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Diboson signals via Fermi scale spin-one states

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ATLAS and CMS observe deviations from the expected background in diboson invariant mass searches of new resonances around 2 TeV. We provide a general analysis of the results in terms of spin-one resonances and find that Fermi scale composite dynamics can be the culprit. The analysis and methodology can be employed for future searches at run two of the Large Hadron Collider.

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I. INTRODUCTION

The ATLAS search for diboson resonances using boson-tagged jets [1] finds local excesses of 3.4 , 2.6 and 2.9σ in the WZ , WW and ZZ tagged boosted dijets with invariant mass spectrum around 2 TeV. This leads to a global significance of 2.5σ . Similarly CMS finds an excess of 1.9σ global significance in a boosted search for WH with the Higgs decaying hadronically [2].

The ATLAS results in the VV' gauge boson channels suggest a reconstructed mass of around 2 TeV and a cross-section $\sigma(pp \rightarrow R \rightarrow VV') \equiv \sigma_{VV'}$ of the order of (6–10) fb. Here R denotes a new intermediate massive vector boson and V , V' weak gauge bosons. The lower value corresponds to the maximization of the likelihood model based on Poisson statistics for the W' peak described in the ATLAS reference.¹ The upper value is extracted from the most stringent upper limits provided by other diboson searches in ATLAS and CMS, which is provided by ATLAS semileptonic channel [3] and is still in good agreement with the peak excess found by ATLAS. $\sigma_{VV'} = (6–10)$ fb will be our *region of interest*.

Here we employ a minimal description of spin-one resonances and study their phenomenology in the narrow width approximation. The model encompasses, however, all the needed ingredients to describe the *signal channels* and relevant constraints. We then make contact with time-honored models of minimal composite dynamics [4–7].

¹We use the likelihood

$$L = \prod_i P_{\text{poi}}(n_{\text{obs}}^i | n_{\text{exp}}^i) \quad (1)$$

where $P_{\text{poi}}(n|\lambda)$ is the Poisson distribution of n for a mean value λ , the index i correspond to the bins in the reconstructed diboson mass distribution of the WZ channel, Fig. 5(a) in Ref. [1], centered from $m_{jj} = 1750$ GeV to 2350 GeV, n_{obs} are the number of events observed in the respective bin and n_{exp} is the number of events expected for the sum of standard model processes plus the peak contribution from the W' model used in the experimental paper scaled with a signal strength μ .

Weinberg and Susskind's minimal models of weak scale, also known as the Fermi scale, composite dynamics [4,5] are based on QCD-like dynamics and are at odds with experiments. On the other hand, modern incarnations that are still minimal but employ non-QCD like dynamics are phenomenologically viable [8–10]. The associated *signal channels* have been investigated in more complete model implementations, e.g. in [11].

Although the data are not yet conclusive, the general features, regarding resonance mass, cross-section and decay patterns are very much in line with models of weak scale compositeness [11].

II. SPIN-ONE LAGRANGIAN

Since we are interested in the hadronic production and diboson decays of spin-one resonances R we consider the simplified effective Lagrangian

$$\mathcal{L}^R = \mathcal{L}_{qq}^R + \mathcal{L}_{VV}^R + \mathcal{L}_{VH}^R + \mathcal{L}_{\mathcal{X}}^R \quad (2)$$

where qq denotes quarks, and H is the observed Higgs state while \mathcal{X} is everything else, e.g leptons and dark matter. Correspondingly the width of the resonance can be written as

$$\Gamma_R = \Gamma_{qq} + \Gamma_{VV} + \Gamma_{VH} + \Gamma_{\mathcal{X}}. \quad (3)$$

The vertices linking the spin-one resonances with the standard model fermions are

$$\mathcal{L}_{qq}^R = \sum_{u,d} \bar{u} R^+ (g_{ud}^V - g_{ud}^A \gamma_5) d + \text{H.c.} \quad (4)$$

$$+ \sum_{ij} \bar{q}_i R^0 (g_{ij}^V - g_{ij}^A \gamma_5) q_j, \quad (5)$$

where u (d) runs over all up-type (down-type) quarks and q runs over all quark flavors, and we also further make the simplifying assumption $g_{ij}^{V/A} = g_{ud}^{V/A} = g_{V/A}$.

Neglecting CP -violating terms (see e.g. [12] for a more complete discussion) the couplings of neutral $R \equiv R_\mu^0$ to standard model gauge fields can be written as

$$\begin{aligned} \mathcal{L}_{VV}^R = & g_{WW1}^R [[RW^+W^-]]_1 + g_{WW2}^R [[RW^+W^-]]_2 \\ & + g_{WW3}^R [[RW^+W^-]]_\epsilon + g_{ZZ}^R [[RZZ]]_\epsilon \end{aligned} \quad (6)$$

where

$$\begin{aligned} [[RW^+W^-]]_1 & \equiv 2i[\partial_\mu W_\nu^+ W^{\mu-} R^\nu - \partial_\mu W_\nu^- W^{\mu+} R^\nu], \\ [[RW^+W^-]]_2 & \equiv \frac{i}{2}(\partial_\mu R_\nu - \partial_\nu R_\mu)(W^{\mu+}W^{\nu-} - W^{\mu-}W^{\nu+}), \\ [[RV_1V_2]]_\epsilon & \equiv \epsilon^{\mu\nu\rho\sigma}(V_{1\mu}\partial_\rho V_{2\nu} - \partial_\rho V_{1\mu}V_{2\nu})R_\sigma. \end{aligned} \quad (7)$$

We also have $L_{ZH}^R = g_{ZH}^R R_\mu Z^\mu H$. An equivalent Lagrangian can be defined for the charged spin-one resonances which we omit for brevity.

III. PRODUCTION CROSS SECTION AND WIDTHS

We now show in Fig. 1 the reference production cross-sections for a neutral and charged spin-one resonance with $g_V = 1$ and $g_A = 0$. We used the NNPDF2.3 set of parton distribution functions [13] and the MADGRAPH5_AMC@NLO framework [14] to get cross sections, with the Universal Feynrules Model [15] described in [16]. For a 2 TeV vector mass we have

$$\begin{aligned} \sigma_{\text{Ref}}(pp \rightarrow R^0) &= 1.5 \times 10^3 \text{ fb}, & m_R^0 &= 2 \text{ TeV} \\ \sigma_{\text{Ref}}(pp \rightarrow R^\pm) &= 1.7 \times 10^3 \text{ fb}, & m_R^\pm &= 2 \text{ TeV}. \end{aligned} \quad (8)$$

The partial widths for the neutral vector are well approximated by

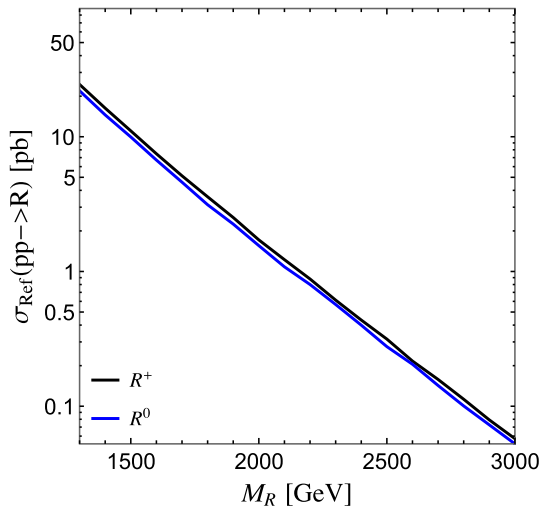


FIG. 1 (color online). Production cross-section $\sigma_{\text{Ref}}(pp \rightarrow R)$ with $g_V = 1$ and $g_A = 0$ for the couplings to fermions.

$$\begin{aligned} \Gamma(R \rightarrow q\bar{q}) & \simeq \frac{m_R N_c}{12\pi} [(g^V)^2 + (g^A)^2], \\ \Gamma(R \rightarrow W^+W^-) & \simeq \frac{1}{192\pi} m_R \left(\frac{m_R}{m_W}\right)^4 (g_{WW2}^R)^2, \\ \Gamma(R \rightarrow ZW) & \simeq \frac{1}{192\pi} m_R \left(\frac{m_R}{m_W}\right)^4 (g_{ZW2}^R)^2, \\ \Gamma(R \rightarrow ZZ) & \simeq \frac{(g_{ZZ}^R)^2}{96\pi} m_R \frac{m_R^2}{m_Z^2}, \\ \Gamma(R \rightarrow ZH) & \simeq \frac{(g_{ZH}^R)^2}{192\pi m_Z^2} m_R. \end{aligned} \quad (9)$$

Here $N_c = 3$ is the number of colors and in the following we will set $g_A = 0$.

IV. CONSTRAINTS

Given the production cross-sections (8), the total cross-section into diboson final states is

$$\begin{aligned} \sigma_{VV} &= g_V^2 \times Br[R \rightarrow VV] \times \sigma_{\text{Ref}}(pp \rightarrow R) \\ &= c_{VV} \sigma_{\text{Ref}}(pp \rightarrow R), \end{aligned} \quad (10)$$

where e.g. $c_{WZ} \sim (3-6) \times 10^{-3}$, in order to have $\sigma_{WZ} \sim (6-10)$ fb. In particular we must have $g_V^2 \gtrsim c_{WZ}$ once we require a certain diboson cross-section.

We now determine the lower limit on the resonance width, required to explain the excesses using both dijets and dibosons (into boosted jets) searches. The dijet cross-section is

$$\begin{aligned} \sigma_{qq} &= g_V^2 \times Br[R \rightarrow qq] \times \sigma_{\text{Ref}}(pp \rightarrow R) \\ &= c_{qq} \sigma_{\text{Ref}}(pp \rightarrow R) \end{aligned} \quad (11)$$

for which the current limits at 2 TeV require $c_{qq} \lesssim 0.1$ [17]. It follows that the total width is

$$\begin{aligned} \frac{\Gamma_{R_V}}{m_{R_V}} & \gtrsim \frac{\Gamma_{qq}}{m_{R_V}} \left(1 + \frac{c_{WZ} + c_{WH}}{c_{qq}}\right) + \frac{\Gamma_{\mathcal{X}}}{m_{R_V}} \\ & \simeq g_V^2 \frac{N_f \times 3}{12\pi} \left(1 + \frac{c_{WZ} + c_{WH}}{c_{qq}}\right) + \frac{\Gamma_{\mathcal{X}}}{m_R}, \end{aligned} \quad (12)$$

where we have taken N_f light quark flavors. So again requiring $\sigma_{WZ} \sim (6-10)$ fb we find $\frac{\Gamma_{R_V}}{m_{R_V}} \gtrsim (8-15) \times 10^{-4}$. Taking into account the dilepton decay modes, included in $\Gamma_{\mathcal{X}}$, will add a subdominant contribution to the width. The resulting limit shows that the narrow width approximation is justified.

A second constraint arises from dilepton searches. Of course this is a model dependent constraint that is, for

example, absent in leptophobic models. The current LHC limits for a single charged or neutral vector resonance are given in Fig. 2. We refer to [18] for a discussion of the details of the plot.

We conclude from the plot that the current limits on $\sigma(pp \rightarrow R \rightarrow \ell\nu)$ at $m_R = 2$ TeV translate into the constraint $g_V^2 \times Br[R \rightarrow \ell\nu] \lesssim 5 \times 10^{-4}$. Using that $g_V^2 \gtrsim c_{WZ} \sim (3-6) \times 10^{-3}$ we have that $Br[R \rightarrow \ell\nu] \lesssim (16-8) \times 10^{-2}$. The conclusions are similar for the neutral resonances, with respect to dijet and dilepton final states. Later we shall see that such a constraint is naturally abided by minimal models of weak composite dynamics.

Other leptonic and semileptonic searches are weaker than the ATLAS semileptonic limit on σ_{WZ} . The ATLAS fully leptonic search for resonances decaying into WZ [25] limits the cross section $\sigma(pp \rightarrow R \rightarrow WZ) \lesssim 22$ fb at $m_R = 2$ TeV. However it does not extend beyond 2 TeV. ATLAS semileptonic and CMS fully hadronically boosted analyses for resonances decaying into WZ give upper limits for the total cross section of around 10 fb and 12 fb respectively. Similarly in WW and ZZ semi-leptonic searches the upper limit on the total cross sections are higher than the semi-leptonic limits.

V. ZZ-CHALLENGE

Given that the spin-one states are weak triplets it is natural to expect signals in WZ and WW channels. On the other hand, large contributions to the ZZ decay mode requires anomalous couplings violating P and/or CP . Given that the mass resolution in the jet mass reconstruction of the W and Z is ± 13 GeV, it is logically possible that the ZZ reconstructed events do in fact involve W s. In the future with sufficient statistics in leptonic final states, this channel will be a diagnostic of the resonance nature. This possibility is demonstrated in [26].

VI. FERMION (GOLDSTONE) COMPOSITE DYNAMICS AND LATTICE RESULTS

We now argue that models of spin-one resonances from composite dynamics at the electroweak scale can have the required mass, production cross section and partial widths to explain the observed excesses at ATLAS.

For the spectrum we use first principle lattice results [27–30]. Specifically in models based on $SU(3)$ with fermions in the 2-index symmetric representation [8] recent lattice results indicate that the lightest vector and axial

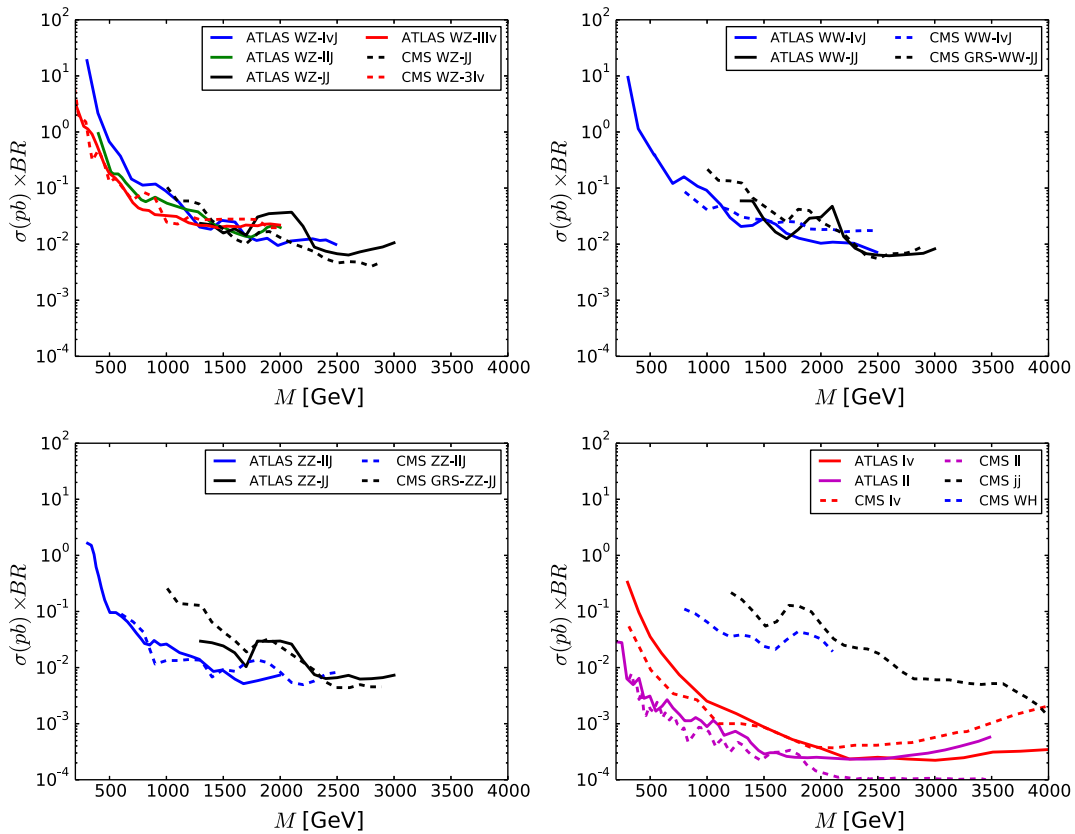


FIG. 2 (color online). ATLAS and CMS 95% CL exclusion limits on the production cross section times branching ratio, $\sigma \times BR$, for a new neutral or charged vector resonance. A capital J indicates W or Z reconstructed via jets. Upper left: WZ searches. Upper right: WW searches. Bottom left: ZZ searches. Bottom right: dilepton, dijets and WH searches. The experimental references are [1–3,17,19–25].

triplets have masses of [28] $m_{R_V} \approx 1.75 \pm 0.1$ TeV, $m_{R_A} \approx 2.3 \pm 0.1$ TeV.

Lattice results for $SU(2)$ driven composite dynamics with fermions in the fundamental representation [10,31,32] yield [30] $m_{R_V} \approx 2.5 \pm 0.5$ TeV, $m_{R_A} \approx 3.3 \pm 0.7$ TeV.

It is worth mentioning that even the simple scaled up version of QCD suggests the existence of a spin-one vector of mass around 2 TeV.

The examples above assume the electroweak condensate scale of 246 GeV to be identified with the decay constant of the composite pions, i.e. the ones absorbed in the longitudinal degrees of freedom of the massive weak gauge bosons. For the $SU(2)$ fundamental dynamics case with chiral symmetry breaking pattern $SU(4) \rightarrow Sp(4)$, however, one can imagine a more general electroweak embedding, parametrized by an angle $0 \leq \theta \leq \pi/2$, that for θ close to zero yields the composite Goldstone Higgs scenario [6,7]. Recent realizations were considered in [33–38]. As a function of the embedding angle we have $m_R(\theta) = m_R / \sin \theta$ with m_R given above for $\theta = \pi/2$ [38]. If, for example, one requires some typical values of $\theta \approx 10^{-1}$, the lattice predicted vector resonances states would be too heavy to be observable at the LHC. A comprehensive phenomenological analysis of the intriguing scalar sector of the theory [38] appeared in [39] for any value of theta.

Other underlying theories are also being studied on the lattice, see for example [40].

Further, in minimal weak scale composite models spin-one resonances couple to standard model fermions via mixing with the electroweak gauge bosons. Additional model building can also yield direct couplings to the fermions. From the mixing alone one deduces $g_V \sim \frac{g^2}{\sqrt{2}g}$ where one power of g comes from the W gauge eigenstate coupling to ud while $\frac{g}{\sqrt{2}} \sim \frac{m_W}{m_R}$ arises from the mass mixing. \tilde{g} is the self-coupling of the new spin-one mesons. Further $\mathcal{O}(1)$ corrections depend on the parameters of the mass

mixing Lagrangian [11]. Using the value of the weak coupling we have $g_V \sim 0.25/\tilde{g}$ and therefore to achieve $g_V^2 Br[R \rightarrow WZ] = c_{WZ} \sim 3-6 \times 10^{-3}$ implies $\tilde{g} \lesssim 4.5$. This is a very natural value of \tilde{g} since for a composite spin-one state we expect $1 \lesssim \tilde{g} \lesssim 4\pi$. In QCD, for example its value is about 2π .

Early studies of weak scale minimal composite dynamics [11] find the spin-one resonance width $\Gamma_R \sim 10^2$ GeV at 2 TeV for $\tilde{g} \sim 5$ with dilepton branching ratios at the level of 10^{-3} such that $g_V^2 Br[R \rightarrow \ell\ell] \sim 10^{-5}-10^{-4}$ [11].

Moreover we expect both vector and axial spin-one triplets with (mainly) axial weak triplets having further significant decay modes into HW and HZ final states [11].

Flavor constraints have been discussed in much detail in [41] for theories featuring spin-one resonances.

Amusingly the lattice results for the spin-one spectrum encompass the one needed to explain the experimental excesses and the deduced couplings to standard model fermions are naturally of the expected order of magnitude.

The run two experiment at the LHC will be able to either confirm or dismiss this intriguing possibility.

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Note added.—While this work was being finalized the analysis [26] appeared. Our discussion of spin-one resonances is fairly general and does not rely on a specific underlying model. Moreover we use lattice data that provide the spectrum of spin-one resonances relevant to experiments.

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