_W^2)$ and $\mathrm{O}(\alpha_W\alpha_S)$ WW fusion amplitudes (the former computed with a 100 GeV Higgs boson). Since jet tagging was not assumed, $\overline{q}q \to WZ$ is also included as a background to the WZ fusion signal.

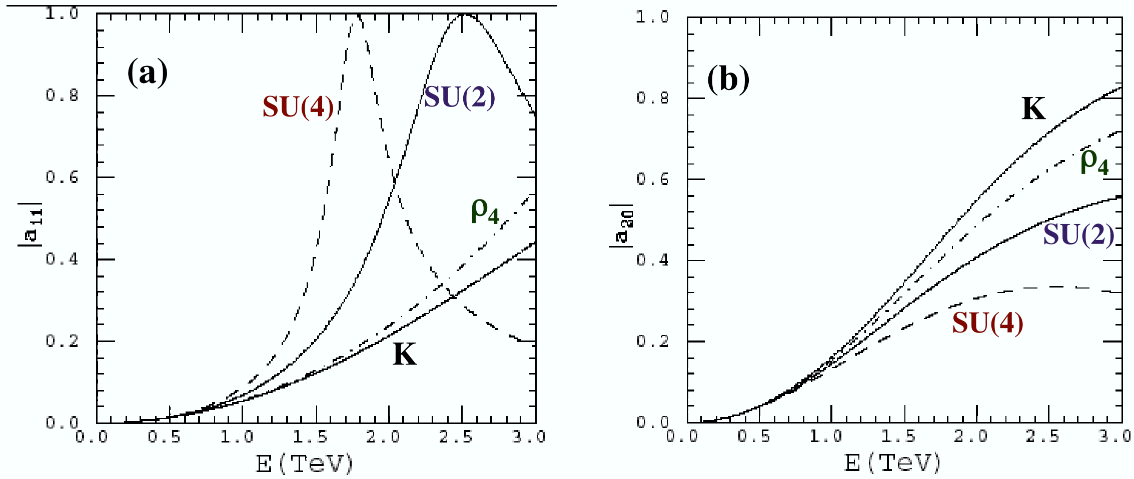

Figure 6: Partial wave amplitudes from technicolor and K-matrix models.[17]

Surprisingly, $\overline{q}q \to W^+Z$ is also a big background for W^+W^+ fusion, due to events in which the negative lepton from the Z escapes detection.[20] Similarly, it turns out that $\overline{q}q \to W^+\gamma^*$ is an equally important background.[18] Top quark backgrounds, $\overline{t}t$ and $\overline{t}tW$, have also been considered and are easily controlled.

While a forward jet tag may eventually prove to be more effective, the results quoted below from [17, 18] rely only on hard lepton cuts and a central jet veto (CJV). The CJV vetos events containing a jet with central rapidity, $\eta_J < 2.5$, and high transverse momentum, $P_T(J) > 60$ GeV; it reduces backgrounds from transversely polarized W bosons, which are emitted at larger transverse momenta than the longitudinally polarized W bosons of the signal. The hard lepton cuts rely on the general property that the strong scattering cross sections are proportional to s_{WW} while the backgrounds scale like $1/s_{WW}$, and on the differing polarization of the signal and background WW pairs. If this strategy suffices it has the advantage of being cleaner than relying on forward jet tagging, which is subject to QCD corrections and to detector-specific jet algorithms and acceptances in the forward region. The results quoted below incorporate reasonable estimates of the experimental efficiencies — see [17, 18] for details. The leptonic cuts are optimized separately for each set of model parameters.

A robust observability criterion is defined and the cuts are optimized by searching over the cut parameter space for the set of cuts that satisfy the observability criterion with the smallest integrated luminosity. The criterion is

$$\sigma^{\uparrow} = S/\sqrt{B} \geq 5 \tag{6.1}$$

$$\sigma^{\downarrow} = S/\sqrt{S+B} \geq 3 \tag{6.2}$$

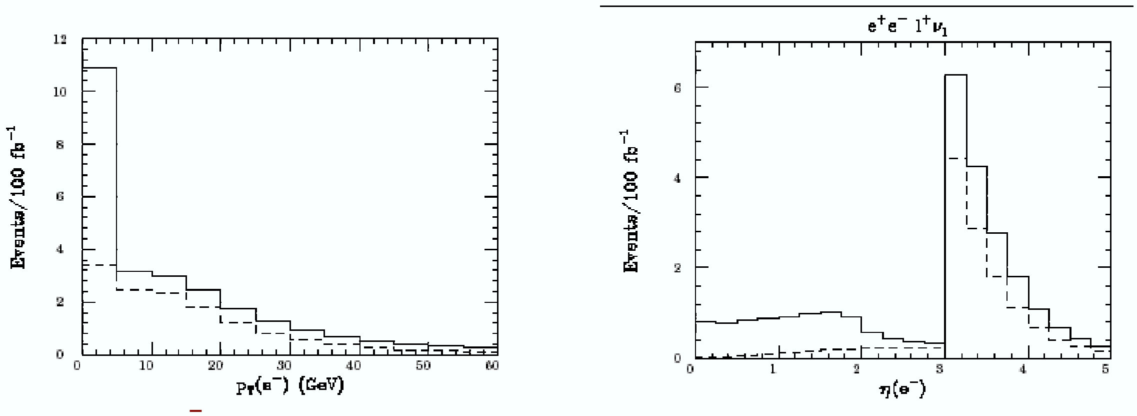

Figure 7: p_T and η distributions for wrong-sign leptons that escape veto.[18]

$$S \geq B, \tag{6.3}$$

where S and B are the number of signal and background events, and $\sigma^{\uparrow}$ and $\sigma^{\downarrow}$ are respectively the number of standard deviations for the background to fluctuate up to give a false signal or for the signal plus background to fluctuate down to the level of the background alone. The requirement $S \geq B$ is imposed so that the signal is unambiguous despite the systematic uncertainty in the size of the backgrounds; this condition might eventually be replaced by a less conservative one, since background studies *in situ* with real LHC data should substantially reduce the systematic uncertainties.

To obtain the W^+W^+ background the complete amplitude for $\overline{q}q \to l^+\nu l^+ l^-$ was computed in [18], including amplitudes which are neither $\overline{q}q \to WZ$ or $\overline{q}q \to W\gamma^*$. They are an inevitable background because any detector has unavoidable blind spots at low transverse momentum and at high rapidity where the l^- escapes detection. At very low p_T, muons will not penetrate the muon detector, electrons or muons may be lost in minimum bias pile-up, and for low enough p_T in a solenoidal detector they will curl up unobservably within the beam pipe. Muon and electron coverage is also not likely to extend to the extreme forward, high rapidity region.

In reference [18] an attempt was made to employ reasonable though aggressive assumptions about the observability of the extra electron or muon. Rapidity coverage for electrons

Table 6.1 Minimum integrated luminosity $\mathcal{L}_{MIN}$ to satisfy significance criterion for $W^+W^+ + W^-W^-$ and $W^{\pm}Z$ scattering.

m_ρ(TeV)	1.8	2.5	4.0
$\mathcal{L}_{MIN}(WW)$ (fb^{-1})	200	150	110
$\mathcal{L}_{MIN}(WZ)$ (fb^{-1})	44	320	NS

and muons was assumed for $\eta(l) < 3$. Within this rapidity range it was assumed that isolated e^- and μ^- leptons with $p_T(l) > 5$ GeV can be identified in events containing two isolated, central, high p_T e^+'s and/or μ^+'s. It was also assumed that electrons (but not muons) with $1 < p_T(l) < 5$ GeV can be identified if they are sufficiently collinear ($m(e^+e^-) < 1$ GeV) with a hard positron in the central region. For $p_T(e^-) < 1$ GeV electrons were considered to be unobservable. These assumptions should be reconsidered in view of the now finalized designs of ATLAS and CMS. Figure 7 shows the p_T and η distributions of the wrong-sign lepton for events in which it escapes the veto and the two like-sign leptons are in the signal region.

The results are summarized in table 6.1, where it is assumed that a high p_T electron or muon can be detected with 85% efficiency, a high p_T Z decaying to electrons or muons with 95% efficiency, and that the $\overline{q}q \rightarrow WZ/W\gamma^*$ background to the like-sign WW signal is rejected with 95% efficiency when the wrong-sign

lepton falls within the acceptance region defined above. (In [18] veto efficiencies of 90 and 98% were also considered.) The 1.8 TeV ρ meson of $N_{TC} = 4$ technicolor provides a signal in the WZ channel that meets the observability criterion with 44fb^{-1}. For the 4 TeV ρ model, the resonance enhancement is unobservable and even for arbitrarily large luminosity there is no set of cuts that meet the observability criterion in the WZ channel. But for this case the like sign WW signal meets the criterion with 110 fb^{-1}. The worst case is intermediate between these two extremes: the $N_{TC} = 2$ model with $m_\rho = 2.5$ TeV is best detected in the like sign WW channel for which 150 fb^{-1} is required. In this sense we may say that 150 fb^{-1} is the "No-lose" luminosity needed to assure the discovery of electroweak symmetry breaking in the worst case.

For the $m_\rho = 2.5$ TeV worst case, the dominant background to the $W^+W^+ + W^-W^-$ signal is from the $\overline{q}q \rightarrow WZ/W\gamma^*$ process: it accounts for 70% of the background while the $O(\alpha_W^2)$ and $O(\alpha_W \alpha_S)$ amplitudes contribute 26 and 3% respectively. Forward jet tagging would eliminate the $\overline{q}q \rightarrow WZ/W\gamma^*$ background, so a lower "No-lose" luminosity might well be attainable. The WZ signal might also be improved if it turns out that the mixed modes are practicable, in which the W decays hadronically while the Z decays leptonically.

Theorist-level estimates are of course oversimplified and optimistic. For instance, apart from guesses at the efficiencies there is no attempt to simulate the detector, to model the effect of pileup or charge misidentification.... Nevertheless it is likely after the LHC begins to run that experimenters will devise aggressive methods that will ultimately yield even better results: they will be free to try bold strategies because they will be able to test them with real data. An example from the past is provided by the experience at LEP I, where experimenters rapidly exceeded the reach for the Higgs boson that was projected in the first CERN yellow book study for LEP, because they were able to show with real data that $Z^* \rightarrow H\overline{\nu}\nu$ is a viable detection mode, which must have seemed too adventurous to the yellow book authors.

7. The bottom line

The origin of electroweak symmetry breaking is completely unknown — many possibilities are open, among them undoubtedly some not yet imagined. With enough luminosity and experimental ingenuity, the LHC at design luminosity and energy is sure to lead us to the answer. To cover all possibilities it is essential to develop the capability to observe strong WW scattering if it exists or to exclude it if it does not.

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