Possible new resonance from $W_L W_L - h h$ interchannel coupling

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We propose and theoretically study a possible new resonance caused by strong coupling between the Higgs-Higgs and the $W_L W_L (Z_L Z_L)$ scattering channels, without regard to the intensity of the elastic interaction in either channel at low energy (that could be weak as in the Standard Model). We expose this channel-coupling resonance from unitarity and dispersion relations encoded in the Inverse Amplitude Method, applied to the Electroweak Chiral Lagrangian with a scalar Higgs.

The LHC experiments CMS and ATLAS [1] have seemingly found what looks like a Higgs boson (mainly an excess of four-lepton events and two-photon events at 125 GeV suggestive of scalar quantum numbers). This finding has been widely discussed, but less recognized is the equally interesting fact that no new physics beyond the Standard Model (SM) appears up to energies of 600-700 GeV for generic searches [2], as shown in Fig. 1. Therefore the unknown system giving rise to the electroweak symmetry breaking (the Symmetry Breaking Sector (SBS)) from $SU(2)_L \times U(1)_Y$ to $U(1)_{em}$ should contain four low-mass states: the three would-be Goldstone bosons $\omega^a (a = 1,2,3)$ responsible for the $W^\pm$ and $Z$ masses and the recently discovered particle $h$. Because this Higgs-like boson turns out to be light and the spectrum is gapped up to the scale of any new physics, it is natural to think of it also as an approximation (composite) Goldstone boson (GB) itself [3]. For instance, as one of GB corresponding to the spontaneous symmetry breaking from a group $G$ to a group $H$ with $\dim(G) - \dim(H) = 4$. This is for example the case of the so-called MCHM (Minimal Composite Higgs Model [4], with $G = SO(5)$ and $H = SO(4)$). Another exciting possibility is that the Higgs is the dilaton [5] (the Goldstone boson associated with spontaneous breaking of the scale symmetry of the SBS).

Assuming the approximate, well established $SU(2)_L + R$ custodial symmetry, the low-energy GB dynamics can be properly described by a $SU(2)_L \times U(1)_Y$ gauged nonlinear effective Lagrangian [6] [7], which is an extension of the former Higgsless Electroweak Chiral Lagrangian [8]. Thus the three $\omega^a$ GB fields parametrize the coset $SU(2)_L \times SU(2)_R / SU(2)_{L+R}$ and the Higgs-like boson $h$ is a custodial isospin singlet.

In this work we are concerned with the $W_L W_L, Z_L Z_L$ and $h h$ scattering. This is because thanks to the Equivalence Theorem [9], we can get information about the unknown SBS of the SM by studying the GB dynamics, whose amplitudes approximate well those of the longitudinal vector bosons $W_L$ and $Z_L$ ($W_L$ for short in the following) of the SM for energies well above the $W$ mass ($E >> M_W$). In this regime we can also neglect the $h$ boson mass since $M_h = 125 \text{ GeV} \sim M_W$. Then the relevant Lagrangian for $W_L W_L$ and $h h$ scattering is [10]:

$$L = \left(1 + 2a \frac{h}{v} + b \frac{h^2}{v^2} \right) \frac{\partial_a \omega^a \partial^\mu \omega^b}{2} \left(\delta^{ab} + \frac{\omega^a \omega^b}{v^2}\right) + \frac{4a_1}{v^4} (\partial_a \omega^a \partial^\mu \omega^b)^2 + \frac{4a_2}{v^4} (\partial_a \omega^a \partial^\nu \omega^b)^2 + \frac{2d}{v^3} \partial_a h \partial^\mu h \partial_a \omega^b \partial^\nu \omega^a + \frac{2e}{v^3} (\partial_a h \partial^\nu \omega^a)^2 + \frac{1}{2} \partial_{\mu} h \partial^\mu h + \frac{g}{v^4} (\partial_{\mu} h \partial^\mu h)^2$$

which should be valid for $M_W, M_h << E << 4\pi v \approx 3 \text{ TeV}$. Thus we have set $M_h = M_W = M_Z \approx 0$. Different SBS dynamics can be modeled by a proper tuning of the parameters $a, b, a_1, a_2, d, e$ and $g$. The last five of them must be renormalized to some scale $\mu$ to absorb the one-loop divergencies coming from the lowest order (LO), i.e. the first term in the Lagrangian, in a similar way as in Chiral Perturbation Theory [11]. In the SM $a^2 = b = 1$ and the rest of the tree level parameters vanish ($a_1 = a_2 = d = e = g = 0$). As is well known in this particular case we get a linear theory which is renormalizable in the standard way and weakly interacting for light $h$.

From the Lagrangian density in Eq. (1), the LO partial waves with $I = J = 0$ ($I$ being the custodial isospin) can

![Figure 1: Electroweak-symmetry breaking sector of the Standard Model after the LHC run I: there are four “low-energy” bosons and any new physics is split by a mass gap.](image-url)
be easily computed and one finds:

\begin{align*}
A_0(\omega \omega \rightarrow \omega \omega) &= \frac{s}{16\pi v^2}(1 - a^2) \\
T_0(hh \rightarrow hh) &= 0 \\
M_0(\omega \omega \rightarrow hh) &= \frac{\sqrt{3}s}{32\pi v^2}(a^2 - b).
\end{align*}

From these low-energy theorems we can expect strong \( W_L W_L \) elastic scattering whenever \( a^2 \neq 1 \) (as Eq. \( 2 \) grows with \( s = E_{cm}^2 \)). For \( b \neq a^2 \) we have strong mixing between the \( W_L W_L \) and \( hh \) channels [12, 13] even in the case \( a = 1 \) (no LO contribution to \( W_L W_L \) elastic scattering). As we will show below, this strong mixing gives rise to a new resonance in the \( I = J = 0 \) channel for an important region of the available \((a, b)\) parameter space after taking into account the experimental information coming from the LHC. At next to leading order (NLO) all three amplitudes in Eq. \( 2 \) acquire one-loop contributions of order \( O(s^2) \). The elastic ones \( A_1 \) and \( T_1 \) are accompanied by logarithmic left and right cuts (LC and RC respectively) in the s-plane, entailing an imaginary part for physical energy corresponding to \( s \) just above the RC. Those amplitudes have already been reported in recent literature [10, 14]. The divergencies appearing for massless \( W \) and \( h \) can be absorbed by renormalization of the \( a_4, a_5, d, e \) and \( g \) parameters but no \( a \) or \( b \) renormalization is needed in this case.

In this work we want to focus on the very interesting phenomenon of the strong mixing appearing whenever \( a \neq b^2 \). In order to emphasize this point we will concentrate first in the particular case where \( a = 1 \) (no direct \( W_L W_L \) LO elastic scattering) and the rest of the parameters except \( b \) vanish [15]. As the renormalized parameters depend on the renormalization scale \( \mu \), that particularly simple choice requires to set this scale to some given value that now plays the role of an ultraviolet cutoff. For definiteness we will take \( \mu = 4\pi v \approx 3 \) TeV which is roughly the limit of applicability of our effective theory. The relevant NLO \( I = 0 \) partial waves read [10]:

\begin{align*}
A_1(\omega \omega \rightarrow \omega \omega) &= \frac{s^2(1-b)^2}{256\pi^4 v^4} \times (3) \\
\left( \frac{17}{9} - \frac{1}{6} \log \left( \frac{s}{\mu^2} \right) \right) - \frac{3}{4} \log \left( \frac{-s}{\mu^2} \right) \\
T_1(hh \rightarrow hh) &= \frac{s^2(1-b)^2}{32\pi^3 v^4} \times (4) \\
\left( \frac{1}{3} - \frac{1}{16} \log \left( \frac{s}{\mu^2} \right) \right) - \frac{3}{32} \log \left( \frac{-s}{\mu^2} \right) \\
M_1(\omega \omega \rightarrow hh) &= \frac{\sqrt{3}(1-b)s}{32\pi v^2} + (5) \\
\frac{\sqrt{3}s^2(1-b)}{9216\pi^3 v^4} \left( \frac{1}{2} - \log \left( \frac{s}{\mu^2} \right) \right).
\end{align*}

These partial waves have adequate analytical properties featuring a LC and also a RC just under the physical region \( s = E_{cm}^2 + i\epsilon \). However, unitarity is satisfied only perturbatively, with \( \text{Im}A_1 = \text{Im}T_1 = |M_0|^2 \) on the RC. Notice also that the cross-channel amplitude \( M_1 \) has only a LC and is thus purely real \( (\text{Im}(M_0 + M_1) = 0) \) on the RC. As the amplitudes grow with \( s \) they will eventually violate the unitarity bound (for example \( \text{Im}A \leq 1 \)). Grouping the two coupled channels in matrix form, \( F = F_0 + F_1 + \ldots \)

\begin{align*}
F &= \begin{pmatrix} 0 & M_0 \\ M_0 & 0 \end{pmatrix} + \begin{pmatrix} A_1 & M_1 \\ M_1 & T_1 \end{pmatrix} + \ldots
\end{align*}

the perturbative unitarity relation satisfied is \( \text{Im}F_1 = F_0^* F_1 \); but exact unitarity requires \( \text{Im}F = F^* F \) instead. However there is a very well known method, based on dispersion relations, called the Inverse Amplitude Method (IAM) [16], that allows to fully unitarize the perturbative partial waves, even in the coupled channel case [17]. The resulting amplitudes are given by \( F^{IAM} := F_0(F_0^* - F_1)^{-1} F_0 \). The IAM amplitudes still have the analytical properties found above but now they fulfill exact unitarity. In addition, the determinant of \( F_0 - F_1 \) appearing in the denominators allows for the possibility of having poles in the second Riemann sheet for some regions of the parameter space. When they are close enough to the physical region, those poles have the natural interpretation of dynamical resonances. The IAM amplitudes have been extensively used to fit meson-meson scattering data [18]. The method has also been applied to \( W_L W_L \) elastic scattering [19] where resonances were found in different channels in terms of the \( a_4 \) and \( a_5 \) parameters. In this work we are rather interested in the pure coupled channel resonances appearing even for \( a = 1 \). The IAM amplitudes are given in this case by:

\[ F^{IAM} = \frac{M_0^2}{(M_0 - M_1)^2 - A_1 T_1} \left( \frac{A_1}{M_0 - M_1} - \frac{M_0 - M_1}{T_1} \right) \]

which perturbative expansion coincides with Eq. \( 4 \) up to and including order \( s^2 \) but that satisfies exact unitarity, \( \text{Im}F^{IAM} = (F^{IAM})^* F^{IAM} \) as it can be easily checked. The IAM method resums the imaginary parts of diagrams like the one shown in Fig. 2.

In Fig. 3 we show the square moduli of two distinct matrix elements of \( F^{IAM} (A^{IAM} \text{ and } M^{IAM}) \). For \( b \) nearly 1 the amplitude vanishes in the LHC region of interest (with \( W_L W_L \rightarrow W_L W_H \) reducing to the small SM amplitude). For larger \( b \) there is a resonant structure
most clearly seen in the inelastic process $\omega \omega \to hh$, but that leaves a trace also in the elastic amplitudes that grow almost vertically for large $b$ and even peak slightly up to the resonance mass, growing more calmly afterwards.

In order to explore in more detail the peaks in Fig. 3, we have analytically extended the complex log($-s/\mu^2$) of Eq. (3) to (5) to the second Riemann sheet, and indeed we found a pole (see Fig. 4) at the same point in all the channels. This pole is naturally interpreted as a dynamical resonance whenever it is close enough to the physical, real $s$. Then the position of the pole $s_0$ is related to the parameters of the resonance (mass $M$ and width $\Gamma$) as $\sqrt{s_0} = M - i\Gamma/2$. In the complex $s$ plane, the motion of the pole $s_0$ with $b$ is displayed in Fig. 5. For this, we numerically find the zeroes of the denominator of Eq. (7) in the second Riemann sheet. We show values for $b > 1$. Since that denominator is an even function of $b - 1$, the plot for $b < 1$ falls on top of the shown trajectory. The pole escapes to infinity as $b$ approaches 1, then returning along the same trajectory as $b$ increases beyond 1.

In Fig. 6 we show the mass and the width of the resonance associated to the pole. Currently there is no constraint on the $b$ parameter but, if other parameters are set
to the SM values, we can provide the bound $b \in (-1,3)$ (around the SM $b = 1$) because otherwise the resonance moves below 700 GeV where it would have already been seen by ATLAS and CMS. These exclusion limits on $b$ are of course uncertain by $(\omega a \sim W_f)$.

Once we have identified the new coupled-channel resonance by switching off the LO elastic $W_L W_L$ channel by setting $a = 1$, we can consider now the more general case of arbitrary $a$. From the LHC data we know this parameter must roughly belong to the interval $(0.7,1.3)$ while $b$ is much more unconstrained. As commented in the introduction, the simplest composite model where the two $\phi$ and $h$ show up as composite (pseudo) GB is the so called MCHM featuring the symmetry breaking pattern $SO(5)$ to $SO(4)$. In this model the $a$ and $b$ parameters are given by $a = \sqrt{1 - \xi}$ and $b = 1 - 2\xi$, where $\xi = \omega^2/f^2$ and $f$ is a new symmetry breaking higher scale. The relevant IAM partial waves in the general case $a \neq 1$ may be retrieved from $[10, 20]$. In Figs. 5 and 6 we show the position of the pole and the corresponding mass and width in terms of $\xi$ for the MCHM model. As can be seen this model features the new resonance studied here for an important range of the allowed $\xi$ parameter range ($0 < \xi < 0.5$). Even for the region where the pole is too far away from the real axis to be considered a resonance, it will produce a huge increment of the cross section for $W_L W_L$ and $hh$ production that could be probed at the LHC. This conclusion probably applies to other composite models beyond the MCHM too. For finite $a$ the resonance receives strength from the channel-coupling and also from elastic scattering, but we have shown $[13, 20]$ that the $\sigma$-like structure from elastic dynamics alone is much broader; for finite $b$ the resonance, as shown in figure 6, is significantly narrower and lighter due to the coupled-channel dynamics.

To conclude, we have made the case for an interesting potential phenomenon to be sought at the LHC run II and beyond; a new resonance in the $W_L W_L − hh$ coupled channels, caused by the channel-mixing interaction independently of the elastic interaction in either channel.

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