Constraints on vector resonances from a strong Higgs sector.

A. E. Cárcamo Hernández,1 Bastián Díaz Sáez,1 Claudio O. Dib,1 and Alfonso Zerwekh1  
1Department of Physics and Centro Científico-Tecnológico de Valparaíso  
Universidad Técnica Federico Santa María, Valparaíso, Chile

We consider a scenario of a composite Higgs arising from a strong sector. We assume that the lowest lying composite states are the Higgs scalar doublet and a massive vector triplet, whose dynamics below the compositeness scale are described in terms of an effective Lagrangian. Electroweak symmetry breaking takes place through a vacuum expectation value just as in the Standard Model, but with the vector resonances strongly coupled to the Higgs field. We determine the constraints on this scenario imposed by (i) the Higgs diphoton decay rate, (ii) the electroweak precision tests and (iii) searches of heavy resonances at the LHC in the final states $l^+l^−$ and $lν_l$ ($l = e, μ$), $τ^+τ^−$, $jj$, $t\bar{t}$, $WZ$, $WW$, $WH$ and $ZH$. We find that the heavy vector resonances should have masses that are constrained to be in the range 2.1 - 3 TeV. On the other hand, the mixing of the heavy vectors with the Standard Model gauge bosons is constrained to be in the range $\tan\theta \sim 0.1 - 0.3$, which is consistent with the assumption that the Higgs couples weakly to the Standard sector, even though it couples strongly to the heavy vector resonances.

I. INTRODUCTION.

The recent discovery of the Higgs boson at the LHC [11,2] provides the opportunity to directly explore the mechanism of electroweak symmetry breaking (EWSB). While this remarkable achievement implies severe constrains on many proposed extensions of the Standard Model (SM), an additional sector beyond our current knowledge is still needed in order to explain the dynamical origin of the electroweak scale and its stability [3]. A specific question in this context is whether this new sector is weakly or strongly interacting [4]. In the latter case, the Higgs boson is viewed as a composite state which must be accompanied by a plethora of new heavy composite particles [5-7]. In general, it is expected that the lightest states produced by the strong dynamics would correspond to spin-0 and spin-1 particles [5-8]. In these models the lightness of the Higgs can be explained in two different ways. One way is to consider the Higgs boson as a pseudo-Goldstone boson that appears after the breakdown of a suitable global symmetry [8,9,26]. A second way is to consider the Higgs boson as the modulus of an effective SU(2) doublet, where its lightness is due to particularities of the dynamics of the underlying theory [10,27,57]. For instance, there are evidences that quasi-conformal strong interacting theories such as walking technicolor may provide a light composite scalar [51-56]. It has also been shown that, in the effective low energy theory, the composite scalar may develop a potential that reproduces the standard Higgs sector [27]. In this scheme, the electroweak symmetry breaking is effectively described by a non-zero vacuum expectation value of the scalar arising from the potential, just as in the Standard Model. However, additional composite particles, like vector resonances, may also be expected to appear in the spectrum [58].

The main reason to consider strongly interacting mechanisms of EWSB as alternatives to the Standard Model mechanism based on a fundamental scalar is the so called hierarchy problem that arises from the Higgs sector of the SM [5-7]. This problem is indicative that, in a natural scenario, new physics should appear at scales not much higher than the EWSB scale, say around a few TeV, in order to stabilize the Higgs mass at a value much lower than the Planck scale ($\sim 10^{19}$ GeV). An underlying strongly interacting dynamics without fundamental scalars, which becomes non-perturbative somewhere above the EW scale, is a possible scenario that gives an answer to this problem.

While a composite Higgs boson is theoretically attractive because the underlying strong dynamics provides a comprehensive and natural explanation for the origin of the Fermi scale [5-7], the presence of additional composite states such as the vector triplets previously mentioned may, in principle, produce phenomenological problems. For instance one could expect that, at one loop level, they may produce sensible corrections to observables involving the Higgs boson. Consequently, an interesting quantity which can eventually reveal the influence of additional states is $\Gamma(h \rightarrow \gamma\gamma)$. In a previous work this decay channel was studied in a simple model with vector resonances and found it in general agreement with current experimental measurements in the limit where the Higgs boson is weakly coupled to the new resonances [42]. However, if the Higgs boson arises from a strongly interacting sector together with other heavy resonances, one should expect a strong coupling among them.
In this work, we want to investigate whether this strong coupling hypothesis is still compatible with the currently known phenomenology and, in general, whether composite models are viable alternatives to electroweak symmetry breaking, given the current experimental success of the Standard Model \[3\]. To be concrete, we describe the new sector by means of an effective model with minimal particle content, without referring to the details of the underlying strong dynamics. We use an effective chiral Lagrangian to describe the theory below the cutoff scale of the underlying strong interaction, assumed to be \( \Lambda = 4\pi v \sim 3 \text{ TeV} \). This low energy effective theory must contain the Standard Model spectrum and the extra composite scalar and vector multiplets.

The content of this paper goes as follows. In section II, we introduce our effective Lagrangian that describes the spectrum of the theory. Section III deals with the constraints arising from the Higgs diphoton decay rate and dijet exclusion limits. The constraints on the model parameter space arising from the oblique \( T \) and \( S \) parameters are discussed in Section IV. In section V, we describe the different decay channels of the heavy vector resonances. Finally, in section VI, we present the constraints of our model arising from LHC searches of heavy vector resonances. Finally, in section VII, we state our conclusions.

II. LAGRANGIAN FOR A HIGGS DOUBLET AND HEAVY VECTOR TRIPLET.

We want to formulate the scenario of EWSB triggered by a strongly coupled sector without referring to specific details of the underlying theory. This underlying theory, as it becomes strong at low energies, should generate the Higgs scalar multiplet as a composite field below a scale \( \Lambda \). This underlying theory, as it becomes strong at low energies, should generate the Higgs scalar multiplet as a composite field below a scale \( \Lambda \). The effective Lagrangian is expressed as

\[
\mathcal{L} = -\frac{1}{2} \left( F_{\mu
u}^{(1)} F^{(1)\mu\nu} \right) - \frac{1}{2} \left( F_{\mu
u}^{(2)} F^{(2)\mu\nu} \right) - \frac{1}{2} \left( B_{\mu
u} B^{\mu\nu} \right) + \frac{f^2}{2} \left( (D_\mu \Sigma)^\dagger D^\mu \Sigma \right) + \left( D_\mu \Phi \right)^\dagger D^\mu \Phi - \mu^2 \Phi^\dagger \Phi + \frac{\lambda}{4} (\Phi^\dagger \Phi)^2
\]

+ \frac{\beta}{2} \left( \Phi^\dagger \Phi \right)^2 \left( (D_\mu \Sigma)^\dagger D^\mu \Sigma \right) + \bar{\psi}_{iL} \gamma^\mu D_\mu \psi_{iL} + \bar{\psi}_{iR} \gamma^\mu D_\mu \psi_{iR} + y_{ij} \bar{\psi}_{iL} \Sigma \Phi \psi_{jR} + \bar{y}_{ij} \bar{\psi}_{iL} \Sigma \Phi^c \psi_{jR} + h.c.
\]

(1)

| Fields | \( SU(2)_1 \) | \( SU(2)_2 \) | \( U(1)_Y \) |
|--------|---------------|---------------|--------------|
| \( \Sigma \) | 2 | 2 | 0 |
| \( \Phi \) | 1 | 2 | 1/2 |
| \( Q_L^i \) | 1 | 1 | 2/3 |
| \( U_R^i \) | 1 | 1 | -1/3 |
| \( D_R^i \) | 2 | 1 | -1/2 |
| \( L_L^i \) | 1 | 2 | -1 |
| \( e_R^i \) | 1 | 1 | 0 |
| \( N_R^i \) | 1 | 1 | 0 |

Table I. Field charge assignments under the full gauge group \( SU(2)_1 \times SU(2)_2 \times U(1)_Y \). The \( i \) index runs from 1 to 3.
Here $F^{(1)\mu
u}$, $F^{(2)\mu
u}$ and $B^{\mu
u}$ are the gauge field tensors, the brackets $\langle \rangle$ denote the trace in the corresponding group indices, $f_\Sigma$ is the analog of a decay constant for the extra would-be Goldstones expressed non-linearly in the field $\Sigma$ (these are absorbed as the longitudinal components of the heavy vector composites), while $\lambda$ and $\mu$ are the SM parameters of the Higgs potential, and $\beta$ is the coefficient of a mixing term allowed by the symmetry. The value of $\beta$ is not easy to isolate from other parameters in observable quantities, so for the sake of simplicity from now on we fix $\beta = 2$. Finally, the covariant derivatives are:

$$D_\mu \Sigma = \partial_\mu \Sigma - ig_1 A^{(1)}_\mu \Sigma + ig_2 A^{(2)}_\mu \Sigma, \quad D_\mu \Phi = \partial_\mu \Phi - ig_2 A^{(2)}_\mu \Phi - i g' B_\mu \Phi,$$

$$D_\mu \psi_{1L} = \partial_\mu \psi_{1L} - ig_1 A^{(1)}_\mu \psi_{1L} - ig' Y_{fL} B_\mu \psi_{1L}, \quad D_\mu \psi_{1R} = \partial_\mu \psi_{1R} - ig' Y_{fR} B_\mu \psi_{1R},$$

(2)

where $A^{(n)}_\mu = \frac{1}{2} \tau^a A^{(n)a}_\mu$ ($n = 1, 2$) are the Hermitian gauge field matrices corresponding to the $SU(2)_n$ gauge fields $A^{(n)}_\mu$, $n = 1, 2$ respectively.

Notice that $\Phi$ is coupled to $A^{(2)}_\mu$ but not to $A^{(1)}_\mu$, and so it is more strongly coupled to the heavy vectors than to the SM gauge fields. In addition, left handed SM fermionic fields will couple mainly to SM gauge fields, which are primarily contained in $SU(2)_1$. The scalar doublet will correspond to the SM Higgs field, which can be expressed as usual by:

$$\Phi = \left( \frac{1}{\sqrt{2}} \begin{pmatrix} G^+ \\ (v + h + iG^0) \end{pmatrix} \right),$$

(3)

where the field $h$ is the Higgs boson, while $G^+$ and $G^0$ are the would-be Goldstones that will be absorbed after EWSB.

The spontaneous breaking of the extra gauge symmetry can be formulated by taking $\Sigma = 1$ (in the unitary gauge). The Lagrangian then takes the following form:

$$\mathcal{L} = \mathcal{L}_{\text{gauge}} + \frac{f_\Sigma^2 g^2}{4 \sin^2 2\theta} V^{(2)\mu} V^{(2)\mu} + (D_\mu \Phi)^\dagger D^\mu \Phi - \mu^2 \Phi^\dagger \Phi + \frac{\lambda}{4} (\Phi^\dagger \Phi)^2 + \frac{\beta g^2}{\sin^2 2\theta} (\Phi^\dagger \Phi) V^{(2)\mu} V^{(2)\mu} + i \bar{\psi}_{1L} \gamma^\mu D_\mu \psi_{1L} + i \bar{\psi}_{1R} \gamma^\mu D_\mu \psi_{1R} + y_{ij} \bar{\psi}_{1L} \Phi \psi_{1R} + \bar{y}_{ij} \psi_{1L} \Phi^c \psi_{1R} + h.c.,$$

(4)

where the covariant derivate is now rewritten as follows:

$$D_\mu \Phi = \partial_\mu \Phi - ig \tau^a \frac{1}{2} V^{(1)\mu} \Phi - ig' \tau^a \frac{1}{2} V^{(2)\mu} \Phi - i g' B_\mu \Phi,$$

(5)

with the vector fields given by:

$$V^{(1)}_\mu = \cos \vartheta A^{(1)}_\mu + \sin \vartheta A^{(2)}_\mu, \quad V^{(2)}_\mu = - \sin \vartheta A^{(1)}_\mu + \cos \vartheta A^{(2)}_\mu, \quad \tan \vartheta = \frac{g_1}{g_2},$$

(6)

and the couplings:

$$g = \frac{g_1 g_2}{\sqrt{g_1^2 + g_2^2}}, \quad \tilde{g} = \frac{g_2^2}{\sqrt{g_1^2 + g_2^2}}.$$

(7)

At this stage, the fields $V^{(1)}_\mu$ remain massless but $V^{(2)}_\mu$ acquire mass proportional to $f_\Sigma$, as shown in Eq. 4. When the Higgs boson acquires a vacuum expectation value $\langle \Phi \rangle = \frac{1}{\sqrt{2}} (0, v)^T$, from Eq. 4 it follows that the squared mass matrices for the neutral and charged gauge bosons are given by:

$$M_{N_\pm}^2 = \frac{v^2}{4} \begin{pmatrix} g^2 & g g' & g \tilde{g} \\ -g g' & g' \tilde{g} & -g \tilde{g} \\ g \tilde{g} & -g' \tilde{g} & g^2 k^2 \end{pmatrix}, \quad M_{C_\pm}^2 = \frac{v^2}{4} \begin{pmatrix} g^2 & g \tilde{g} & g' \tilde{g} \\ g \tilde{g} & k^2 g' \tilde{g} \\ g' \tilde{g} & k^2 g^2 \end{pmatrix},$$

(8)

with $k^2 = 1 + \frac{2 f_\Sigma^2 + \beta v^2}{v^2 \cos^4 \vartheta}$.

The masses of the gauge bosons are given by diagonalization of these mass matrices:
\[ M_A = 0, \]
\[ M_Z = \frac{v}{2\sqrt{2}} \sqrt{g^2 + g'^2 + \tilde{g}^2 k^2 - \sqrt{g^4 k^4 + (g^2 + g'^2)(g'^2 + g'^2 + 2(2 - k^2) \tilde{g}^2)},} \]
\[ M_{\rho} = \frac{v}{2\sqrt{2}} \sqrt{g^2 + g'^2 + \tilde{g}^2 k^2 + \sqrt{g^4 k^4 + (g^2 + g'^2)(g'^2 + g'^2 + 2(2 - k^2) \tilde{g}^2)},} \]
\[ M_{W^{\pm}} = \frac{v}{2\sqrt{2}} \sqrt{k^2 \tilde{g}^2 + g^2} - \sqrt{(g^2 - k^2 \tilde{g}^2)^2 + 4g^2 g'^2}, \]
\[ M_{\rho^0} = \frac{v}{2\sqrt{2}} \sqrt{k^2 \tilde{g}^2 + g^2 + \sqrt{(g^2 - k^2 \tilde{g}^2)^2 + 4g^2 g'^2},} \]
\[ \text{and the physical neutral and charged gauge bosons are given by:} \]
\[ \begin{pmatrix} A_\mu \\ Z_\mu \\ \rho^0_\mu \end{pmatrix} = \begin{pmatrix} \cos \theta_W & \sin \theta_W & 0 \\ -\cos \gamma \sin \theta_W & \cos \gamma \cos \theta_W & -\sin \gamma \\ -\sin \gamma \sin \theta_W & \cos \theta_W \sin \gamma & \cos \gamma \end{pmatrix} \begin{pmatrix} B^0_\mu \\ V^{(1)}_\mu \\ V^{(2)}_\mu \end{pmatrix}, \]
\[ \begin{pmatrix} W^{\pm}_\mu \\ \rho^\pm_\mu \end{pmatrix} = \begin{pmatrix} \cos \kappa & -\sin \kappa \\ \sin \kappa & \cos \kappa \end{pmatrix} \begin{pmatrix} V^{(1)}_\mu \\ V^{(2)}_\mu \end{pmatrix}, \]
\[ \text{where, besides the standard } \theta_W, \text{ the additional mixing angles } \kappa \text{ and } \gamma \text{ are:} \]
\[ \tan 2\kappa = \frac{\sqrt{4M^2_{\rho^0} - g^2 v^2}}{2 \left( M^2_{\rho^0} + M^2_W \right) - g^2 v^2}, \]
\[ \tan 2\gamma = \frac{2g \sqrt{g^2 + g'^2}}{g^2 k^2 - g^2 - g'^2}. \]

At this stage the electroweak symmetry is finally broken and the only remaining massless vector boson is the photon field \( A_\mu \).

### III. CONSTRAINTS FROM HIGGS DECAY INTO TWO PHOTONS

In the Standard Model, the \( h \to \gamma \gamma \) decay is dominated by \( W \) loop diagrams which can interfere destructively with the subdominant top quark loop. In our strongly coupled model, the \( h \to \gamma \gamma \) decay receives additional contributions from loops with charged \( \rho^0_\mu \), as shown in Fig. II. The explicit form for the \( h \to \gamma \gamma \) decay rate is:

\[ \Gamma (h \to \gamma \gamma) = \frac{\alpha^2 \mu m_h^3}{256 \pi^2 v^2} \left| \sum_f N_f Q_f^2 F_{1/2} (x_f) + a_{hW} F_1 (x_W) + a_{h\rho^+\rho^-} F_1 (x_\rho) \right|^2, \]

where:

\[ a_{hW} = \frac{vg^2}{2} \left\{ \frac{\cos \kappa - \cot \vartheta \sin \kappa}{\sin^2 \vartheta \cos^2 \vartheta} \right\} \frac{v}{2M^2_W}, \]
\[ a_{h\rho^+\rho^-} = \frac{vg^2}{2} \left\{ \frac{\sin \kappa + \cot \vartheta \cos \kappa}{\sin^2 \vartheta \cos^2 \vartheta} \right\} \frac{v}{2M^2_{\rho^0}}. \]

Here \( x_f \) are the mass ratios \( x_f = m^2_f / M^2_W \), with \( M_1 = m_f, M_W \) and \( M_{\rho^0} \), respectively, \( \alpha_{em} \) is the fine structure constant, \( N_C \) is the color factor \( (N_C = 1 \text{ for leptons, } N_C = 3 \text{ for quarks}) \), and \( Q_f \) is the electric charge of the fermion in the loop. From the fermion loop contributions we will keep only the dominant term, which is the one involving the top quark.

The dimensionless loop factors \( F_{1/2} (x) \) and \( F_1 (x) \) (for particles of spin 1/2 and 1 in the loop, respectively) are:

\[ F_{1/2} (x) = 2 \left[ x + (x - 1) f (x) \right] x^{-2}, \]
\[ F_1 (x) = - \left[ 2x^2 + 3x + 3(2x - 1) f (x) \right] x^{-2}. \]
with
\[
f(x) = \begin{cases} \arcsin^2 \sqrt{x}, & \text{for } x \leq 1 \\ -\frac{1}{4} \left[ \ln \left( \frac{1 + \sqrt{1-x}}{1 - \sqrt{1-x}} \right) - i\pi \right], & \text{for } x > 1. \end{cases}
\] (17)

In what follows, we want to determine the range of values for the mass $M_\rho$ of the heavy vector resonances and the mixing angle $\vartheta$, consistent with the Higgs diphoton signal strength measured by the ATLAS and CMS collaborations at the LHC. To this end, we introduce the ratio $R_{\gamma\gamma}$, which corresponds to the Higgs diphoton signal strength that normalises the $\gamma\gamma$ signal predicted by our model relative to that of the SM:

\[
R_{\gamma\gamma} = \frac{\sigma(pp \to h) \Gamma(h \to \gamma\gamma)}{\sigma(pp \to h)_{SM} \Gamma(h \to \gamma\gamma)_{SM}} \simeq \frac{\Gamma(h \to \gamma\gamma)}{\Gamma(h \to \gamma\gamma)_{SM}}.
\]

This normalization for $h \to \gamma\gamma$ was also done in Refs. [66–68]. Here we have used the fact that in our model, single Higgs production is also dominated by gluon fusion as in the Standard Model.

Fig. 2 shows the sensitivity of the ratio $R_{\gamma\gamma}$ as a function of $M_\rho$ for several values of $\tan \vartheta$. The solid curves from top to bottom correspond to $\tan \vartheta = 0.1, 0.2$ and 0.3. The horizontal dashed lines are the minimum and maximum values of the ratio $R_{\gamma\gamma}$ inside the 1σ experimentally allowed range by CMS and ATLAS, namely $1.14 \pm 0.23$ and $1.17 \pm 0.27$, respectively [69, 70].

Fig. 2 shows the sensitivity of the ratio $R_{\gamma\gamma}$ under variations of $M_\rho$ for several values of $\tan \vartheta$. The curves from top to bottom correspond to $\tan \vartheta = 0.1, 0.2, 0.3$. The ratio $R_{\gamma\gamma}$ decreases slowly when the heavy vector masses are increased.

As shown, our model successfully accommodates the current Higgs diphoton decay rate constraints.

Figure 1. One loop Feynman diagrams in the Unitary Gauge contributing to the $h \to \gamma\gamma$ decay.

Figure 2. The ratio $R_{\gamma\gamma}$ as a function of $M_\rho$ for several values of $\tan \vartheta$. The solid curves from top to bottom correspond to $\tan \vartheta = 0.1, 0.2$ and 0.3. The horizontal dashed lines are the minimum and maximum values of the ratio $R_{\gamma\gamma}$ inside the 1σ experimentally allowed range by CMS and ATLAS, namely $1.14 \pm 0.23$ and $1.17 \pm 0.27$, respectively [69, 70].
Figure 3. Lower bound on $\tan \vartheta$ vs. $M_\rho$, according to the constraint imposed by the Higgs diphoton decay rate $h \to \gamma \gamma$ at the LHC.

A more exhaustive study of the allowed values of $\vartheta$ for different $M_\rho$ is shown in Fig. 3. The observed Higgs diphoton decay rate at the LHC excludes the white region below the curve in the figure, corresponding to too small values of $\vartheta$: for such small $\vartheta$ values the Higgs boson would couple too strongly to the heavy vector resonances, increasing the Higgs diphoton decay rate beyond the observed values. In addition, the heavy vector contribution to the Higgs diphoton decay rate scales as $1/M_\rho^2$ due to the heavy vector propagator and consequently, as Fig. 3 shows, the tightest lower bound $\tan \vartheta \gtrsim 0.15$ is obtained at the lower end of $M_\rho$; for larger masses of the vector resonances the $\vartheta$ values are less restricted.

IV. CONSTRAINTS FROM THE T AND S PARAMETERS

The inclusion of the extra composite particles also modifies the oblique corrections of the SM, the values of which have been extracted from high precision experiments. Consequently, the validity of our model depends on the condition that the extra particles do not contradict those experimental results. These oblique corrections are parametrized in terms of the two well known quantities $T$ and $S$. The $T$ parameter is defined as [71–73]:

$$T = \frac{\Pi_{33} (0) - \Pi_{11} (0)}{M_W^2 \alpha_{em} (m_Z)},$$  

(18)

where $\Pi_{33} (0)$ and $\Pi_{11} (0)$ are the vacuum polarization amplitudes at $q^2 = 0$ for the propagators of the gauge bosons $A^{(1)}_{\mu}, A^{(2)}_{\mu},$ and $A^{(3)}_{\mu}$, respectively, which are those that couple to the external fermions in the process $e^+ e^- \to f \bar{f}$ [73].

Figure 4. One loop Feynman diagrams contributing to the $T$ parameter. The fields are those in Eqs. (2) and (3).

In turn, the $S$ parameter is defined as [71–73]:

$$S = \left. \frac{4 \sin^2 \theta_W}{\alpha_{em} (m_Z)} \frac{g}{g'} \frac{d}{dq^2} \Pi_{30} (q^2) \right|_{q^2=0},$$  

(19)
where $\Pi_{30}(q^2)$ is the vacuum polarization for the propagator mixing of $A_{3}^{(1)}$ and $B_{3}$. The most important Feynman diagrams contributing to the $T$ and $S$ parameters are shown in Figures 4 and 5. We computed these oblique $T$ and $S$ parameters in the Landau gauge for the SM gauge bosons and would-be-Goldstone bosons, where the global $SU(2)_L \times U(1)_Y$ symmetry is preserved. We can separate the contributions to $T$ and $S$ from the SM and extra physics as $T = T_{SM} + \Delta T$ and $S = S_{SM} + \Delta S$, where

$$T_{SM} = -\frac{3}{16\pi} \ln \frac{m_T^2}{m_W^2} + \frac{3m_t^2}{32\pi^2 \alpha(Z)v^2},$$

$$S_{SM} = \frac{1}{12\pi} \ln \left( \frac{m_T^2}{m_W^2} \right) + \frac{1}{2\pi} \left[ 3 - \frac{1}{2} \ln \left( \frac{m_T^2}{m_b^2} \right) \right],$$

while $\Delta T$ and $\Delta S$ contain all the contributions involving the extra particles. The dominant one-loop contribution to $\Delta T$ and $\Delta S$ in our model are:

$$\Delta T = \tan^2 \vartheta \, T_{SM} - \frac{3\beta M_W^2}{16\pi \cos^2 \vartheta \sin^2 \vartheta \cos^2 \theta_W} F(M_B, m_h, M_{A^2}),$$

$$\Delta S = \frac{1 - \cos \vartheta}{\cos \vartheta} \, S_{SM} + \frac{2\beta M_W^2}{\pi \sin^2 \vartheta \cos \vartheta} \left[ G_1(M_{A^2}, m_h) - \frac{1}{4M_{A^2}^2} G_2(M_{A^2}, m_h) \right],$$

where

$$F(m_1, m_2, m_3) = \frac{m_3^2}{(m_1^2 - m_2^2)(m_2^2 - m_3^2)} \ln \left( \frac{1 + \Lambda^2/m_3^2}{1 + \Lambda^2/m_2^2} \right) + \frac{m_2^2}{(m_2^2 - m_3^2)(m_3^2 - m_1^2)} \ln \left( \frac{1 + \Lambda^2/m_3^2}{1 + \Lambda^2/m_1^2} \right),$$

$$G_1(m_1, m_2) = \int_0^1 dx \, (1 - x) \left[ \frac{1}{2} \left[ -\frac{m_3^2}{(m_1^2 - m_2^2)} x + m_1^2 \right] + \Lambda^2 \left[ -\frac{m_3^2}{(m_1^2 - m_2^2)} x + m_1^2 \right] + \Lambda^2 \right] - \frac{1}{2} \left[ -\frac{m_3^2}{(m_1^2 - m_2^2)} x + m_1^2 \right],$$

$$G_2(m_1, m_2) = \int_0^1 dx \, (1 - x) \left[ \ln \left( \frac{-\frac{m_3^2}{(m_1^2 - m_2^2)} x + m_1^2}{-\frac{m_3^2}{(m_1^2 - m_2^2)} x + m_1^2} \right) - \frac{3}{2} \right]$$

$$+ \int_0^1 dx \, (1 - x) \frac{4}{2} \left[ -\frac{m_3^2}{(m_1^2 - m_2^2)} x + m_1^2 \right] \Lambda^2 + \frac{3}{2} \left[ -\frac{m_3^2}{(m_1^2 - m_2^2)} x + m_1^2 \right]^2 + \frac{\Lambda}{2} \left[ -\frac{m_3^2}{(m_1^2 - m_2^2)} x + m_1^2 \right],$$

$$M_{A^2} = \frac{g_2}{\sqrt{2}} \sqrt{f_s^2 + \frac{\beta}{2} v^2} = \frac{4\cos^2 \vartheta}{\sin \vartheta} \left( \frac{M_W}{v} \right)^2 M_p.$$
It is worth mentioning that we do not consider the tree level contribution of the heavy vectors to the $\Delta T$ and $\Delta S$ parameters, since they are of the form $\sim v^2/M^2_\rho$, which are subleading compared to the loop contributions.

As a result, the experimental constraints on the $T$ and $S$ parameters \cite{73} impose an upper bound on our mixing parameter $\tan \theta \lesssim 0.47$, for heavy vector masses from 2 TeV up to 3 TeV.

\section{Production and Decays of the Heavy Vectors}

The current important period of LHC exploration of the Higgs properties and discovery of heavier particles may provide crucial steps to unravel the electroweak symmetry breaking mechanism. Consequently, we complement our work by studying the production and decay channels of the heavy vector resonances which are relevant for the LHC.

At a hadron collider like the LHC, the most important production channel for the heavy vector resonance is quark–anti-quark annihilation. In our construction, the coupling of heavy vectors to quarks goes through a term which has its origin in the mixing between the gauge fields $A^{(1)}_\mu$ and $A^{(2)}_\mu$. Consequently, the $\rho$ production amplitude is proportional to $\tan \theta$ which acts as a suppression factor. The influence of $\tan \theta$ can be seen in Figure 6 where we show the $\rho$ production cross section, computed with CalcHEP \cite{75}, for different values of $M_\rho$ and $\tan \theta$.

![Figure 6](image.png)

\textbf{Figure 6}. Heavy vector production cross section $\sigma(pp \to \rho) \text{ vs. } M_\rho$, for $\tan \theta = 0.2, 0.3$ and 0.4.

Additionally, we compute the two-body decay rates of the heavy vectors. These rates, up to corrections of order
\( m_{\rho_0}^2/M_{\rho_0}^2 \) and \( M_W^2/M_{\rho_0}^2 \) are:

\[
\Gamma (\rho^0 \rightarrow q\bar{q}) \simeq \frac{3g^2 \tan^2 \vartheta}{96\pi} M_{\rho},
\]

\[
\Gamma (\rho^+ \rightarrow u_d \bar{d}) = \Gamma (\rho^- \rightarrow u_d \bar{d}) \simeq \frac{3g^2 \tan^2 \vartheta}{96\pi} |V_{ij}|^2 M_{\rho},
\]

\[
\Gamma (\rho^\pm \rightarrow W^\pm h) = \Gamma (\rho^0 \rightarrow Zh) \simeq \frac{g^2 \cot^2 \vartheta}{96\pi} M_{\rho},
\]

\[
\Gamma (\rho^0 \rightarrow W^+W^-) = \Gamma (\rho^+ \rightarrow W^+Z) \simeq \frac{g^2 \cot^2 \vartheta}{96\pi} M_{\rho}.
\]  

(27)

Fig. 7 displays the branching ratios of the neutral (a) and charged (b) heavy vectors to a quark-antiquark pair and to a SM-like Higgs in association with a SM gauge boson, as a function of \( \tan \vartheta \). This angle controls the strength of the coupling of the heavy vector resonances with fermions. Clearly the largest decay rates of the heavy vectors are into a pair of SM Gauge bosons as well as into a SM-like Higgs and SM gauge boson, for all values of \( \tan \vartheta \). The decays into quark-antiquark pairs are much smaller in the relevant region of parameter space. This is a direct consequence of the gauge structure of the model and the representations of the fermions and the Higgs doublet under the full gauge symmetry group.

VI. BOUNDS FROM LHC SEARCHES

The ATLAS and CMS collaborations have performed several searches for heavy resonances decaying into different final states \[76,80\]. These searches are based on upper limits in the resonant cross section for different heavy vector particles. We use those limits to set restrictions on the model parameter space thus complementing the diphoton and the electroweak precision test constraints described above. As stated at the end of Sections III and IV, the allowed mixing parameter \( \tan \vartheta \) is restricted to the range \( 0.15 \lesssim \tan \vartheta \lesssim 0.47 \). In what follows, we will use as benchmark points the values \( \tan \vartheta = 0.15, 0.20, 0.30 \) and 0.47.

![Figure 8](image-url)

**Figure 8.** Left: predicted \( pp \rightarrow \rho^0 \rightarrow l^+l^- \) resonant production at \( \sqrt{s} = 13 \) TeV, for the combined channel \( ee \) and \( \mu\mu \), as a function of \( M_{\rho_0} \), for different values of \( \tan \vartheta \); the solid line is the 95% C.L. upper limit obtained by ATLAS \[76\]. Right: *idem*, but for the \( pp \rightarrow \rho^0 \rightarrow \tau^+\tau^- \) channel; the solid line is the 95% C.L. upper limit obtained by CMS \[77\].

We now focus on the LHC upper limits to constrain the model parameter space using the final states \( l^+l^- \), \( l\nu \) \((l = e, \mu)\), \( \tau^+\tau^- \), \( jj \), \( tt \), \( WZ \), \( WW \), \( WH \) and \( ZH \), assumed to be produced through a resonant \( \rho^0 \) or \( \rho^\pm \) decay. For example, the observation of the combined dilepton modes \( e^+e^- \) and \( \mu^+\mu^- \) \[76\] provides a bound for a neutral resonance, which we identify here with the neutral state \( \rho^0 \). Fig. 8 (left) shows the cross section prediction for \( \rho^0 \) decaying into dileptons \((l = e, \mu)\), together with the upper limit obtained by ATLAS, thus setting restrictions on the \( \tan \vartheta \) and \( M_{\rho_0} \) parameter space. The CMS upper bounds in the \( \tau^+\tau^- \) final state \[77\] are less restrictive than those of \( e^+e^- \) and \( \mu^+\mu^- \), thus providing no further constraints as shown in Fig. 8 (right). The experimental bound in the \( l\nu \) final state, with \( l = e, \mu \), is as stringent as in the \( l^+l^- \) channel (see Fig. 8). In contrast to dileptons, the \( tt \) \[79\] and dijet \[80\] experimental upper bounds do not restrict our parameter space, as can be seen in Fig. 10.
Figure 9. Predicted $pp \rightarrow \rho^{\pm} \rightarrow l\nu$ resonant production at $\sqrt{s} = 13$ TeV, for $l = e, \mu$, as a function of $M_\rho$, for different values of $\tan \vartheta$; the solid line is the 95% C.L. upper limit obtained by ATLAS \cite{76}.

Figure 10. Left: predicted $pp \rightarrow \rho^{0,\pm} \rightarrow jj$ resonant production at $\sqrt{s} = 13$ TeV as a function of $M_\rho$, for different values of $\tan \vartheta$; the solid line is the 95% C.L. upper limit obtained by ATLAS \cite{80}. Right: idem, but for $pp \rightarrow \rho^0 \rightarrow t\bar{t}$; the solid line is the observed 95% C.L. upper limit obtained by ATLAS \cite{79}.

Now, the associated $hZ$ and $hW^{\pm}$ production estimates and upper bounds \cite{81} are shown in Fig. 11. Finally, the ATLAS constraints from $ZW$ and $W^+W^-$ production are shown in Fig. 12, where again we contrast the experimental upper bound \cite{82} with the resonant production of $W^+W^-$ (left) and $ZW$ (right).

Combining all these restrictions in the $\tan \vartheta$-$M_\rho$ plane, we arrive at Fig. 13 where the allowed region of the parameter space is shown. Here we include also the $h \rightarrow \gamma\gamma$ (diphoton) constraint—which provides the lower bounds on $\tan \vartheta$, and the upper bound from electroweak precision tests (EWPT)—which turns out to be less restrictive than the upper bounds from the dilepton and diboson channels, as shown in the figure.

VII. CONCLUSIONS

We studied a framework of strongly interacting dynamics where the Higgs (a scalar doublet), and also a heavy vector triplet, appear as composite fields below a scale $\Lambda \sim 4\pi v \sim 3$ TeV. Without addressing details of the strong dynamics, we focus on the effective theory below the scale $\Lambda$, assumed to be the Standard Model, with its $SU(2)_L \times U(1)_Y$ gauge group, with the addition of a $SU(2)_L$ triplet of heavy vectors. The inclusion of the composite fields in the effective Lagrangian, i.e. the Higgs and the heavy vectors, is done by considering the vectors as gauge fields of a hidden local $SU(2)_2$ symmetry, and the Higgs as a doublet under this same symmetry. On the other hand, the SM gauge group at this stage is a $SU(2)_1 \times U(1)$. The SM fermions transform only under the latter group. By the mechanism of
hidden local symmetry, the $SU(2)_1 \times SU(2)_2$ breaks down to the diagonal $SU(2)$ subgroup, which will be effectively the $SU(2)_L$ of the SM. This spontaneous breakdown is formulated in terms of a non-linear sigma model, where the “would-be Goldstone” bosons are absorbed into the massive vector triplets. In this process, the Higgs $SU(2)_2$ doublet, which originally interacts with the composite vectors only, now acquires interactions with the SM fields. In this way, the composite Higgs maintains a rather strong interaction with the composite vector triplets, and a weaker interaction with the SM fields.

We put to test the resulting spectrum and interactions, in view of the existing experimental data: we determined the constraints arising from the measured Higgs diphoton decay rate, electroweak precision tests and the searches of heavy resonances at the LHC in the final states $l^+l^-$ and $l\nu_l$ ($l = e, \mu$), $\tau^+\tau^-$, $jj$, $t\bar{t}$, $WZ$, $WW$, $WH$ and $ZH$.

As a consequence of these constraints, we find that heavy vector masses in the range 2.1 - 3.0 TeV are consistent with the data, together with a mixing of the heavy vectors with the SM gauge bosons in the range $\tan \vartheta \sim 0.1 - 0.3$. These values are also consistent with the assumption that the Higgs couples weakly to the Standard sector and strongly to the heavy vector resonances. In other words, the current experimental data still allows for a Higgs boson that is strongly coupled to a composite sector, here assumed as triplet of vector resonances.
Figure 13. Allowed and excluded regions in the model parameter space after including all the constraints, i.e. the $h \rightarrow \gamma\gamma$ (diphoton) constraint, the electroweak precision tests (EWPT) and the bounds from the LHC searches in the channels $l^+l^-$, $l\nu_l$ ($l = e, \mu, \tau$, $jj$, $t\bar{t}$, $ZH$, $WH$, $WW$ and $WZ$. The allowed region which is consistent with all the constraints is shown in white, while the grey regions are excluded.

Acknowledgements

We would like to thank F. Rojas, M. Schmaltz and J. Urbina for useful discussions. This work was supported in part by Conicyt (Chile) grants ACT-146 and PIA/Basal FB0821, and by Fondecyt (Chile) grants No. 1130617, 1170171, 1120346, 1160423 and 1170803. B.D was partially supported by Conicyt Becas-Chile and PIIC/DGIP.

[1] G. Aad et al. [ATLAS Collaboration], Phys. Lett. B 716, 1 (2012) doi:10.1016/j.physletb.2012.08.020 [arXiv:1207.7214 [hep-ex]].
[2] S. Chatrchyan et al. [CMS Collaboration], Phys. Lett. B 716, 30 (2012) doi:10.1016/j.physletb.2012.08.021 [arXiv:1207.7235 [hep-ex]].
[3] K. A. Olive et al. [Particle Data Group], Chin. Phys. C 38, 090001 (2014). doi:10.1088/1674-1137/38/9/090001
[4] R. Contino, Nuovo Cim. 32, 11 (2009) doi:10.1393/ncc/i2009-10427-3 [arXiv:0908.3578 [hep-ph]].
[5] C. Grojean, PoS EPS -HEP2009, 008 (2009) [arXiv:0910.4976 [hep-ph]].
[6] R. Contino, arXiv:1005.4269 [hep-ph].
[7] G. Panico and A. Wulzer, Lect. Notes Phys. 913, pp.1 (2016) doi:10.1007/978-3-319-22617-0 [arXiv:1506.01961 [hep-ph]].
[8] A. Arbey, G. Cacciapaglia, H. Cai, A. Deandrea, S. Le Corre and F. Sannino, Phys. Rev. D 95, no. 1, 015028 (2017) doi:10.1103/PhysRevD.95.015028 [arXiv:1502.04718 [hep-ph]].
[9] K. Agashe, R. Contino and A. Pomarol, Nucl. Phys. B 719, 165 (2005) doi:10.1016/j.nuclphysb.2005.04.035 [hep-ph/0412089].
[10] B. Gripaios, A. Pomarol, F. Riva and J. Serra, JHEP 0904, 070 (2009) doi:10.1088/1126-6708/2009/04/070 [arXiv:0902.1483 [hep-ph]].
[11] R. Barbieri, B. Bellazzini, V. S. Rychkov and A. Varagnolo, Phys. Rev. D 76, 115008 (2007) doi:10.1103/PhysRevD.76.115008 [arXiv:0706.0432 [hep-ph]].
[12] C. Csaki, A. Falkowski and A. Weiler, JHEP 0809, 008 (2008) doi:10.1088/1126-6708/2008/09/008 [arXiv:0804.1954 [hep-ph]].
[13] R. Contino, D. Marzocca, D. Pappadopulo and R. Rattazzi, JHEP 1110, 081 (2011) doi:10.1007/JHEP10(2011)081 [arXiv:1109.1570 [hep-ph]].
[14] J. Mrazek, A. Pomarol, R. Rattazzi, M. Redi, J. Serra and A. Wulzer, Nucl. Phys. B 853, 1 (2011) doi:10.1016/j.nuclphysb.2011.07.008 [arXiv:1105.5403 [hep-ph]].
[15] A. Pomarol and F. Riva, JHEP 1208, 135 (2012) doi:10.1007/JHEP08(2012)135 [arXiv:1205.6434 [hep-ph]].
[57] M. Chala, G. Durieux, C. Grojean, L. de Lima and O. Matsedonskyi, JHEP **1706**, 088 (2017) doi:10.1007/JHEP06(2017)088 [arXiv:1703.10624 [hep-ph]].

[58] D. D. Dietrich, F. Sannino and K. Tuominen, Phys. Rev. D **72**, 055001 (2005) doi:10.1103/PhysRevD.72.055001 [hep-ph/0505059].

[59] M. A. Shifman, A. I. Vainshtein, M. B. Voloshin and V. I. Zakharov, Sov. J. Nucl. Phys. **30**, 711 (1979) [Yad. Fiz. **30**, 1368 (1979)].

[60] M. B. Gavela, G. Girardi, C. Malleville and P. Sorba, Nucl. Phys. B **193**, 257 (1981) doi:10.1016/0550-3213(81)90529-0.

[61] P. Kalyniak, R. Bates and J. N. Ng, Phys. Rev. D **33**, 755 (1986) doi:10.1103/PhysRevD.33.755.

[62] J. F. Gunion, H. E. Haber, G. L. Kane and S. Dawson, Front. Phys. **80**, 1 (2000).

[63] M. Spira, Fortsch. Phys. **46**, 203 (1998) doi:10.1002/(SICI)1521-3978(199804)46:3<203::AID-PROP203>3.0.CO;2-4 [hep-ph/9705337].

[64] A. Djouadi, Phys. Rept. **459**, 1 (2008) doi:10.1016/j.physrep.2007.10.005 [hep-ph/0503173].

[65] W. J. Marciano, C. Zhang and S. Willenbrock, Phys. Rev. D **85**, 013002 (2012) doi:10.1103/PhysRevD.85.013002 [arXiv:1109.5304 [hep-ph]].

[66] L. Wang and X. F. Han, Phys. Rev. D **86**, 095007 (2012) doi:10.1103/PhysRevD.86.095007 [arXiv:1206.1673 [hep-ph]].

[67] M. D. Campos, A. E. Cárcamo Hernández, H. Päs and E. Schumacher, Phys. Rev. D **91**, no. 11, 116011 (2015) doi:10.1103/PhysRevD.91.116011 [arXiv:1408.1652 [hep-ph]].

[68] A. E. Cárcamo Hernández, I. de Medeiros Varzielas and E. Schumacher, Phys. Rev. D **93**, no. 1, 016003 (2016) doi:10.1103/PhysRevD.93.016003 [arXiv:1509.02083 [hep-ph]].

[69] V. Khachatryan et al. [CMS Collaboration], Eur. Phys. J. C **74**, no. 10, 3076 (2014) doi:10.1140/epjc/s10052-014-3076-z [arXiv:1407.0558 [hep-ex]].

[70] G. Aad et al. [ATLAS Collaboration], Phys. Rev. D **90**, no. 11, 112015 (2014) doi:10.1103/PhysRevD.90.112015 [arXiv:1408.7084 [hep-ex]].

[71] M. E. Peskin and T. Takeuchi, Phys. Rev. D **46**, 381 (1992) doi:10.1103/PhysRevD.46.381.

[72] G. Altarelli and R. Barbieri, Phys. Lett. B **253**, 161 (1991) doi:10.1016/0370-2693(91)91378-9.

[73] R. Barbieri, A. Pomarol, R. Rattazzi and A. Strumia, Nucl. Phys. B **703**, 127 (2004) doi:10.1016/j.nuclphysb.2004.10.014 [hep-ph/0405040].

[74] M. Baak, M. Goebel, J. Haller, A. Hoecker, D. Ludwig, K. Moenig, M. Schott and J. Stelzer, Eur. Phys. J. C **72**, 2003 (2012) doi:10.1140/epjc/s10052-012-2003-4 [arXiv:1107.0975 [hep-ph]].

[75] A. Belyaev, N. D. Christensen and A. Pukhov, Comput. Phys. Commun. **184**, 1729 (2013) doi:10.1016/j.cpc.2013.01.014 [arXiv:1207.6082 [hep-ph]].

[76] The ATLAS collaboration [ATLAS Collaboration], ATLAS-CONF-2017-027.

[77] V. Khachatryan et al. [CMS Collaboration], JHEP **1702**, 048 (2017) doi:10.1007/JHEP02(2017)048 [arXiv:1611.06594 [hep-ex]].

[78] M. Aaboud et al. [ATLAS Collaboration], arXiv:1706.04786 [hep-ex].

[79] The ATLAS collaboration, ATLAS-CONF-2016-014.

[80] The ATLAS collaboration [ATLAS Collaboration], ATLAS-CONF-2016-069.

[81] The ATLAS collaboration [ATLAS Collaboration], ATLAS-CONF-2017-018.

[82] The ATLAS collaboration [ATLAS Collaboration], ATLAS-CONF-2016-062.