



# Spin-1 resonances

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**Abstract** Spin-1 resonances are naturally present in composite Higgs frameworks. We first review a model independent approach to parameterize a single additional heavy triplet and then we consider more realistic models arising in composite Higgs scenarios where a larger number of spin-1 resonances is expected. In these cases, finite width and interference effects can heavily affect the bounds extracted from the data.

## 1 Introduction

The presence of additional spin-1 particles, triplets or singlets of SU(2), is a common feature of several beyond the Standard Model (BSM) frameworks. They appear as Kaluza–Klein excitations of  $W$ ,  $Z$  and  $\gamma$  in flat or warped extra-dimension extensions of the SM, as new vector and axial-vector resonances in walking technicolor inspired effective Lagrangians, and are naturally present in composite Higgs scenarios (for references see, for example, the recent reviews [1,2]). They are also expected in weakly interacting BSM theories like the Little Higgs models or Grand Unified theories. In this chapter, we review some of the topics related to the theoretical approaches to describe the interactions of the new spin-1 resonances with the SM particles and with a composite Higgs. The new vectors could give clear signals in the Drell–Yan channel at the LHC, in the dilepton or in the diboson channels. From the LHC measurements one could then derive bounds on their masses and couplings. A very convenient tool to get these limits is a general model independent approach [3]. This allows, from a comparison with the data, to get bounds on the couplings to fermions and bosons of a new spin-1 triplet at fixed mass. This simplified approach however is not able to reproduce extended models containing more than one triplet. We consider then possible proposals describing two new triplets which can describe the low-lying resonances relevant at the LHC and outline the importance of the interference and finite width effects.

In Sect. 2 we review a model independent approach, to parameterize the presence of an additional heavy triplet of spin-1 resonances, that makes use of an effective Lagrangian description. Experimental data from CMS and ATLAS collaborations, using this parameterization, are used in Sect. 3 to get bounds on two partic-

ular models. In Sect. 4 we review a general method to build effective Lagrangians describing two new heavy triplets, replicas of  $W$  and  $Z$ , analyze two particular models and review the bounds on their parameter space using the LHC measurements. Finally in Sect. 5 we outline, by considering a composite Higgs model, the relevance of the interference and finite width effects in the analysis of the data to extract bounds on the model itself.

## 2 Heavy vector triplet simplified model and bounds from the LHC measurements

Following [3] (see also [4]) we can parameterize the presence of a new heavy vector triplet, interacting with SM gauge bosons and fermions by the following simple effective phenomenological Lagrangian containing operators up to dimension  $d = 4$

$$\begin{aligned} \mathcal{L}_V = & -\frac{1}{4}D_{[\mu}V_{\nu]}^a D^{[\mu}V^{\nu]a} + \frac{m_V^2}{2}V_{\mu}^a V^{\mu a} \\ & + i g_V c_H V_{\mu}^a H^{\dagger} \tau^a \overleftrightarrow{D}^{\mu} H + \frac{g^2}{g_V} c_F V_{\mu}^a J_F^{\mu a} \\ & + \frac{g_V}{2} c_{VVV} \epsilon_{abc} V_{\mu}^a V_{\nu}^b D^{[\mu}V^{\nu]c} \\ & + g_V^2 c_{VVHH} V_{\mu}^a V^{\mu a} H^{\dagger} H \\ & - \frac{g}{2} c_{VW} \epsilon_{abc} W^{\mu\nu a} V_{\mu}^b V_{\nu}^c. \end{aligned} \quad (1)$$

Here  $V_{\mu}^a$ ,  $a = 1, 2, 3$ , is the new heavy vector triplet, in the adjoint representation of SU(2)<sub>L</sub> and with vanishing hypercharge, describing the charged and the neutral heavy spin-1 particles:  $V_{\mu}^{\pm} = (V_{\mu}^1 \mp iV_{\mu}^2)/\sqrt{2}$  and  $V_{\mu}^0 = V_{\mu}^3$ . The covariant derivatives in Eq. (1) are defined by

$$D_{[\mu}V_{\nu]}^a = D_{\mu}V_{\nu}^a - D_{\nu}V_{\mu}^a, \quad D_{\mu}V_{\nu}^a = \partial_{\mu}V_{\nu}^a + g \epsilon^{abc} W_{\mu}^b V_{\nu}^c, \quad (2)$$

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where  $g$  denotes the  $SU(2)_L$  gauge coupling. The last two terms of the first row of Eq. (1) contain the interactions of the new vector bosons with the Higgs and the three Goldstones (encoded in  $H$ ) and therefore, by the Equivalence Theorem, with the longitudinal components of  $W$  and  $Z$  which are their relevant degrees of freedom at high energies, and with the SM fermions. Specifically:

$$i H^\dagger \tau^a \overleftrightarrow{D}^\mu H = i H^\dagger \tau^a D^\mu H - i D^\mu H^\dagger \tau^a H, \quad (3)$$

$$D^\mu H = \left( \partial^\mu + ig \frac{\sigma^a}{2} W^a - ig' \frac{\sigma^3}{2} Y^\mu \right) H, \quad (4)$$

with  $g'$  the weak hypercharge coupling, while  $J_F^{\mu a}$  are the SM left-handed fermionic currents

$$J_F^{\mu a} = \sum_f \bar{f}_L \gamma^\mu \frac{\sigma^a}{2} f_L. \quad (5)$$

The operators in the second line of Eq. (1) are not relevant for LHC because they do not contain vertices with SM gauge bosons.

Analysis at ATLAS and CMS on new heavy vector triplet models are performed by using the first line operators, in terms of  $m_V$  and two parameters  $g_H$  and  $g_F$  defined as

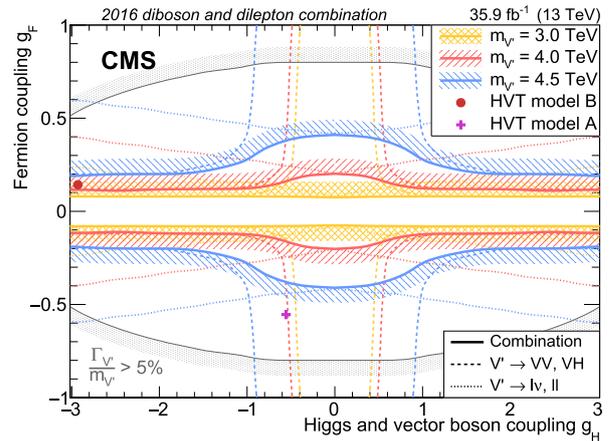
$$g_H = g_V c_H, \quad g_F = \frac{g^2 c_F}{g_V}. \quad (6)$$

In Fig. 1 the CMS bounds on  $(g_H, g_F)$ , obtained by using an integrated luminosity of  $35.6 \text{ fb}^{-1}$  at the LHC with a center of mass energy of 13 TeV [5] are shown. A combination of searches for a heavy triplet of resonances decaying into pairs of vector bosons, a vector boson and a Higgs boson, two Higgs bosons, or pairs of leptons, has been considered. The bounds are obtained for narrow resonances, the areas bounded by the thin gray contour lines correspond to regions where the resonance widths are predicted to be larger than the average experimental resolution (5%). The red dot and the violet cross correspond to model A and B that we review in the next subsection. A similar analysis has been performed by ATLAS [6].

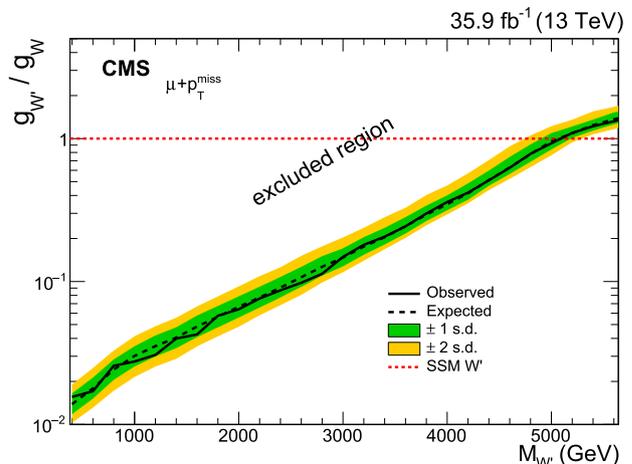
In Fig. 2 we present a further result from a CMS analysis corresponding of the same integrated luminosity corresponds to  $35.9 \text{ fb}^{-1}$ , where only the charged channel in Drell–Yan process of new resonances has been considered to obtain 95%CL limits on the coupling strength  $g_{W'}/g_W$  as a function of the  $W'$  mass [7]. Notice that for the Sequential Standard Model (SSM), where the coupling of the  $W'$  is set at the SM value, the bound on  $M_{W'}$  is around 5 TeV.

### 3 Two reference models

In this section we review two reference models that have been proposed in the past to describe additional triplets and have been used in the ATLAS and CMS analyses.



**Fig. 1** Exclusion limits on the couplings of heavy vector resonances to fermions  $g_F$  and SM vector bosons and the Higgs boson  $g_H$  for narrow resonances obtained from the statistical combination of all the channels (solid lines). The dilepton (dotted lines) and the diboson searches (dashed lines) constrain different regions. Three resonance mass hypotheses (3.0, 4.0, and 4.5 TeV) are considered. The hatched bands indicate the regions excluded. The analysis uses an integrated luminosity of  $35.9 \text{ fb}^{-1}$ . From [5]



**Fig. 2** Expected (dashed line) and observed (solid line) 95%CL limits on the coupling strength  $g_{W'}/g_W$  as functions of the  $W'$  mass, for the muon channel. The area above the limit line is excluded. The SSM  $W'$  couplings are shown as a dotted line. The analysis uses an integrated luminosity of  $35.9 \text{ fb}^{-1}$ . Similar bounds are extracted from the electron channel. From [7]

They are a model, (A), with an extended electroweak symmetry [8] and a strongly coupled composite model (B) [9].

#### 3.1 Model A

This model corresponds to an extension of the electroweak symmetry to  $SU(2)_1 \otimes SU(2)_2 \otimes U(1)_Y$  broken to  $SU(2) \otimes U(1)_Y$ . Fermions are supposed to transform

only under  $SU(2)_1 \otimes U(1)_Y$ . Two scalar fields are introduced: the SM Higgs  $H$ , a doublet of  $SU(2)_1$ , and an additional one,  $\Phi$ , belonging to the representation  $(\mathbf{2}, \mathbf{2})$  of  $SU(2)_1 \otimes SU(2)_2 \otimes U(1)_Y$ , the latter acquiring a vacuum expectation value:

$$\langle \Phi \rangle = \begin{pmatrix} f & 0 \\ 0 & f \end{pmatrix}. \tag{7}$$

This bidoublet breaks the  $SU(2)_1 \otimes SU(2)_2$  symmetry to the diagonal  $SU(2)$ . The gauge couplings  $g_V$  and  $g$  are identified as  $g_V = g_2$ ,  $1/g^2 = 1/g_1^2 + 1/g_2^2$ , with  $g_{1,2}$  the gauge couplings of  $SU(2)_{1,2}$ . Other details can be found in [8] and [3]. In Fig. 1 the violet cross represents this model. Only masses of the new resonances higher than 4.5–5 TeV are allowed.

### 3.2 Model B

This is the simplest model where the Higgs boson is realized as a pseudo Nambu–Goldstone from the breaking, at the scale  $f$ , of an extended symmetry,  $SO(5) \rightarrow SO(4)$ , the so-called Minimal Composite Higgs Model [9]. The construction of the Lagrangian is based on the Callan–Coleman–Wess–Zumino formalism. The Higgs is identified as a  $(2,2)$  under the unbroken  $SU(2)_L \otimes SU(2)_R \sim SO(4)$  and new vectors are introduced as representations of the unbroken  $SO(4)$ . In particular a new triplet field  $\rho_\mu$ ,  $(\mathbf{3}, \mathbf{1})$  of  $SU(2)_L \otimes SU(2)_R \sim SO(4)$  is present in the spectrum. The Lagrangian describing the Higgs gauge boson sector is given by

$$\begin{aligned} \mathcal{L}_\rho = & -\frac{1}{4\tilde{g}^2}(B_{\mu\nu})^2 - \frac{1}{4\tilde{g}^2}(W_{\mu\nu}^a)^2 + \frac{f^2}{4}d_\mu^i d^{\mu i} \\ & - \frac{1}{4g_\rho^2}(\rho_{\mu\nu}^a)^2 + \frac{m_\rho^2}{2g_\rho^2}(\rho_\mu^a - e_\mu^a)^2. \end{aligned} \tag{8}$$

The  $\rho$  field strength is given by  $\rho_{\mu\nu}^a = \partial_\mu \rho_\nu^a - \partial_\nu \rho_\mu^a - \epsilon^{abc} \rho_\mu^b \rho_\nu^c$ , while the  $d$  and  $(\rho - e)$  terms for the  $SO(5)/SO(4)$  coset, in the large  $f$  limit, are given by (see also Ref. [10]):

$$\begin{aligned} d_\mu^i d^{\mu i} = & \frac{4}{f^2}|D_\mu H|^2 + \frac{2}{3f^4}[(\partial_\mu |H|^2)^2 - 4|H|^2|D_\mu H|^2] \\ & + O\left(\frac{1}{f^6}\right), \end{aligned} \tag{9}$$

and

$$\begin{aligned} \rho_\mu^a - e_\mu^a = & \rho_\mu^a + W_\mu^a - \frac{i}{f^2}H^\dagger \tau^a \overleftrightarrow{D}_\mu H \\ & + \frac{i}{f^4}|H|^2 H^\dagger \tau^a \overleftrightarrow{D}_\mu H + O\left(\frac{1}{f^6}\right). \end{aligned} \tag{10}$$

The heavy vector triplet  $V_\mu^a$  of Eq. (1), is identified as

$$V_\mu^a \equiv \rho_\mu^a + W_\mu^a. \tag{11}$$

After the breaking of  $SU(2) \otimes U(1)_Y$ , the couplings of the physical Higgs to pairs of SM gauge bosons are rescaled by

$$\sqrt{1 - \xi}, \quad \text{with} \quad \xi = \frac{v^2}{f^2}, \quad v = 245 \text{ GeV}. \tag{12}$$

The current limit on  $\xi$ , using  $k_V = 1.05 \pm 0.04$  from ATLAS [32], is  $\xi \lesssim 0.06$  at 95%CL.

The coupling  $g_V$  is here identified with  $g_\rho$ , and it turns out [3] that

$$c_H \sim c_F \sim 1. \tag{13}$$

The red dot of Fig. 1 corresponds to  $c_H = -1$  and  $g_V = 3$ , for which resonances lighter than 4–4.5 TeV are excluded. Other vector resonances like  $\rho_R$  belonging to the representation  $(\mathbf{1}, \mathbf{3})$  of  $SU(2)_L \otimes SU(2)_R$  could be included in the model, see for instance [11].

## 4 Effective Lagrangians for new spin-1 triplets

A general procedure for building effective Lagrangians describing new vector resonances is the so-called *hidden local symmetry approach* [12–14]. We review here the main ingredients for two new triplets of vector resonances. For similar proposals in the framework of walking technicolor models see [15,16]. In these approaches the Higgs is described as an isosinglet scalar state.

We start considering the following group structure:  $G' = SU(2)_L \otimes SU(2)_R \otimes SU(2)'_L \otimes SU(2)'_R$  broken to the diagonal  $SU(2)_V$ . The nine Goldstone bosons resulting from the spontaneous breaking can be described by three independent  $SU(2)$  fields  $L(x)$ ,  $R(x)$  and  $M(x)$ , transforming under the extended symmetry group  $G'$  as follows:

$$L' = g_L L h_L, \quad R' = g_R R h_R, \quad M' = h_R^\dagger M h_L, \tag{14}$$

with  $g_{L,R} \in SU(2)_{L,R}$  and  $h_{L,R} \in SU(2)'_{L,R}$ .

In order to derive the most general Lagrangian up to second order in the derivatives, it is convenient to consider the following covariant quantities

$$\mathcal{V}_0^\mu = L^\dagger D^\mu L, \quad \mathcal{V}_1^\mu = M^\dagger D^\mu M, \quad \mathcal{V}_2^\mu = M^\dagger (R^\dagger D^\mu R) M, \tag{15}$$

with

$$\begin{aligned} D_\mu L = & \partial_\mu L - L \mathbf{L}_\mu, \quad D_\mu R = \partial_\mu R - R \mathbf{R}_\mu, \\ D_\mu M = & \partial_\mu M - M \mathbf{L}_\mu + \mathbf{R}_\mu M \end{aligned} \tag{16}$$

where  $\mathbf{L}_\mu = ig_L L_\mu^a \tau^a / 2$ ,  $\mathbf{R}_\mu = ig_R R_\mu^a \tau^a / 2$  are gauge fields of  $SU(2)'_L \otimes SU(2)'_R$ . Using Eq. (15) we can build

six independent invariants under the transformations of Eq. (14):

$$\text{Tr}\mathcal{V}_0^2, \text{Tr}\mathcal{V}_1^2, \text{Tr}\mathcal{V}_2^2, \text{Tr}(\mathcal{V}_0\mathcal{V}_1), \text{Tr}(\mathcal{V}_1\mathcal{V}_2), \text{Tr}(\mathcal{V}_0\mathcal{V}_2). \tag{17}$$

In conclusion, the most general Lagrangian for the vector boson interactions is a combination of these six invariants. The gauging of the Lagrangian with respect to the electroweak symmetry is obtained by substituting the covariant derivatives of Eq. (16) with

$$\begin{aligned} D_\mu L &\rightarrow D_\mu L + \tilde{\mathbf{W}}_\mu L, & D_\mu R &\rightarrow D_\mu R + \tilde{\mathbf{Y}}_\mu R, \\ D_\mu M &\rightarrow D_\mu M, \end{aligned} \tag{18}$$

where  $\tilde{\mathbf{W}}_\mu = i\tilde{g}\tilde{W}_\mu^a\tau^a/2$ ,  $\tilde{\mathbf{Y}}_\mu = i\tilde{g}'\tilde{Y}_\mu\tau^3/2$ , and the “tildes” indicate quantities which are not the SM ones due to mixing. The non-linear formulation of the SM is obtained in terms of the matrix  $U \equiv LM^\dagger R^\dagger$  and the scalar singlet  $h$ . The effective Higgs Lagrangian is then written in terms of the quantities given in Eq. (15) as

$$\begin{aligned} &\frac{v^2}{4} \left( 1 + 2a\frac{h}{v} + b\frac{h^2}{v^2} \right) \text{Tr}(D^\mu U)^\dagger D^\mu U + \dots \\ &= -\frac{v^2}{4} \left( 1 + 2a\frac{h}{v} + b\frac{h^2}{v^2} \right) \text{Tr}[(\mathcal{V}_0 - \mathcal{V}_1 - \mathcal{V}_2)^2] + \dots \end{aligned} \tag{19}$$

The SM couplings correspond to  $a = b = 1$  and dots stand for interaction terms between the scalar singlet  $h$  and the other invariants of Eq. (17). For example, by adding the two invariants  $\text{Tr}\mathcal{V}_0^2$  and  $\text{Tr}\mathcal{V}_2^2$ , one obtains a model with two spin-1 vector triplets, with  $\mathbf{L}_\mu$  and  $\mathbf{R}_\mu$  unmixed, when neglecting weak interactions.

Following the *hidden local symmetry* approach, here briefly reviewed, the SM fermions are coupled to the extra spin-1 resonances via mixing with the SM gauge bosons. In addition, direct couplings, for example, of the  $\psi_L$  fermions to  $\mathbf{L}_\mu$  are allowed [13]. In fact, for each,  $\psi_L$  we can construct a SU(2) doublet

$$\chi_L = L^\dagger \psi_L, \tag{20}$$

and consider the following invariant term

$$b_L \bar{\chi}_L i\gamma^\mu (\partial_\mu + i\mathbf{L}_\mu + \frac{i}{2}\tilde{g}'(B - L)Y_\mu)\chi_L \tag{21}$$

where  $b_L$  is a dimensionless parameter. In the unitary gauge  $L = I$  and, shifting  $\psi_L \rightarrow \psi_L/\sqrt{1 + b_L}$ , one gets a fermion-heavy triplet interaction term

$$-\frac{b_L}{1 + b_L} \bar{\psi}_L \gamma^\mu \mathbf{L}_\mu \psi_L \tag{22}$$

contributing, together with the mixing term, to  $g_F = g^2_{CF}/g_V$  in Eq. (1). The procedure can be generalized to several triplets [17].

### 4.1 Explicit models with two new spin-1 triplets

In the following we will concentrate on two different models describing two extra triplets of spin-1 resonances. We have seen that their interactions are described, in general, by six independent invariants. They are reduced to four by requiring the invariance under the discrete left-right transformation, denoted by  $P : L \leftrightarrow R, M \leftrightarrow M^\dagger$ . Namely, the most general  $G' \otimes P$  invariant Lagrangian is given by [14]

$$\mathcal{L}_G = -\frac{v^2}{4} [a_1 I_1 + a_2 I_2 + a_3 I_3 + a_4 I_4] + \mathcal{L}_{\text{kin}} \tag{23}$$

with

$$\begin{aligned} I_1 &= \text{Tr}[(\mathcal{V}_0 - \mathcal{V}_1 - \mathcal{V}_2)^2], & I_2 &= \text{Tr}[(\mathcal{V}_0 + \mathcal{V}_2)^2], \\ I_3 &= \text{Tr}[(\mathcal{V}_0 - \mathcal{V}_2)^2], & I_4 &= \text{Tr}[\mathcal{V}_1^2] \end{aligned} \tag{24}$$

and the kinetic terms (we take  $g_L = g_R = g''/\sqrt{2}$ ):

$$\mathcal{L}_{\text{kin}} = \frac{1}{g''^2} \text{Tr}[F_{\mu\nu}(\mathbf{L})]^2 + \frac{1}{g''^2} \text{Tr}[F_{\mu\nu}(\mathbf{R})]^2. \tag{25}$$

#### 4.1.1 4-site model with a composite Higgs

As said, new vector resonances appear in five-dimensional extensions of the SM as Kaluza–Klein (KK) excitations of the SM gauge bosons [18–20]. When deconstructed [21–24], these theories appear as gauge theories with extended SU(2) symmetries.

We review here a simple four-dimensional model [25–30], corresponding to a deconstruction including two additional copies of the SU(2) symmetry, namely  $G' = \text{SU}(2)_L \otimes \text{SU}(2)_R \otimes \text{SU}(2)_1 \otimes \text{SU}(2)_2$  with the particular choice:

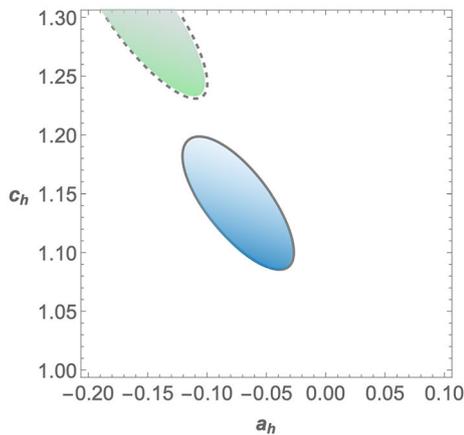
$$a_1 = 0, \quad a_2 = a_3 = \frac{2f_1^2}{v^2}, \quad a_4 = \frac{4f_2^2}{v^2}, \quad \frac{4}{v^2} = \frac{2}{f_1^2} + \frac{1}{f_2^2}. \tag{26}$$

Following the general construction illustrated above, we add to the Lagrangian the scalar–scalar and scalar–vector interactions as in Eq. (19) (with scalar we here refer both the the Higgs particle and the Goldstones):

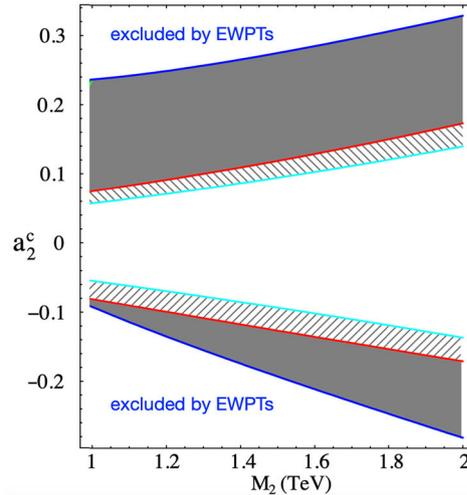
$$\begin{aligned} \mathcal{L}_{hG} &= \left( 2a_h \frac{h}{v} + b_h \frac{h^2}{v^2} \right) f_1^2 [\text{Tr}(D_\mu \Sigma_1)^\dagger D^\mu \Sigma_1 \\ &\quad + \text{Tr}(D_\mu \Sigma_3)^\dagger D^\mu \Sigma_3] \\ &\quad + \left( 2c_h \frac{h}{v} + d_h \frac{h^2}{v^2} \right) f_2^2 \text{Tr}(D_\mu \Sigma_2)^\dagger D^\mu \Sigma_2, \end{aligned} \tag{27}$$

with the identification of the chiral fields:  $\Sigma_1 = L, \Sigma_2 = M^\dagger, \Sigma_3 = R$ . The covariant derivatives are defined as in Eq. (18) with the additional gauge fields identified as:  $\mathbf{L}_\mu \rightarrow -\tilde{\mathbf{A}}_\mu^1, \mathbf{R}_\mu \rightarrow -\tilde{\mathbf{A}}_\mu^2, \tilde{\mathbf{A}}_\mu^i = ig_i \tilde{A}_\mu^i \tau^a/2$  and we assume  $g_1 = g_2$  for the couplings.

In the unitary gauge, this model predicts two new triplets of gauge bosons, which acquire mass through



**Fig. 3** Left: 95%CL bounds on the plane  $(a_h, c_h)$  for  $z = 0.95$  (continuous blue) and  $z = 0.9$  (dashed green) from ATLAS measurements [32] on  $k_V = 1.05 \pm 0.04, k_\gamma = 1 \pm 0.06$  based on  $79.8 \text{ fb}^{-1}$ . Right:  $5\sigma$ -discovery plot from D–Y pro-



duction at LHC with  $L=100 \text{ fb}^{-1}$  for  $z = 0.8$  in the plane  $(a_2^c, M_2)$ . Inside the grey (dashed) (central white) regions both  $W_{1,2}$  (only  $W_2$ ) (no resonance) can be detected. From [25]

the same non-linear symmetry breaking mechanism giving mass to the SM gauge bosons. For large  $g_1, W^\pm, Z$  mass eigenvalues are given by

$$\begin{aligned} M_{W^\pm}^2 &\simeq \frac{\tilde{g}^2 v^2}{4} \left[ 1 - \frac{\tilde{g}^2}{2g_1^2} (1 + z^4) \right], \\ M_Z^2 &\simeq \frac{(\tilde{g}^2 + \tilde{g}'^2) v^2}{4} \left[ 1 - \frac{\tilde{g}^2}{2g_1^2} z_\zeta \right], \end{aligned} \quad (28)$$

where

$$z = \frac{f_1}{\sqrt{f_1^2 + 2f_2^2}}, \quad z_\zeta = \frac{1}{2} \frac{z^4 + \cos^2 2\tilde{\theta}}{\cos^2 \tilde{\theta}}, \quad \tan \tilde{\theta} = \frac{\tilde{g}'}{\tilde{g}}. \quad (29)$$

The heavy vector triplets,  $\mathbf{A}^1 = (W_1^\pm, Z_1)$  and  $\mathbf{A}^2 = (W_2^\pm, Z_2)$ , have masses approximately given by  $M_1, M_2$  with

$$M_1 \simeq \frac{g_1 v}{\sqrt{2(1 - z^2)}}, \quad M_2 \simeq \frac{1}{z} M_1, \quad (30)$$

As said, direct couplings of the extra spin-1 resonances to the SM fermions are allowed by the symmetries. They are present also in the small mixing limit and can be introduced by following the construction described above.

From Eq. (27) we derive the couplings of the Higgs  $h$  to the charged gauge bosons:

$$\frac{2h}{v} [a_M M_W^2 W^+ W^- + a_h M_1^2 W_1^+ W_1^- + (a_h z^2 + c_h (1 - z^2) M_2^2 W_2^+ W_2^-)], \quad (31)$$

$$a = a_h (1 - z^2) + c_h z^2, \quad (32)$$

and similarly for the neutral sector. The SM limit is obtained for  $z \rightarrow 1$  (namely  $g_1, f_1 \rightarrow \infty$ ),  $c_h \rightarrow 1$  and  $a_h \rightarrow 0$ . In this limit  $a \rightarrow 1$ .

Bounds on  $hW^+W^-, hZZ, h\gamma\gamma$  couplings from recent analysis at LHC, can be translated in limits on  $a_h, c_h$  for different values of  $z$ . Assuming no deviation in the top Higgs couplings ( $c_t = 1$ ), we can compute [31]

$$\begin{aligned} \Gamma/\Gamma_{SM}(h \rightarrow \gamma\gamma) &\sim \left[ 1 + \frac{9}{8} [a_h (1 + z^2) + c_h (1 - z^2)] + \frac{9}{7} (a - 1) \right], \end{aligned} \quad (33)$$

and use ATLAS measurements [32] on  $k_V = 1.05 \pm 0.04, k_\gamma = 1 \pm 0.06$  based on  $79.8 \text{ fb}^{-1}$  of integrated luminosity to get the 95%CL bounds shown in Fig. 3 (Left).

Unitarity bounds on the 4-Site model extra spin-1 resonance masses are approximately 2–4 TeV depending on the  $z$  value. For  $z = 0.95$  unitarity bounds increases to approximately 4 TeV [29].

Bounds from  $S, T, U$  parameters on the charged couplings of the heavy triplets to SM fermions,  $a_1^c$  and  $a_2^c$ , and on the corresponding neutral couplings, can be found in [25–30]. For example, for  $z = 0.95, M_{W_1} = 2 \text{ TeV}$ , one obtains  $-0.35 \leq a_1^c \leq 0.4$ . An extended analysis of Drell–Yan production of these new vectors in the dilepton and diboson channel at LHC was also performed in [25–30]. The  $5\sigma$ -discovery contours for a foreseen LHC integrated luminosity of  $100 \text{ fb}^{-1}$  in the plane  $(a_2^c, M_2)$ , for  $z = 0.8$ , where  $a_2^c$  is the  $W_2$  charged coupling and  $M_2$  is approximately  $W_2$  mass, is plotted in Fig. 3 (Right). The upper and lower parts of the plot are excluded by the EWPTs. Inside the grey regions both  $W_{1,2}$  are visible; inside the dashed ones only  $W_2$

can be detected. Inside the central uncolored region no resonance is visible in the Drell-Yan channel.

However, in presence of several new triplets, as in this model, we cannot extract bounds on the couplings of the new resonances from Figs. 1 and 2 because as shown in [27,33,34] there are significant interference effects that cannot be neglected. In Sect. 5 we review these effects in a composite Higgs model.

#### 4.1.2 Linear degenerate BESS

A second interesting choice in the general parameterization given in Eq. (23) is  $a_1 = 1, a_4 = 0, a_2 = a_3$ . The Lagrangian then reads

$$\mathcal{L}_G = \frac{v^2}{4} \{ \text{Tr}(\partial_\mu U^\dagger \partial^\mu U) + 2 a_2 [\text{Tr}(D_\mu L^\dagger D^\mu L) + \text{Tr}(D_\mu R^\dagger D^\mu R)] \}. \tag{34}$$

In this model, before turning on electro-weak interactions, the new triplets  $\mathbf{L}_\mu$  and  $\mathbf{R}_\mu$  are degenerate in mass (from here the name *degenerate Breaking the Electro-Weak Symmetry Strongly (BESS) Model*). Notice that each of the three terms in the above expression is invariant under an independent  $SU(2) \otimes SU(2)$  group

$$U' = \omega_L U \omega_R^\dagger, \quad L' = g_L L h_L, \quad R' = g_R R h_R, \tag{35}$$

so that the symmetry of the Lagrangian is enlarged to  $G_{\max} = [SU(2) \otimes SU(2)]^3$ . As a consequence of this enlarged symmetry or equivalently of the degeneracy of  $\mathbf{L}_\mu$  and  $\mathbf{R}_\mu$  fields (or the combination corresponding to the vector and axial-vector fields) the correction to the  $S$  parameter at leading order is zero [35]. The possibility of an enhanced symmetry in near conformal theories and the emergence of parity doublet was suggested also in [15] (see also [16]) in the framework of walking technicolor.

The degenerate BESS model has also a linear (renormalizable) formulation [36,37]. It is based on a gauge group  $G = SU(2)_L \otimes U(1) \otimes SU(2)'_L \otimes SU(2)'_R$  and has a scalar sector consisting of scalar fields  $\tilde{L} \in (2, 0, 2, 0), \tilde{U} \in (2, 2, 0, 0), \tilde{R} \in (0, 2, 0, 2)$  of the group  $G$  (it is a generalized version of Model A in Sect. 3.1). The scalar fields break the gauge symmetries through the following chain

$$\begin{array}{ccc} SU(2)_L \otimes U(1) \otimes SU(2)'_L \otimes SU(2)'_R & & \\ \downarrow u & & \\ SU(2)_{\text{weak}} \otimes U(1)_Y & & (36) \\ \downarrow v & & \\ U(1)_{\text{em}} & & \end{array}$$

The two breakings are induced by the expectation values  $\langle \tilde{L} \rangle = \langle \tilde{R} \rangle = u$ , and  $\langle \tilde{U} \rangle = v$ , respectively and we assume  $u \gg v$ . The first two expectation values make the breaking  $SU(2)_L \otimes SU(2)'_L \rightarrow SU(2)_{\text{weak}}$  and  $U(1) \otimes SU(2)'_R \rightarrow U(1)_Y$ , whereas the second breaks

in the standard way  $SU(2)_{\text{weak}} \otimes U(1)_Y \rightarrow U(1)_{\text{em}}$ . The parameter  $a_2$  of Eq. (34) is related to  $u$  via  $a_2 = u^2/(2v^2)$ .

Standard fermions are supposed to couple only to  $\tilde{U}$ . The model, in the limit of large  $u$ , satisfies decoupling giving back the SM with a light Higgs and suppressed contributions to the parameters  $S, T, U$ .

We parameterize the scalar fields as  $\tilde{L} = \tilde{\rho}_L L, \tilde{R} = \tilde{\rho}_R R, \tilde{U} = \tilde{\rho}_U U$ , with  $L^\dagger L = I, R^\dagger R = I$  and  $U^\dagger U = I$ . The Lagrangian for their kinetic terms and interactions with the gauge bosons is given by  $(\mathbf{L}_\mu \rightarrow \tilde{\mathbf{A}}_\mu^1, \mathbf{R}_\mu \rightarrow -\tilde{\mathbf{A}}_\mu^2)$

$$\mathcal{L}_{\rho G} = \frac{1}{4} \left[ \text{Tr}(D_\mu \tilde{U})^\dagger (D^\mu \tilde{U}) + \text{Tr}(D_\mu \tilde{L})^\dagger (D^\mu \tilde{L}) + \text{Tr}(D_\mu \tilde{R})^\dagger (D^\mu \tilde{R}) \right], \tag{37}$$

$$\begin{aligned} D\tilde{L} &= \partial\tilde{L} + \tilde{\mathbf{W}}_\mu \tilde{L} - \tilde{L} \tilde{\mathbf{A}}_\mu^1 \\ D\tilde{R} &= \partial\tilde{R} + \tilde{\mathbf{Y}}_\mu \tilde{R} - \tilde{R} \tilde{\mathbf{A}}_\mu^2 \\ D\tilde{U} &= \partial\tilde{U} + \tilde{\mathbf{W}}_\mu \tilde{U} - \tilde{U} \tilde{\mathbf{Y}}_\mu. \end{aligned} \tag{38}$$

The scalar potential is expressed in terms of three Higgs fields:

$$V = 2\mu^2(\tilde{\rho}_L^2 + \tilde{\rho}_R^2) + \lambda(\tilde{\rho}_L^4 + \tilde{\rho}_R^4) + 2m^2\tilde{\rho}_U^2 + h\tilde{\rho}_U^4 + 2f_3\tilde{\rho}_L^2\tilde{\rho}_R^2 + 2f\tilde{\rho}_U^2(\tilde{\rho}_L^2 + \tilde{\rho}_R^2). \tag{39}$$

After shifting the fields by their VEV's  $u, v \neq 0$ , the squared mass eigenvalues of the Higgses are:

$$\begin{aligned} m_{\rho_U}^2 &\simeq 8 \left( h - 2 \frac{f^2}{\lambda + f_3} \right) v^2, \\ m_{\rho_L}^2 &\simeq 8(\lambda - f_3)u^2, \\ m_{\rho_R}^2 &\simeq 8(\lambda + f_3)u^2. \end{aligned} \tag{40}$$

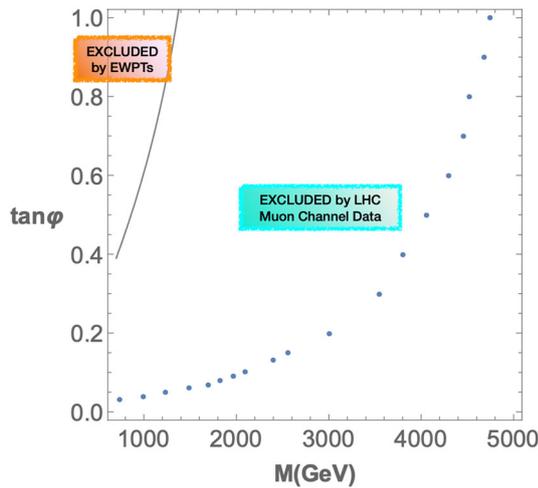
Concerning the gauge sector, let's call the mass eigenstates of the new spin-1 resonances  $\mathbf{A}^i = (W_i^\pm, Z_i)$ ,  $i = 1, 2$ . The fields  $W_2^\pm$  are unmixed and their squared mass is given by

$$M_{W_2}^2 = \frac{1}{4} g_1^2 u^2 \equiv M^2, \tag{41}$$

where  $g_1$  is the common coupling constant of the extra gauge symmetry. The other two charged eigenvalues, in the limit of small  $r = g^2 v^2/(4M^2)$ , are

$$M_W^2 \simeq \frac{v^2}{4} g^2 (1 - r s_\varphi^2), \quad M_{W_1}^2 \simeq \frac{M^2}{c_\varphi^2}, \tag{42}$$

with  $\varphi$  defined by the relation  $g = g_1 s_\varphi$ , with  $1/g^2 = 1/\tilde{g}^2 + 1/g_1^2$ . Notice that the SM gauge boson masses receive corrections  $O(r)$  due to mixing.



**Fig. 4** 95%CL bounds on the plane  $(M, \tan \varphi)$  from electro-weak precision measurements (continuous line) compared with bounds (dotted) from the muon channel at LHC extracted from the analysis of Fig. 2

In the charged sector, at the first order in  $r$ , the couplings are given by

$$\mathcal{L}_{\text{fermions}}^{\text{charged}} = -(a_W W_\mu^- + a_1^c W_{1\mu}^-) J_L^{\mu-} + h.c., \quad (43)$$

with

$$a_W = \frac{g}{\sqrt{2}}(1 - s_\varphi^2 r), \quad a_1^c = -\frac{g}{\sqrt{2}}(1 + c_\varphi^2 r) \tan \varphi, \quad (44)$$

and  $J_L^\pm = \bar{\psi}_L \gamma^\mu \tau^\pm \psi_L$ . There is no coupling of  $W_2^\pm$  to fermions, because they do not mix with  $W^\pm$ . The fermionic couplings of the neutral gauge boson sector are given in [37]. The heavy gauge bosons are coupled to fermions only through mixing with the SM ones, namely these couplings vanish for  $s_\varphi \rightarrow 0$ . Nevertheless one could consider also direct couplings to fermions as shown in the derivation of the general effective Lagrangian.

Since in this set-up, only the charged gauge bosons  $W_1^\pm$  are coupled to fermions, there is no the interference effect between  $W_1^\pm$  and  $W_2^\pm$  and we can use the analysis presented in Fig. 2 for a single heavy triplet resonances to extract bounds on the linear degenerate BESS model. Results on the plane  $(M, \tan \varphi)$  are shown in Fig. 4.

In the neutral sector both  $Z_1$  and  $Z_2$  are coupled to fermions and a dedicated analysis is necessary to get bounds from LHC results. Couplings of the SM-like Higgs  $\rho_U$  to SM fermions are rescaled by  $c_\alpha$ , where  $\alpha$  is the mixing angle in the scalar sector, while the couplings of  $\rho_U$  to SM gauge bosons are given by

$$\rho_U \left[ \frac{v}{2} g^2 c_\alpha (1 - s_\varphi^4 x^2) W^+ W^- + \frac{v}{4 c_\theta^2} g^2 c_\alpha \left[ 1 - 2x^2 \frac{s_\varphi^4}{c_\theta^4} (1 - 2c_\theta^2 s_\theta^2) \right] Z^2 \right]. \quad (45)$$

Due to the stringent limits on the plane  $(M, \tan \varphi)$ , current bounds from LHC on the parameter  $k_V$  [32] can be directly translated on a bound on  $\alpha$ , namely:  $0.97 \leq c_\alpha \leq 1$  at  $2\sigma$  level.

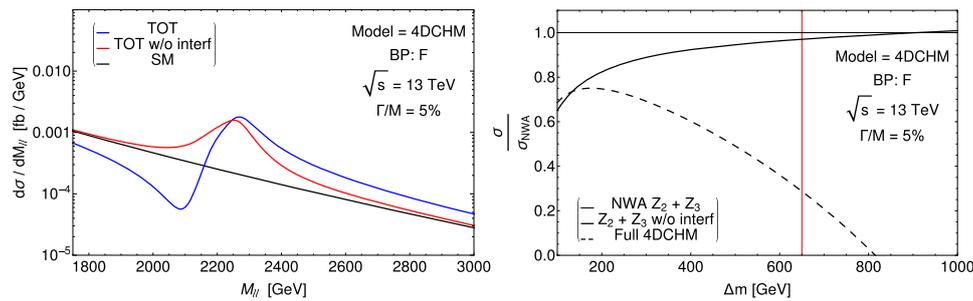
### 5 Drell–Yan production and interference effects in models with multi-triplet structure

Drell–Yan production in the dilepton channels is a very good tool for searching new heavy triplets as we have already discussed in Sect. 2. A class of BSM scenarios, like composite Higgs models and the SM model in five dimensions, predict several new spin-1 resonances. In these cases the interpretation of the experimental results and related bounds on the parameter space of the models is complicated by the presence of finite width and interference effects. We briefly review in this Section some of such effects for the D–Y process of the neutral channels (new  $Z$ 's) in composite Higgs models [34]. The four-dimensional composite Higgs model (4DCHM), that we are going to review, is the one proposed in [38], based on two non-linear  $\sigma$  models, one for  $\text{SO}(5)/\text{SO}(4)$  and the second for  $\text{SO}(5)_L \otimes \text{SO}(5)_R / \text{SO}(5)_{L+R}$ . After the breaking one is left with 4 neutral and 6 charged spin one heavy particles, or two triplets  $(\mathbf{3}, \mathbf{1}) + (\mathbf{1}, \mathbf{3})$ , almost degenerate in mass, and, in addition, two neutral and one charged coset resonances.

The extra neutral resonances, that couple to fermions, are denoted by  $Z_2, Z_3, Z_5$  with squared masses, at order  $O(\xi)$ , given by

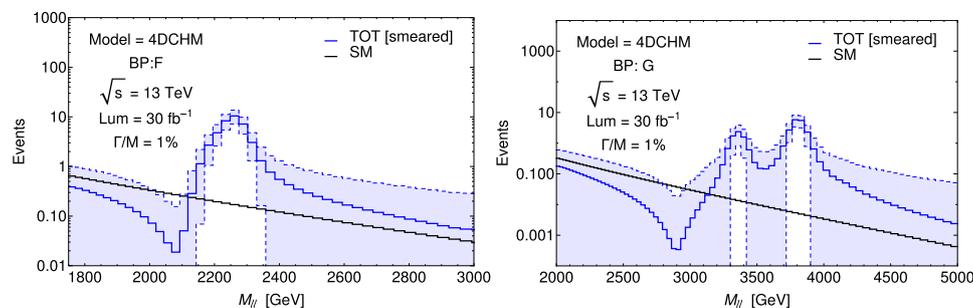
$$\begin{aligned} M_{Z_2}^2 &\simeq \frac{m_\rho^2}{c_\psi^2} \left( 1 - \frac{s_\psi^2 c_\theta^4}{4c_{2\psi}} \xi \right), \\ M_{Z_3}^2 &\simeq \frac{m_\rho^2}{c_\theta^2} \left( 1 - \frac{s_\theta^2 c_\theta^4}{4c_{2\theta}} \xi \right) \\ M_{Z_5}^2 &\simeq 2m_\rho^2 \left[ 1 + \frac{1}{16} \left( \frac{1}{c_{2\theta}} + \frac{1}{2c_{2\psi}} \right) \xi \right], \end{aligned} \quad (46)$$

with  $\xi = v^2/f^2$ ,  $\tan \theta = \tilde{g}/g_\rho$  and  $\tan \psi = \sqrt{2}\tilde{g}'/g_\rho$  where  $g_\rho$  is the gauge coupling of the  $\text{SO}(5)$  group,  $f$  is the scale of the spontaneous strong symmetry breaking and  $m_\rho = fg_\rho$ . The  $Z_5$  resonance, having a mass  $\sim \sqrt{2}$  higher than  $Z_2$  and  $Z_3$ , gives a negligible contribution to the Drell–Yan process at LHC RunII and it has been neglected in the analysis. The change in the differential D–Y cross section in the dilepton invariant mass produced by the interference is shown in Fig. 5 for the benchmark F, corresponding to  $M_{Z_2} = 2192$  GeV and  $M_{Z_3} = 2258$  GeV. The interference effects produce the dip before the resonant peak, and spoils the analysis performed within the narrow width approximation (NWA) or in the finite width approach (FWA). The same conclusion can be derived by plotting the ratio of the cross section over the NWA as a function



**Fig. 5** (Left) Differential cross section in the dilepton invariant mass for the benchmark point F within the 4DCHM, i.e. double-resonant  $Z_{2,3}$  scenario. (Right) Ratio of the full signal cross section for the  $Z_{2,3}$  bosons corresponding to benchmark F within the 4DCHM scenario (dashed line) and the two resonances FWA (solid line) over the NWA

result as a function of the symmetric integration interval around the peak. The vertical red line represents the CMS adopted optimal cut,  $\delta m = |M_{ll} - M_{Z'}| \leq 0.05 E_{LHC}$ , which keeps the interference and FW effects below 10% in the case of narrow single  $Z'$  models [33]. From [34]



**Fig. 6** Left(Right): Differential cross section in the dilepton invariant mass for the BP F(G) after the smearing due to finite detector resolution. The width of the Gaussian is assumed 26 GeV (38 GeV). The statistical error, visualized as the blue bands, is evaluated for  $30 \text{ fb}^{-1}$ . From [34]

**Table 1** 4DCHM parameter space points associated to the benchmarks F, G mentioned in the text. From [34]

Benchmark	$f$ (GeV)	$g_\rho$	$M_{Z_2}$ (GeV)	$M_{Z_3}$ (GeV)	$M_{Z_5}$ (GeV)
F	1200	1.75	2192	2258	2972
G	2900	1.00	3356	3806	4107

of  $\delta m = M_{ll} - M_{Z'}$ . In the NWA the signal rate is typically estimated modelling the signal as a Breit–Wigner function convoluted with a Gaussian function, which is used to describe the dilepton mass resolution. This example shows that one should avoid to use NWA or FWA within the 4DCHM or similar models like the 4-Site model of Sect. 4.1 to extract bounds on  $Z'$  couplings and masses.

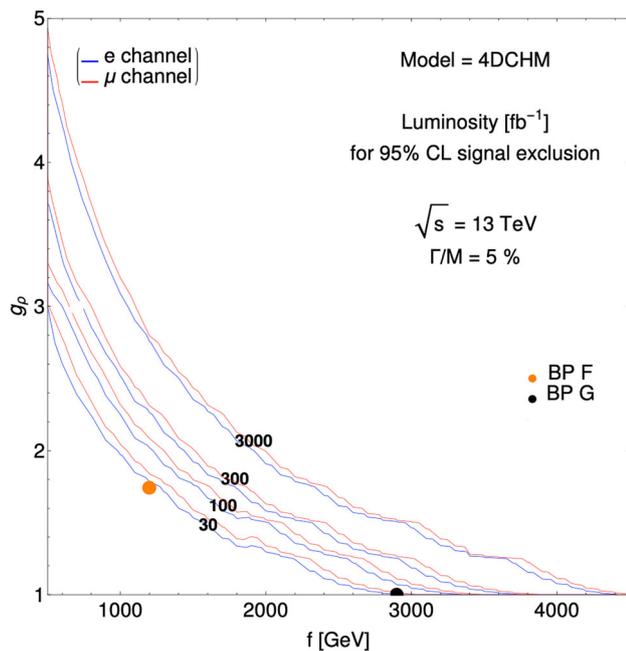
The two resonances  $Z_2$  and  $Z_3$  are, in some regions of the parameter space, almost degenerate. Therefore they could appear as a single bump in the dilepton invariant mass after convoluting the signal with a Gaussian distribution to simulate detector resolution. This is shown in Fig. 6 (Left) for the benchmark F of Table 1. In Fig. 6 (Right) the two resonances corresponding to the benchmark G of Table 1 are separated enough so that the two peak structure is not washed out after the inclusion of the detector smearing.

In Fig. 7 we show the exclusion limits at LHC RUN II for different values of integrated luminosity ( $30\text{--}3000 \text{ fb}^{-1}$ ) in the parameter space ( $f, g_\rho$ ).

The 4DCHM can be considered as a concrete realization of the Model B, described in Sect. (3.2) with the feature that the two almost degenerate triplets of spin-1 resonances corresponding to the SO(4) symmetry (broken by the EW interactions) are both phenomenologically relevant. Their interference cannot be neglected as well as their widths which is naturally non-narrow. A direct comparison of the results shown in Fig. 7 with the results shown in Fig. 1 is not possible, because the analysis of CMS refers to a single heavy vector triplet and takes into account, in addition to the D–Y processes, diboson and H-boson productions.

## 6 Conclusions

The presence of new spin-1 particles, heavier than  $W$  and  $Z$ , is a common feature of several proposals of BSM physics, like for example extensions of the SM in flat or warped extra dimensions, composite Higgs or tech-



**Fig. 7** 95%CL exclusion limits from D–Y production in dilepton channel at the LHC RUNII with 13 TeV for different values of integrated luminosity. The blue (red) contour refers to the electron (muon) channel. From [34]

nicolor models. In this note we have first reviewed a general model independent approach based on a simplified effective Lagrangian describing the interactions of a new heavy triplet with the SM particles. This parameterization has been used, using CMS and ATLAS results, to get bounds on the parameters of the effective Lagrangian. We have then reviewed a general approach to built effective Lagrangians describing the interactions in presence of two new triplets of heavy resonances and considered two particular models and their present limits. Finally we have discussed, in the framework of a composite Higgs model the relevance of the interference and finite width effects, in the Drell–Yan differential cross section in the dilepton neutral channel. The interference effect, coming from the contributions of the neutral spin-1 resonances exchanged in the s-channel, generate a dip before the resonant peak, that spoils the simple analysis performed within the narrow width approximation or in the finite width approach. In these cases a model dependent approach is therefore necessary to get bounds on the parameters of the model.

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