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Can the electroweak symmetry-breaking sector be hidden?

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In a recent paper, Chivukula and Golden claimed that the electroweak symmetry-breaking sector could be hidden if there were many inelastic channels in the longitudinal WW scattering process. They presented a model in which the W's couple to pseudo-Goldstone bosons, which may be difficult to detect experimentally. Because of these inelastic channels, the WW interactions do not become strong in the TeV region. We demonstrate that, despite the reduced WW elastic amplitudes in this model, the total event rate (~ 5000 extra longitudinal W^+W^- pairs produced in one standard SSC year) does not decrease with an increasing number of inelastic channels, and is roughly the same as in a model with a broad high-energy resonance and no inelastic channels.

The "no-lose theorem" states that if light Higgs bosons do not exist, elastic longitudinal W scattering becomes strong at or above 1 TeV, and that the new strong interactions can be detected by observing WW scattering via leptonic decays of W's [1]. (We use W to denote either the $W^\pm$ or Z^0 boson.) In a recent paper [2], Chivukula and Golden have argued that the "no-lose theorem" breaks down if there are many inelastic channels into which the W's can scatter. They presented a toy model in which the W's couple to a large number of pseudo-Goldstone bosons, which may be difficult to detect experimentally. Because of the large number of inelastic channels, there are light resonances in the elastic WW scattering amplitudes, which are too broad to be discernible as peaks. Moreover, the growth of the elastic scattering amplitudes is cut off at the scale of the light resonances, so the WW interactions do not become strong in the TeV region. They conclude that, unless the pseudo-Goldstone bosons themselves can be observed, the electroweak-symmetry breaking sector will remain hidden.

In this note, we point out that the total event rate for elastic WW scattering in the model of ref. [2] does not decrease as the number of inelastic channels increases. The elastic amplitude is smaller, but the resonance is at a lower energy, where the parton luminosity is greater. We give a simple scaling argument to show that the total rate (~ 5000 W^+W^- pairs and ~ 2500 Z^0Z^0 pairs produced in one standard SSC year) is roughly independent of the number of inelastic channels, and is about the same as that in the standard O(4) model [3,4] with no inelastic channels and a broad TeV scale resonance. We also briefly comment on possible methods to detect the signal.

The model presented in ref. [2] to demonstrate the possible effects of inelastic channels in the electroweak sector contains both exact Goldstone bosons and a large number of pseudo-Goldstone bosons. This model possesses an approximate $O(j+n)$ symmetry which is explicitly broken to $O(j)\times O(n)$. The exact $O(j)$ symmetry is spontaneously broken to $O(j-1)$, yielding $j-1$ massless Goldstone bosons, ϕ, and one massive scalar boson. The $O(n)$ symmetry remains unbroken, and there are n degenerate pseudo-Goldstone bosons, ψ, with mass m_ψ.

To use this model to describe the scattering of longitudinal W's, one applies the equivalence theorem, replacing the longitudinal W with its corresponding Goldstone boson ϕ in the S-matrix. This equivalence holds only when $E_W \ll M_W$, where E_W is the energy of the W boson in the WW center-of-mass frame. Therefore, strictly speaking, one should not use this model to describe WW scattering when the invariant mass M_{WW} is less than a couple of times the mass threshold $2M_W$. The amplitudes for the scattering of longitudinal W's are given by [5] [#1]

[#1] For footnote see next page.

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$$\mathcal{M}(Z^0Z^0 \to W^-W^+) = A(s,t,u)\,,$$
$$\mathcal{M}(W^-W^+ \to Z^0Z^0) = A(s,t,u)\,,$$
$$\mathcal{M}(W^-W^+ \to W^-W^+) = A(s,t,u) + A(t,s,u)\,,$$
$$\mathcal{M}(Z^0Z^0 \to Z^0Z^0) = A(s,t,u) + A(t,s,u) + A(u,t,s)\,,$$
$$\mathcal{M}(W^\pm Z^0 \to W^\pm Z^0) = A(t,s,u)\,,$$
$$\mathcal{M}(W^\pm W^\pm \to W^\pm W^\pm) = A(t,s,u) + A(u,t,s)\,. \quad (1)$$

In the models we consider in this note, $A(s, t, u) = A(s)$ depends only on s.

Before turning to the model of ref. [2], we recall some relevant features of the $O(N)$ model [3,4]. To leading order in $1/N$, with the parameter v and Λ_c held fixed as $N\to\infty$, the amplitude $A(s)$ in the $O(N)$ model is given by

$$A(s) = \frac{s}{N}\left\{v^2 - \frac{s}{32\pi^2}\left[\ln\left(\frac{e\Lambda_c^2}{|s|}\right) + i\pi\Theta(s)\right]\right\}^{-1}. \quad (2)$$

Here $\Theta(s>0)=1$ and $\Theta(s<0)=0$, and Λ_c is the cut-off scale of the theory [#2], related to the tachyon mass μ_t through

$$\Lambda_c = \frac{\mu_t}{\sqrt{e}}\exp\left(\frac{-16\pi^2 v^2}{\mu_t^2}\right). \quad (3)$$

(For $\mu_t^2 \gg v^2$, the tachyon and cut-off scales are roughly the same, $\Lambda_c \simeq \mu_t/\sqrt{e}$.) Because of the presence of the tachyon, the $O(N)$ model must be regarded as an effective theory, valid only at energy scales well below μ_t. With v, Λ_c, and s held fixed, the amplitude (2) evidently scales as $1/N$.

To extract physical predictions from the $O(N)$ model, one must set N equal to some finite value; $N=4$ corresponds to the electroweak sector with its three Goldstone bosons and one massive Higgs boson. To ensure that low-energy theorems for the scattering amplitudes are satisfied, one must then set $v=f/\sqrt{N}$, where $f=250$ GeV characterizes the symmetry-breaking scale. The amplitude is then given by

$$A(s) = s\left\{f^2 - \frac{sN}{32\pi^2}\left[\ln\left(\frac{e\Lambda_c^2}{|s|}\right) + i\pi\Theta(s)\right]\right\}^{-1}. \quad (4)$$

With f held fixed, the scaling property of the amplitude differs slightly from that described above; $A(s)$ scales as $1/N$, but only if we simultaneously scale s with $1/N$ and Λ_c with $1/\sqrt{N}$. In other words,

$$A(s) = \frac{1}{N}F(\tilde{s}, \tilde{\Lambda}_c)\,, \quad \tilde{s} = Ns\,, \quad \tilde{\Lambda}_c = \sqrt{N}\Lambda_c\,, \quad (5)$$

where $F(\tilde{s}, \tilde{\Lambda}_c)$ only depends on N through $\tilde{s}$ and $\tilde{\Lambda}_c$.

To locate resonances in the scattering amplitudes (1), we look for the position of the (complex) pole of $A(s)$ as a function of the parameters of the theory. The position of the pole s can be parametrized by its "mass" m and "width" Γ through the relation $s=(m-\frac{1}{2}i\Gamma)^2$, though we should not take these terms literally when Γ is comparable to m. The pole traces out a curve in the s-plane as μ_t is varied. When μ_t is very large, the real and imaginary parts of the pole are both small, corresponding to a light, narrow resonance. As μ_t decreases, $\mathrm{Re}(s)$ increases, reaches a maximum and then begins to decrease, while $\mathrm{Im}(s)$ continues to increase. We refer to the pole position with maximum $\mathrm{Re}(s)$ as the "heaviest" resonance. This resonance is very broad, with Γ roughly equal to m. In the $O(4)$ model, the "heaviest" resonance is found [4] to have "mass" $m=845$ GeV and "width" $\Gamma=640$ GeV, and corresponds to a cut-off $\Lambda_c=4.9$ TeV and tachyon mass $\mu_t=8.4$ TeV. From the scaling property of eq. (4), the values of m and Γ corresponding to the heaviest resonance for the $O(N)$ model are $\sqrt{4/N}$ times those for the $O(4)$ model; the cut-off and tachyon mass scale in the same way.

We now turn to the $O(j)\times O(n)$ model of ref. [2]. The amplitudes are calculated in the limit j, $n\to\infty$ with the ratio j/n fixed; only diagrams which contribute to leading order in $1/(j+n)$ are included. The amplitude $A(s)$ is given by

$$A(s) = s\left\{f^2 - \frac{sj}{32\pi^2}\left[\ln\left(\frac{e\Lambda_c^2}{|s|}\right) + i\pi\Theta(s)\right] - \frac{sn}{32\pi^2}\left[\ln\left(\frac{\Lambda_c^2}{em_\psi^2}\right) - F_2(s, m_\psi)\right]\right\}^{-1}, \quad (6)$$

where

[#1] In ref. [5] eq. (2.16) should be corrected in accord with eq. (1) of this paper. The effect of this correction is that the W^-W^+ event rate discussed in this reference for the $O(4)$ model should be multiplied by ~ 2 to 2.5.

[#2] Our Λ_c is $\sqrt{e}$ times the one defined in ref. [4].

$$F_2(s,m)=-2+\sqrt{1-\frac{4m^2}{s}}\ln\left(\frac{\sqrt{4m^2-s}+\sqrt{-s}}{\sqrt{4m^2-s}-\sqrt{-s}}\right)$$

for $s<0$,

$$F_2(s,m)=-2+2\sqrt{-1+\frac{4m^2}{s}}\arctan\sqrt{\frac{s}{4m^2-s}}$$

for $0<s<4m^2$,

$$F_2(s,m)$$
$$=-2+\sqrt{1-\frac{4m^2}{s}}\left[\ln\left(\frac{\sqrt{s}+\sqrt{s-4m^2}}{\sqrt{s}-\sqrt{s-rm^2}}\right)-i\pi\right]$$

for $s>4m^2$, (7)

and the cut-off Λ_c (equal to M of ref. [2]) is related to the tachyon scale by

$$\Lambda_c=\frac{\mu_t}{\sqrt{e}}\exp\left\{\frac{-16\pi^2 f^2}{(j+n)\mu_t^2}+\frac{n}{2(j+n)}\left[\ln\left(\frac{m_\psi^2}{\mu_t^2}\right)\right.\right.$$
$$\left.\left.-\sqrt{1+\frac{4m_\psi^2}{\mu_t^2}}\ln\left(\frac{\sqrt{\mu_t^2+4m_\psi^2}-\mu_t}{\sqrt{\mu_t^2+4m_\psi^2}+\mu_t}\right)\right]\right\}. \tag{8}$$

Qualitatively speaking, for center-of-mass energies well below the ψ mass threshold, $s\ll 4m_\psi^2$, the $O(j)\times O(n)$ model behaves like the $O(j)$ model; the pseudo-Goldstone bosons play a little role. On the other hand, well above the threshold, $s\gg 4m_\psi^2$, the $O(j)\times O(n)$ model behaves like the $O(j+n)$ model.

Having obtained the amplitude (6), one sets $f=250$ GeV and $j=4$; the exact Goldstone bosons in this model correspond to the longitudinal W's. Three independent parameters now specify the model: the number of pseudo-Goldstone bosons n, their mass m_ψ, and the tachyon mass μ_t. (Again, the model is only valid at energy scales well below the tachyon mass.)

We now compare the total event rates in the $O(4)\times O(n)$ model for different values of n. As in the $O(N)$ model, the amplitude (6) has a complex pole, whose real part increases and then decreases as μ_t varies. We choose the parameter m_ψ so that the resonance is well above the pseudo-Goldstone mass threshold, where the model essentially behaves like the $O(4+n)$ model. Thus, using the scaling behavior described earlier, the "heaviest" resonance of the $O(4)\times O(n)$ model has $m\simeq\sqrt{4/(4+n)}\times 845$ GeV and $\Gamma\simeq\sqrt{4/(4+n)}\times 640$ GeV; the corresponding tachyon mass and cut-off also scale as $\sqrt{4/(4+n)}$ relative to the $O(4)$ model. We choose the tachyon mass μ_t for each value of n to correspond to the "heaviest" resonance of that model. Note that as n increases, the resonance moves to smaller mass m; the width to mass ratio of the resonance is of course independent of n.

Above the pseudo-Goldstone mass threshold, where the model behaves like the $O(4+n)$ model, the amplitude (6) has the scaling property

$$A(s)=\frac{1}{4+n}F(\tilde{s}),\quad \tilde{s}=(4+n)s, \tag{9}$$

This follows from eq. (5) because the cut-off Λ_c for the heaviest resonance scales as $1/\sqrt{4+n}$, and so $\tilde{\Lambda}_c$ is independent of n. Since the amplitude (9) scales as $1/(4+n)$, it would seem that the scattering rate becomes smaller as n increases. On the other hand, for larger n, the resonance occurs at lower invariant mass, where the WW parton luminosity L_{WW} is higher. We now show that the two effects cancel each other.

In the effective-W approximation [6,7] we have

$$\sigma_{pp\to WW\to WW}(S)=\int_{\tau_{min}}^{1}d\tau\frac{dL_{WW}}{d\tau}\sigma_{WW\to WW}(\tau S),$$

$$\sigma_{WW\to WW}(\tau S)=\int d\Omega\frac{1}{2\tau S}|\mathcal{M}(\tau S)|^2, \tag{10}$$

where $\sqrt{S}$ is the center-of-mass energy of the pp collider, $\sqrt{s}=\sqrt{\tau S}$ is the invariant mass of the WW pair, $\tau_{min}=4M_W^2/S$, and $d\Omega$ integrates over the direction of the out-going W in the WW center-of-mass frame. The parton luminosity $(dL_{WW}/d\tau)$ scales [#3] approximately as $1/\tau^2$ for $M_{WW}\lesssim 1$ TeV at the SSC. By rewriting eq. (10) in terms of $\tilde{\tau}=(4+n)\tau$, using eq. (9), we find that $\sigma_{pp\to WW\to WW}(S)$ is actually independent of n. Thus we conclude that, although the amplitude decreases as n increases, the total elastic event rate stays the same.

To see how large the event rate actually is, we choose $n=8$ and $m_\psi=125$ GeV. We expect from our scaling arguments that the heaviest resonance for $n=8$ will have $m\simeq\sqrt{\frac{4}{12}}\times 845=490$ GeV and $\Gamma\simeq\sqrt{\frac{4}{12}}\times 640=370$ GeV. Indeed, using eq. (6) explicitly, we find that the heaviest resonance occurs for $m=485$ GeV and $\Gamma=350$ GeV, corresponding to a tachyon

[#3] Based on fig. 5 of ref. [7].

mass $\mu_t = 4.3$ TeV. Because $E_W \gg M_W$, use of the equivalence theorem is probably justified for these parameters.

To obtain the event rate for the $O(4)\times O(8)$ model with parameters $m_\psi = 125$ GeV and $\mu_t = 4.3$ TeV, we fold the amplitudes with the parton luminosities. We find that the elastic W^-W^+ event rate for $M_{WW} \geqslant 350$ GeV at the SSC (with $\sqrt{S} = 40$ TeV and integrated luminosity 10^4 pb^{-1}) is about 0.5 pb. This rate is about the same as the total rate (0.6 pb) for the $O(4)$ model (i.e., the $n=0$ limit) with a resonance with $m = 845$ GeV and $\Gamma = 640$ GeV, as expected from the scaling argument given above. Moreover, this rate is not much smaller than the rate (3.4 pb) for a 500 GeV standard model Higgs boson, with width 64 GeV, produced via the W-fusion process. The Z^0Z^0 event rate in the $O(4)\times O(8)$ model is about half the W^+W^- event rate. We have not included in these rates W pairs produced by either quark or gluon fusion, restricting our consideration to WW scattering.

Many studies have been performed on detecting Higgs bosons at the SSC. It has been shown that a ~ 500 GeV standard model Higgs boson can be detected using the "gold-plated" mode alone, and does not require the application of techniques such as jet-tagging [8] and/or jet-vetoing [9] used in studying TeV WW interactions. However, these techniques, together with others, such as measuring the charged particle multiplicity of the event [10] or testing the fraction of longitudinal W's [5] could be used to further improve the signal-to-background ratio to study a ~ 500 GeV standard model Higgs boson produced via WW fusion processes. We think it is clear that similar strategies could be applied to detect the ~ 500 GeV resonance in the $O(4)\times O(8)$ model discussed above.

For the $O(4)\times O(32)$ model of ref. [2], with parameters $m_\psi = 125$ GeV and $\mu_t = 2.5$ TeV, the resonance ($m = 275$ GeV and $\Gamma = 120$ GeV) is in a region where the energy E_W of the longitudinal W in the WW center-of-mass frame is less than twice M_W [#4]. To see whether such a resonance could be detected would require a detailed Monte Carlo study, which we will not perform in this paper. We would argue, however, that with ~ 5000 extra longitudinal W^+W^- pairs and ~ 2500 extra longitudinal Z^0Z^0 pairs produced in one standard SSC year, this signal could probably be observed with appropriate detectors.

In this note, we have considered the $O(4)\times O(n)$ model presented in ref. [2] containing many inelastic channels in the WW scattering process. We demonstrated through a simple scaling argument that, although the amplitude for elastic WW scattering in this model decreases as the number n of inelastic channels increases, the total elastic event rate remains more or less the same. (We choose the parameters for each model to give the "heaviest" possible resonance.) This rate is about the same as that for the $O(4)$ model, with no inelastic channels and a heavy resonance.

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[#4] The equivalence theorem, which allows us to identify the Goldstone bosons ϕ with the longitudinal gauge bosons W in the WW scattering processes, requires $E_W \gg M_W$.

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