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Interpreting a 2 TeV resonance in *WW* scattering

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A diboson excess has been observed—albeit with very limited statistical significance—in WW, WZ, and ZZ final states at the LHC experiments using the accumulated 8 TeV data. Assuming that these signals are due to resonances resulting from an extended symmetry breaking sector in the standard model and exact custodial symmetry, we determine using unitarization methods the values of the relevant low-energy constants in the corresponding effective Lagrangian. Unitarity arguments also predict the widths of these resonances. We introduce unitarized form factors to allow for a proper treatment of the resonances in Monte Carlo generators and a more precise comparison with experiment.

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I. INTRODUCTION

In a series of recent papers [1–6], the relation between the coefficients of an effective Lagrangian parametrizing an extended electroweak symmetry breaking sector (EWSBS) and the appearance of narrow resonances in several isospin and angular momentum channels involving the scattering of longitudinally polarized W, Z bosons has been clearly established. It was found that, except for a small set of points in the space of parameters very close to the minimal standard model (MSM) values, resonances with these characteristics should appear. In fact, it was argued that detecting such resonances, if ever found, could provide an indirect but effective way of determining anomalous triple and quartic gauge boson vertices.

The connection between resonances and coefficients of the effective EWSBS Lagrangian is not based on a fully rigorous mathematical theorem, but it is amply supported by a wealth of experience on strong interactions and unitarization techniques in effective theories [7]. In the present context, results have been provided by two different groups. In [1,3] some of the present authors found by using the inverse amplitude method (IAM) of unitarization the relation between the characteristics of the first resonance in the various IJ channels ($I =$ custodial isospin) and the value of the coefficients of the effective Lagrangian. The analysis was done making only as minimal as possible usage of the equivalence theorem [8,9] as this is known to be prone to substantial corrections at low values of s. The Madrid group [4–6] making use of the equivalence theorem have also been able to determine the connection between resonances and departures from the MSM at an effective Lagrangian level. The agreement between the two independent sets of calculations is excellent whenever they can be compared. In addition, the Madrid group has done a careful analysis of different unitarization methods [6].

Unitarization leads to various resonances depending on the values of the effective couplings. There is also an ample region of parameter space ruled out as viable effective theories, something that is not a surprise to effective theory practitioners [10]. While there is certainly some room for some quantitative differences between different unitarization methods, the results are generally believed to be fairly accurate.

In the present discussion, by unitarization we refer to the reconstruction of a unitary amplitude using tree-level plus one-loop results. Several works considering the so-called tree-level unitarity (i.e., the requirement that amplitudes of the kind considered here do not grow with s) already exist [11].

Recently, the experimental collaborations ATLAS and CMS have reported [12,13] a modest excess of diboson events peaking around the 2 TeV region. ATLAS looks for the invariant mass distribution of a pair of jets that are compatible with a highly boosted W or Z boson. CMS combines dijet and final states with one or two leptons and concludes that there is a small excess around 1.8 TeV but with less statistical significance. In what follows, we shall use the ATLAS results assuming a mass for a putative resonance in the range 1.8 TeV $< M <$ 2.2 TeV.

In hadronic decays such as the ones used by ATLAS, it is not always possible to establish the nature of the jet (W or Z) [14]. Yet the experimental collaboration feels confident enough to claim that the signal is apparently present in the three channels WW, WZ, and ZZ. Assuming exact custodial symmetry, this would suggest that the resonance could not have $I = 0$, as this would not contribute in the s channel to WZ scattering, where the signal appears to be stronger.

However, elementary isospin arguments forbid a resonant contribution with $I = 1$ in processes with a ZZ final state. Therefore, assuming exact custodial symmetry, whether the resonance has either $I = 0$ or $I = 1$, one of the "observed" channels must have necessarily been misidentified [14].

The alternative to accepting $\mathcal{O}(1)$ custodial breaking would be to contemplate a resonant $I = 2$ state (contributes to all final states), but we regard this as unlikely for the reasons described in detail in [3,15] (but see [16] where an elementary $I = 2$ state is introduced).

As previously mentioned, the IAM and other unitarization prescriptions are widely used in strong interactions to derive poles in the S matrix that are not reachable in perturbation theory. Given a set of low-energy constants accompanying the higher-dimensional operators in chiral perturbation theory, one is able to predict different resonances with great precision. In addition, the unitarization method also gives predictions for their widths and production cross sections. Of course, one could follow a second avenue: If the lowest-lying resonances are provided in the different channels *and* their decay constants are also known, one can integrate these resonances out and derive the low-energy constants in the effective chiral Lagrangian (as was, for instance, done in [17]). If one uses these values for the low-energy constants and unitarizes, the initial resonances masses, decay constants, and widths are approximately recovered [7].

When considering physics beyond the standard model (BSM), one does not know the masses or coupling constants of the new states or resonances presumed to exist (either fundamental or composite), and, therefore, the values of the low-energy constants are *a priori* unknown unless one resorts to specific models [18]. Therefore, the second avenue is not open to us unless one proceeds model by model (and even in this case, if the BSM model is strongly interacting, the coupling constants and masses are output rather than input and, therefore, in general not known precisely). Does this mean that one cannot relate possible poles in the S matrix to the effective theory? Obviously, nothing prevents us from carrying out the same program that one usually does in strong interactions (that works equally well for weakly coupled theories). The IAM provides a direct connection between the poles of the S matrix and the low-energy constants of the effective theory. This relation can be read either way, and while it is normally used to predict resonances given a set of coefficients, we plan to use the information obtained from possible resonances in WW scattering (or the absence thereof) to constrain the effective Lagrangian.

This procedure is possible because in the nonlinear realization in the custodial limit, there are only two independent next-to-leading operators, and projecting on specific channels allows us to constrain two different combinations of the low-energy constants.[1] In addition, the method is capable of providing from the mere knowledge of the masses of the lowest-lying resonances their widths, production cross sections, and decay constants (this last prediction is presented in this paper for the first time in the present BSM context, we believe).

In this article, we shall contemplate the hypothesis that the resonant behavior apparently seen at the LHC is due either to an $I = 0;\ J = 0$ resonance and/or an $I = 1;\ J = 1$ one and make use of the IAM to derive very restrictive bounds on a combination of two coefficients of the effective Lagrangian. In addition, we will be able to approximately determine the widths of these putative resonances. The allowed regions in parameter space partly overlap; namely, there are regions with both a scalar and vector resonances (this would, of course, help to explain the excess in all channels). We will comment on the respective possible widths and masses. We will see that the range of masses contemplated here would lead to a severe reduction in the range of variation of the low-energy constants providing precious information to disentangle the class of underlying physics that one could be contemplating.

It should be mentioned that the method is of interest, independent of whether or not the mentioned signal is confirmed in the future. It is also worth emphasizing that resonances in the TeV region do correspond to very small values of the low-energy constants of the effective Lagrangian. These values are in no way unnatural as simple power-counting arguments suggest. These small low-energy constants lead also to very small modifications of the triple and quartic vector boson vertex, extremely hard to determine otherwise.

One salient characteristic of the resonances found in the mentioned unitarization analysis is that they are *very narrow*, something that runs contrary to the intuition of many practitioners in strongly interacting theories. This comes about because of the strong but partial unitarization that a Higgs at $M_H = 125$ GeV brings about. By construction, these resonances couple *only* to W and Z bosons. Together with the assumption of exact custodial symmetry, this is the only hypothesis in our analysis.

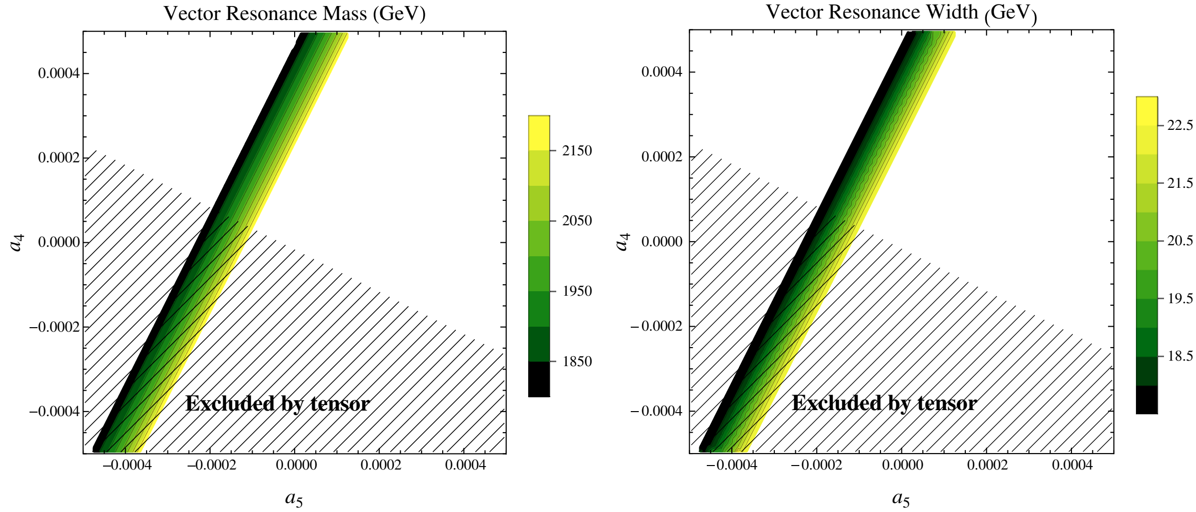
II. CONSTRAINING THE EFFECTIVE LAGRANGIAN COEFFICIENTS

The effective Lagrangian whose unitarized amplitudes we will consider is

$$\begin{aligned}\mathcal{L} = &-\frac{1}{2}\mathrm{Tr}W_{\mu\nu}W^{\mu\nu} - \frac{1}{4}\mathrm{Tr}B_{\mu\nu}B^{\mu\nu} + \frac{1}{2}\partial_\mu h \partial^\mu h - \frac{M_H^2}{2}h^2 \\ &- d_3(\lambda v)h^3 - d_4\frac{\lambda}{4}h^4 \\ &+ \frac{v^2}{4}\left(1 + 2a\left(\frac{h}{v}\right) + b\left(\frac{h}{v}\right)^2 + \cdots\right) \\ &\times \mathrm{Tr}D_\mu U^\dagger D^\mu U + \sum a_i \mathcal{O}_i, \end{aligned} \tag{1}$$

where

[1] The IAM method could, in principle, provide information for higher-order terms in the effective Lagrangian from excited resonances, but this requires knowledge of higher powers of s in the corresponding amplitudes, i.e., higher loops.



(a) (b)

FIG. 1. For $a = 1$ and $b = 1$: (a) Allowed values for a_4, a_5 corresponding to a vector resonance with a mass between 1.8 and 2.2 TeV. Note the extremely limited range of variation that is allowed in the figure for the low-energy constants. (b) The corresponding widths as predicted by unitarity using the IAM method. The characteristic value is 20 GeV—quite narrow for such a large mass. The dashed area is excluded on causality grounds stemming from the $I = 2$ channel.

$$U = \exp\left(i\frac{w\cdot\tau}{v}\right) \quad \text{and}$$
$$D_\mu U = \partial_\mu U + \frac{1}{2} igW^i_\mu \tau^i U - \frac{1}{2} ig' B^i_\mu U\tau^3. \qquad (2)$$

The w's are the three Goldstone of the global group $SU(2)_L \times SU(2)_R \to SU(2)_V$. This symmetry breaking is the minimal pattern to provide the longitudinal components to the $W^\pm$ and Z and emerging from phenomenology. The Higgs field h is a gauge and $SU(2)_L \times SU(2)_R$ singlet, and the $\mathcal{O}_i$ is a set of higher-dimensional operators. In an energy expansion and at the next-to-leading order, it is sufficient to consider the $\mathcal{O}(p^4)$ operators. This formulation is strictly equivalent to others where the Higgs is introduced as part of a complex doublet, as S-matrix elements are independent of the parametrization.

The operators $\mathcal{O}_i$ include the complete set of operators defined, e.g., in [1,19,20]. We will be interested in WW scattering and work in the strict custodial limit. Therefore, only a restrict number of operators have to be considered; namely, of the possible 13 $\mathcal{O}(p^4)$ operators, only two $\mathcal{O}_4$ and $\mathcal{O}_5$ will contribute to $W_L W_L$ scattering[2] in the custodial limit:

$$\mathcal{O}_4 = \mathrm{Tr}[V_\mu V_\nu]\mathrm{Tr}[V^\mu V^\nu], \qquad \mathcal{O}_5 = \mathrm{Tr}[V_\mu V^\mu]\mathrm{Tr}[V_\nu V^\nu], \qquad (3)$$

where $V_\mu = (D_\mu U)U^\dagger$. We could easily extend the analysis to include noncustodial contributions, but we see little or no reason to do so at present.

The parameters a and b control the coupling of the Higgs to the gauge sector [21]. Couplings containing higher powers of h/v do not enter WW scattering, and they have not been included in (1). The two additional parameters d_3 and d_4 parametrize the three- and four-point interactions of the Higgs field.[3] The MSM case corresponds to setting $a = b = d_3 = d_4 = 1$ in Eq. (1). Current LHC results give the following bounds for a, $a_{4,5}$:

$$a = [0.67, 1.33], \qquad a_4 = [-0.094, 0.10],$$
$$a_5 = [-0.23, 0.26] \quad 90\%\ \text{C.L.}; \qquad (4)$$

see [22,23]. Present data clearly favor values of a close to the MSM value ($a = 1$). We shall consider here only this case, leaving the consideration of other values of a to a forthcoming publication.[4] The parameter b is almost totally undetermined at present and actually does not play a very

[2] Actually, there are other custodial invariant operators (see, e.g., the Appendix in [1] for comments and notation). Of these, only the one corresponding to the coefficient a_3 actually appears in $W_L W_L$ scattering. It acts by modifying the value of the diagrams where a W or Z is exchanged with an overall factor that for the range of values of the low-energy constants contemplated here is way too small to be of any influence. Detailed formulas can be found in [1,9]. Only a_4 and a_5 really matter.

[3] This is not the most general form of the Higgs potential, and, in fact, additional counterterms are needed beyond the standard model [4], but this does not affect the subsequent discussion for $W_L W_L$ scattering.

[4] It should be mentioned at this point that considering $a < 1$ leaves the vector cross section almost unchanged (although the range of a_4 a_5 is somewhat modified), it does increase noticeably the scalar cross section.

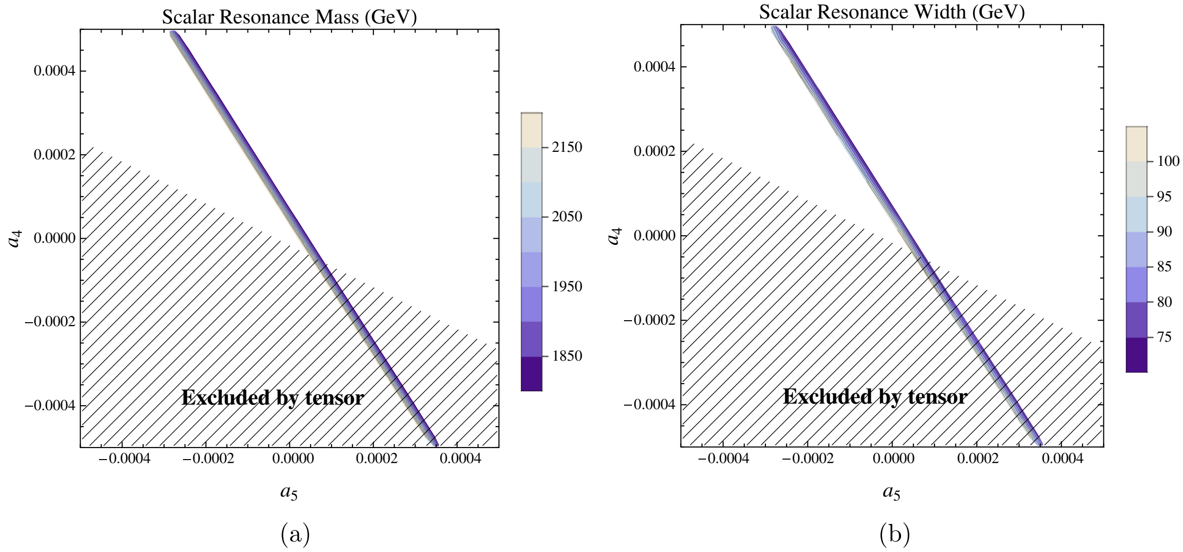


FIG. 2. For $a = 1$ and $b = 1$: (a) Allowed values for a_4, a_5 corresponding to a scalar resonance with a mass between 1.8 and 2.2 TeV. (b) The corresponding widths as predicted by unitarity using the IAM method; characteristic values are in the 70–100 GeV range.

relevant role in the present discussion. We will assume $b = a^2$ without further ado.

Determining the range of parameters a_4 and a_5 allowed by assuming a scalar and/or vector resonance in the range $1.8\ \text{TeV} < M < 2.2\ \text{TeV}$ is the main purpose of the present analysis. It should be mentioned that these two low-energy constants do not affect at all oblique corrections (quite constrained; see, e.g., [24]) or the triple gauge boson coupling: a_1, a_2, and a_3 are the relevant couplings in the custodial limit to consider in these contexts. The effective EWSBS Lagrangian nicely disentangles the two kinds of constraints.

We shall not provide here the technical details of the unitarization method we use, as they have been described in detail elsewhere [1,3].

After requiring a resonance in the vector channel with a mass in the quoted range, one gets in an a_4-a_5 plane the region shown on the left in Fig. 1 for $a = 1$. An analogous

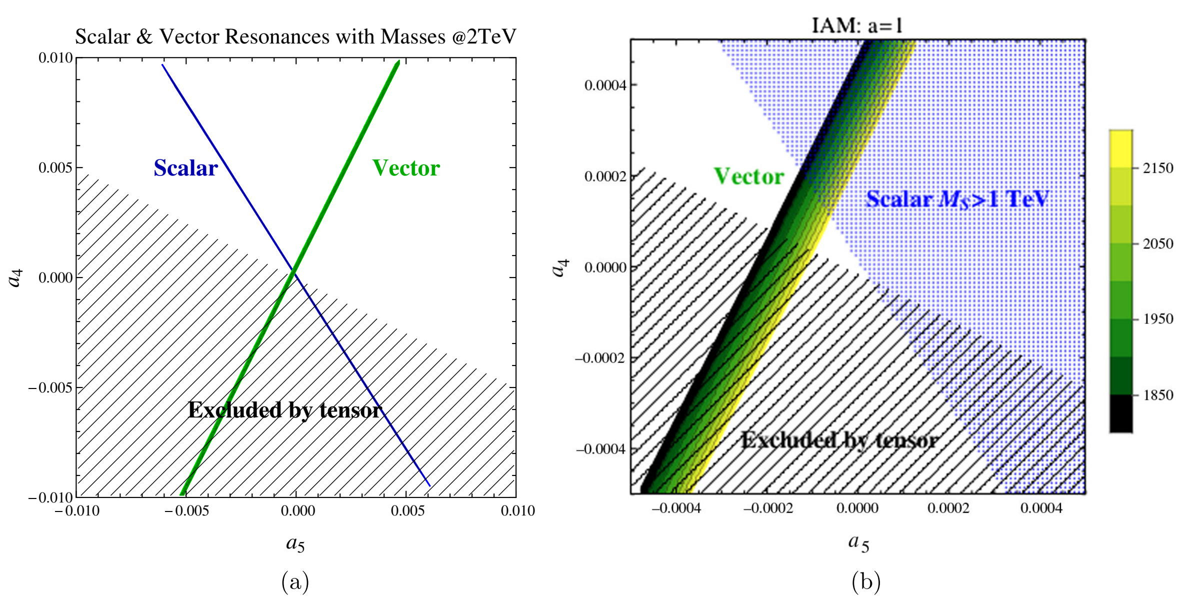


FIG. 3. (a) The existence of a resonance in the $1.8\ \text{TeV} < M < 2.2\ \text{TeV}$ range constraints a lot the allowed values of the low-energy constants of the EWSBS effective Lagrangian. The dashed area is excluded on causality grounds. (b) Blowup of the region of overlap where vector and scalar resonances may coexist. The broad strip shows the region of admissible vector resonances with masses in the 1.8–2.2 TeV range. The shaded area in the upper-right part contains scalar resonances of mass > 1 TeV.

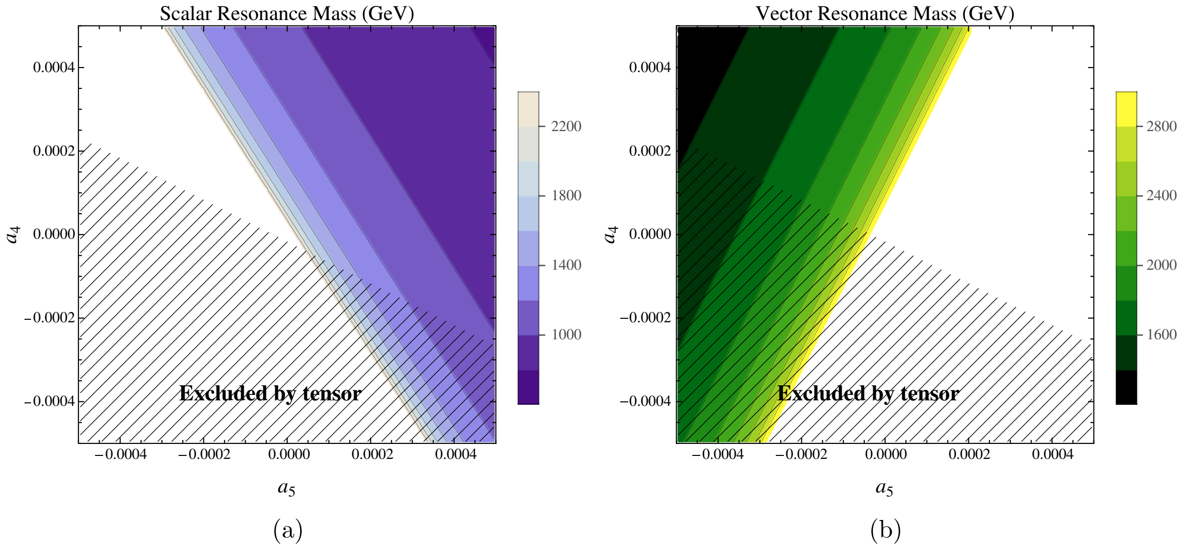


FIG. 4. (a) Viable scalar resonance masses in the region of interest in the a_4-a_5 plane for $a = 1$ assuming a vector resonance in the 1.8 TeV $< M <$ 2.2 TeV range. (b) The reverse situation: assuming a scalar mass in the 1.8 TeV $< M <$ 2.2 TeV range and depicting the possible values for a vector resonance compatible with it.

procedure, but assuming that the resonance is the $I = 0$, $J = 0$ channel, results in the allowed region in the a_4-a_5 plane depicted in Fig. 2.

We would like to emphasize the very limited range of variation for the parameters that is shown in Figs. 1 and 2. The constants a_4 and a_5 lie in the small region $|a_4|, |a_5| < 5 \times 10^{-4}$ (this region includes, of course, the MSM value $a_4 = a_5 = 0$, but, obviously, there are no resonances there).

In order to convey a picture of the sort of predictive power of unitarization techniques, we plot in Fig. 3 the allowed bands in the broader range $|a_4|, |a_5| < 0.01$ that was considered in a previous work [1] as still being phenomenologically acceptable. Indeed, setting even a relatively loose bound for the mass of the resonance restricts the range of variation of the relevant low-energy constants enormously. In the same Fig. 3, we show a blowup of the region where *both* a scalar and a vector resonance in this mass range may coexist. The dashed area is excluded as acceptable for effective EWSBS theories (see [3]). In Fig. 4 we show viable scalar and vector resonances in the a_4-a_5 plan.

III. EXPERIMENTAL VISIBILITY OF THE RESONANCES

The statistics so far available from the LHC experiments are limited. Searching for new particles in the LHC environment is extremely challenging, and analyzing the contribution of possible resonances to an experimental signal is not easy without a well-defined theoretical model with definite predictions for the couplings, form factors, etc. The IAM method is not only able to predict resonance masses and widths but also their couplings to the $W_L W_L$. In [1,3], the experimental signal of the different resonances was compared to that of a MSM Higgs with an identical mass. Because the decay modes are similar (in the vector boson channels that is) and limits on different Higgs masses are very documented, this was a rather intuitive way of presenting the cross section for possible EWSBS resonances, but it is not that useful for heavy resonances as the signal of a hypothetical Higgs of analogous mass becomes very broad and diluted. This point and several others were discussed in detail in [1]. Here we shall give very simple estimates of some cross sections based on the effective W approximation (EWA) [25] in a couple of channels. These estimates should be taken as extremely tentative and only relevant to establish comparisons between different masses and channels. In the last section, we will introduce form factors and vertex functions to allow for a proper comparison with experiment. Please note that amplitudes where scalars contribute the contribution of the 125 GeV Higgs are also included.

Some results for the cross sections are depicted in Fig. 5 for the processes $W_L^+ W_L^- \to W_L^+ W_L^-$ and $Z_L Z_L \to Z_L Z_L$. In the first case, we quote the contribution from a possible vector resonance only (a scalar resonance is also possible in this process). In the second case, only scalar exchange is possible. Note that both diboson production modes are subdominant at the LHC with respect to gluon production mediated by a top-quark loop and that the possible resonances in the scenario discussed here couple *only* to dibosons.

Compared to the preliminary experimental indications, the results quoted for the cross sections of these two specific processes are low, particularly for vector resonances, but there are several caveats. First of all, the EWA tends to underestimate the cross sections, and it is difficult to assess its validity in the present kinematical situation.

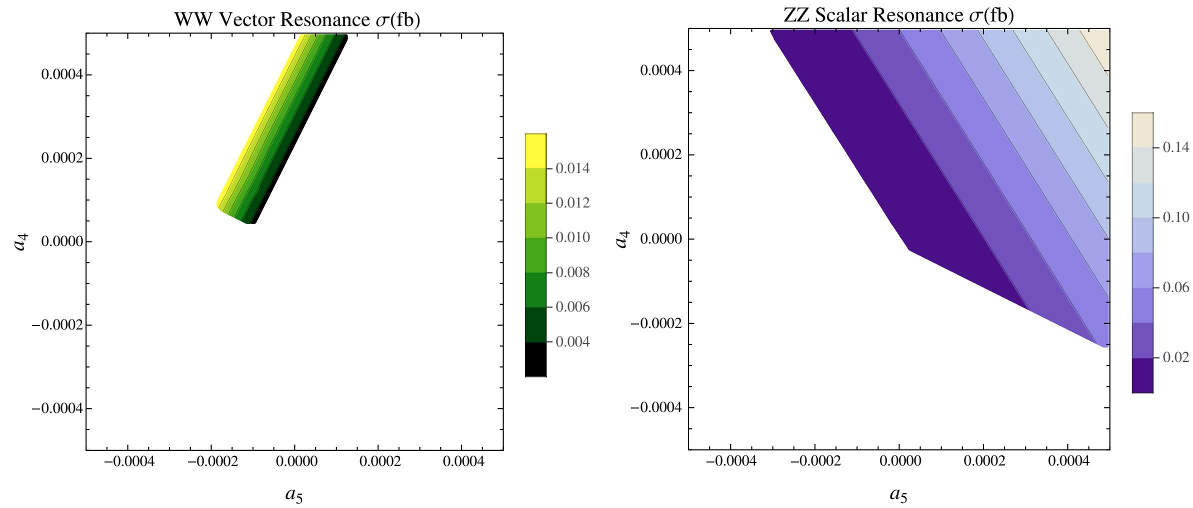


FIG. 5. Experimental signal of resonances for $a = 1$: The resonance cross sections are given in fb, the LHC energy has been taken to be 8 TeV, and the EWA approximation is assumed in this calculation. Left: Estimated cross section for the process $W_L W_L \to W_L W_L$ as a function of the parameters a_4, a_5 due to a vector resonance. Right: Cross section for the process $Z_L Z_L \to Z_L Z_L$ due to a scalar resonance. The contribution from the 125 GeV Higgs is also included in both cases.

Second, in this region of parameter space, the cross sections do change very quickly with only small changes of the parameters, thus, adding an element of uncertainty. Finally, the quoted cross sections correspond to considering only the interval $s \in [M - 2\Gamma, M + 2\Gamma]$ so as to have some intuition on the contribution of the resonance itself. It should also be mentioned that, as discussed in [1], there is an enhancement in the $W^+W^- \to W^+W^-$ channel when both the vector and scalar resonances become nearly degenerate; this is possible in a limited region of parameter space. Also, as previously stated, the scalar channel is enhanced if $a < 1$.

Interesting as partial waves for a given process may be, they are not that useful to implement unitarization in a Monte Carlo generator in order to make a detailed quantitative comparison with experiment. One would need to implement diagrammatically, and for that, one needs vertex functions and propagators wherewith to construct and compute the contribution from different topologies. Our proposal to tackle this problem is presented next.

IV. INTRODUCING FORM FACTORS

The amplitude $A(W_L^a(p^a) + W_L^b(p^b) \to W^c(p^c)_L + W^d(p^d)_L)$ will be denoted by $A^{abcd}(p^a, p^b, p^c, p^d)$. Using isospin and Bose symmetries, this amplitude can be expressed in terms of a universal function as

$$A^{abcd}(p^a, p^b, p^c, p^d) = \delta^{ab}\delta^{cd}A(s,t,u) + \delta^{ac}\delta^{bd}A(t,s,u) + \delta^{ad}\delta^{bc}A(u,t,s), \quad (5)$$

with $A(s,t,u) = A^{+-00}(p^+, p^-, p^0, p'^0)$. The fixed-isospin amplitudes are given by the following combinations:

$$\begin{aligned} T_0(s,t,u) &= 3A(s,t,u) + A(t,s,u) + A(u,t,s), \\ T_1(s,t,u) &= A(t,s,u) - A(u,t,s), \\ T_2(s,t,u) &= A(t,s,u) + A(u,t,s). \end{aligned} \quad (6)$$

In writing these expressions, we assume exact crossing symmetry.[5] We also write the reciprocal relations (also assuming exact crossing symmetry)

$$\begin{aligned} A^{+0+0}(s,t,u) &= \frac{1}{2}T_1(s,t,u) + \frac{1}{2}T_2(s,t,u), \\ A^{+-+-}(s,t,u) &= \frac{1}{3}T_0(s,t,u) + \frac{1}{2}T_1(s,t,u) + \frac{1}{6}T_2(s,t,u), \\ A^{++++}(s,t,u) &= T_2(s,t,u), \\ A^{0000}(s,t,u) &= \frac{1}{3}T_0(s,t,u) + \frac{2}{3}T_2(s,t,u). \end{aligned} \quad (7)$$

Other amplitudes [such as, e.g., $A^{+-00}(s,t,u)$] can be obtained trivially from the previous ones using obvious symmetries (and crossing symmetry too).

The partial wave amplitudes for fixed isospin I and total angular momentum J are defined by

$$t_{IJ}(s) = \frac{1}{64\pi}\int_{-1}^{1} d(\cos\theta) P_J(\cos\theta) T_I(s,t,u), \quad (8)$$

where the $P_J(x)$ are the Legendre polynomials and $t = (1 - \cos\theta)(4M_W^2 - s)/2$, $u = (1 + \cos\theta)(4M_W^2 - s)/2$

[5] This remark is pertinent because amplitudes involving longitudinally polarized bosons are not crossing symmetric. The formulas can be easily extended to this case but become somewhat more involved and will not be reported here. See [1].

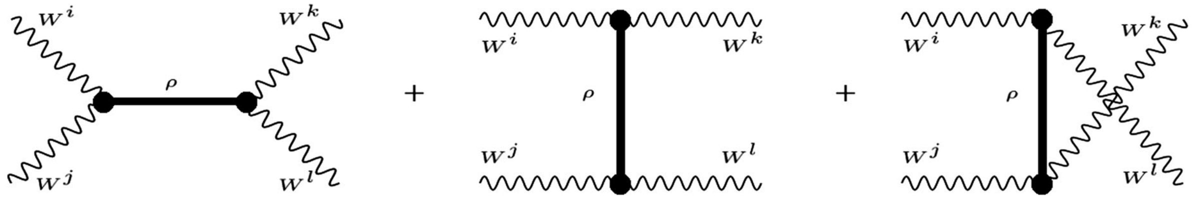


FIG. 6. Decomposition of a process in (unitarized) form factors and resonance propagators.

with M_W being the W, Z mass, and t_{00}, t_{11}, and t_{20} are the first nonvanishing partial waves in the present case. The poles in the respective unitarized partial wave amplitudes dictate the presence or absence of EWSBS resonances in the different channels.

We would like to express any amplitude as the sum of exchanges of resonances in the s, t, and u channels, as it is diagrammatically expressed in Fig. 6. That is, we decompose, say, A^{+0+0},

$$A^{+0+0} = \sum_{IJ} (A_s^{IJ} + A_t^{IJ} + A_u^{IJ}). \tag{9}$$

Not all IJ's receive contributions from all three channels. For example, in the case $A^{+0,+0}$ a possible scalar resonance only contributes to the t channel. In addition, not all processes are resonant in all regions of parameter space, so the above decomposition assumes resonance saturation. Let us now define the vector form factor as[6]

$$\langle W_L^i(p_1) W_L^j(p_2) | J_\mu^k | 0\rangle = (p_1 - p_2)_\mu F_V(s) \epsilon^{ijk}, \tag{10}$$

where J_k^μ is the interpolating vector current with isospin index k that creates the resonance ρ, and $F_V(s)$ is the vector form factor. From this form factor, we derive a vector vertex function K^μ via the relation

$$K^\mu(p_1, p_2) = (p_1 - p_2)^\mu F_V(s)(s - M_{\rm pole}^2). \tag{11}$$

Let us focus, for instance, on the amplitude A^{+0+0} that potentially has contributions from a vector and a tensor. The IAM does exclude the $I = 2$ contribution [3], so let us consider A_s^{11} for this process. It can be expressed as

$$\begin{aligned} A_s^{11} &= K^\mu \frac{g_{\mu\nu} - \frac{k_\mu k_\nu}{k^2}}{s - M_{\rm pole}^2} K^{*\nu} = |F_V(s)|^2 (s - M_{\rm pole}^{*2})(-2t - s) \\ &= |F_V(s)|^2 (s - M_{\rm pole}^{*2})(-s\cos\theta), \end{aligned} \tag{12}$$

where $M_{\rm pole} = M - i\Gamma/2$. Analogous decompositions exist for A_t^{11} and A_u^{11}. In fact, we do not need to consider A_t^{11} and A_u^{11} at all because assuming exact isospin symmetry, $A^{11}(s,t,u) = (-1)^I A^{11}(s,u,t)$. Here we assume, and it is a necessary ingredient of the present approach, that external lines are on shell.

On the other hand, from unitarization we know that

$$A^{11} \simeq 96\pi t_{11}(s) \cos\theta, \tag{13}$$

so neglecting further partial waves, it is natural to identify

$$|F_V(s)|^2 = -\frac{96\pi t_{11}(s)}{s(s - M_{\rm pole}^{*2})}, \tag{14}$$

where for t_{IJ} we can use the IAM approximation

$$t_{IJ} \approx \frac{t_{IJ}^{(0)}}{1 - t_{IJ}^{(2)}/t_{IJ}^{(0)}}. \tag{15}$$

Although $|F_V|^2$ should, of course, be real and positive, when using the identification above we get a tiny imaginary part ($\mathrm{Im}|F_V|^2 \sim 10^{-2}\mathrm{Re}|F_V|^2$) due to the fact that we are missing possible channels (including nonresonant contributions) and terms in the partial wave expansion. However, we can regard the description of the amplitude via vertex functions and resonance propagators as quite satisfactory in the regions where resonances are present.

Neglecting the gauge boson mass (quite justified at 2 TeV), unitarity requires the form factor to obey the following relation within a vector-dominance region [20]

$$\mathrm{Im}F_V(s) = t_{11}^*(s) F_V(s). \tag{16}$$

Equation (16) allows us to extract the phase of $F_V(s)$. Thus, combining the phase and the modulus, we obtain the vector form factor

$$F_V(s) = |F_V(s)| \exp\left(i \arctan \frac{\mathrm{Re}t_{11}}{1 - \mathrm{Im}t_{11}} \right). \tag{17}$$

Similar techniques could allow us to define a unitarized scalar form factor $F_S(s)$ and a vertex function directly derived from the unitarized amplitude that in this channel is

$$A^{00} \simeq 32\pi t_{00}(s) \tag{18}$$

and assuming resonance dominance. In Fig. 7, we plot the vertex functions $K_V(s)$ and $K_S(s)$ obtained by the method just described:

[6]Current vector conversation has been used.

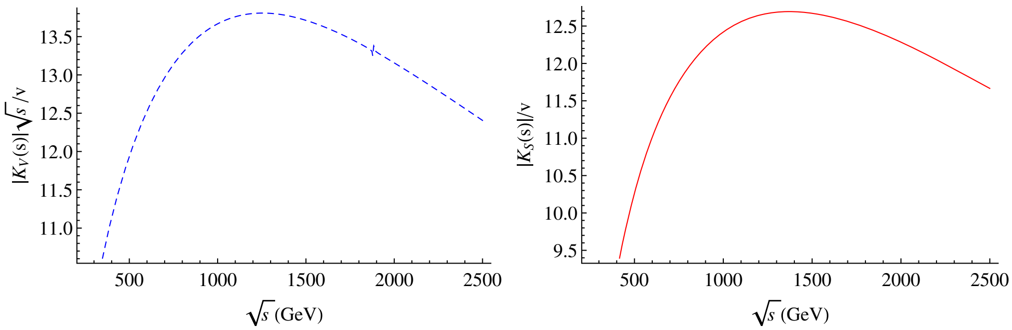


FIG. 7. Left: Plot of the effective coupling of the vector resonance $K_V(s)\sqrt{s}$ for the value $M = 1881$ GeV corresponding to $a = 1$, $a_4 = 0.0002$, $a_5 = -0.0001$. Right: Plot of the effective coupling for a scalar resonance $K_S(s)$ corresponding to the same values of a_4 and a_5 that yields a scalar mass $M = 2064$. Note that in both cases the coupling is quite large, certainly nonperturbative. In fact, on the scalar resonance, the effective coupling is ∼30 times the coupling of a MSM Higgs with identical mass.

$$|K_V(s)| \sim |F_V(s)||s - M^2_{\rm pole}|,$$
$$|K_S(s)| \sim |F_S(s)||s - M^2_{\rm pole}|. \quad (19)$$

Note that the function $K_V(s)$ is dimensionless while $K_S(s)$ has units of energy. However, for vector resonances, the effective coupling is typically $K_V(s)\sqrt{s}$ (see the expression for the form factor and the associated Feynman rule). In the last figure, we plot these effective couplings normalized to the scale v. The contribution to the form factor from the 125 GeV Higgs is negligible around the scalar resonance at 2 TeV.

We notice that, not unexpectedly, the effective couplings are momentum dependent, even though the dependence seems mild in a relatively large region around the resonance. It is not difficult to extract the value of the effective coupling constant at the resonance value. If one assumes an effective operator coupling, the vector resonance to the Goldstone bosons given by the only operator possible of dimensionality four

$$\sim g_V \epsilon_{ijk} w^i \partial^\mu w^k V^k_\mu, \quad (20)$$

the following value is obtained: $g_V(\sqrt{s} = 2\ \text{TeV}) \simeq 1.6$; namely the coupling is certainly nonperturbative but not enormously large.

In a resonance model dominated by vector meson exchange (similar to the one described in [17] in the context of QCD), the couplings a_4 and a_5 would be expected to be of order $g_V^2(\frac{M_W}{M_V})^2$ (plus radiative corrections) where g_V would be the (dimensionless) coupling constant introduced above and M_V the mass of the vector resonance. Introducing the value of g_V found and setting $M_V = 2$ TeV, one gets coefficients of order 10^{-4} as favored by the direct unitarity analysis. Notice that in the IAM approach, everything in the vector channel depends to a very good approximation on the single combination $a_4 - 2a_5$ (and g_V is an output), whereas using a phenomenological resonance model, one needs to know separately g_V and M_V and make a number of further assumptions (such as assuming neglecting the momentum dependence of form factors).

Once we feel confident that the combination of resonant propagators and the vertex functions just given reproduces very satisfactorily the unitarized amplitudes, we can pass on this information to Monte Carlo generator practitioners to implement these form factors in their favorite generator. However, generators include the creation and propagation of longitudinal (and transverse) W, not Goldstone bosons, so the above result cannot be directly carried over and more work is needed. This will be reported elsewhere in due course. Note that the unitarization proposed is manifestly crossing symmetric because crossed diagrams are included manifestly with the replacements $K(s) \to K(t)$, etc.

The expressions for $M_{\rm pole}$, $t_{00}(s)$, and $t_{11}(s)$ needed to reproduce the diagrammatic expansion for the various values of a and a_4, a_5 can be found in [1–3] (and [4,6] if a full use of the equivalence theorem is made[7]). Further details will be provided in a forthcoming extended publication.

V. CONCLUSIONS

To conclude, we extracted the values of the low-energy constants a_4 and a_5 of the effective Lagrangian describing an extended electroweak symmetry breaking sector assuming (iso)vector dominance and/or (iso)scalar dominance with a mass in the range 1.8 TeV $< M <$ 2.2 TeV, as it would be the case if one considers the preliminary results coming from the LHC experiment to be a hint of the existence of new $W_L W_L$ interactions. The calculation was performed in the framework of the inverse amplitude unitarization method. We derived the widths of such resonances, which turn out to be quite narrow. We also speculated on the possibility of more than one resonance being present, compatible with the derived bounds on a_4 and a_5 (something that is favored by

[7]Please note that t-channel W exchange is not included in some of these works.

custodial symmetry considerations). The given range of masses restricts enormously the admissible values for a_4 and a_5, surely a consequence of this mass scale being relatively close to the natural cutoff of the effective theory (∼3 TeV). The cross sections obtained using the effective W approximation are, however, too low, particularly for vector resonances, and this may eventually prove bad news for resonances of the kind considered here. However, we regard estimates based on the EWA as being too preliminary at this point.

To overcome this difficulty, we proposed a diagrammatic method to deal with resonances in regions of parameter space in the effective Lagrangian where the former are assumed to dominate. We derived the corresponding form factors and vertex functions. The agreement with the full amplitude is very good, and we understand that the technique that we introduced here may be useful to deal with the type of resonances that may emerge in EWSBS. We hope that this will trigger interest from our experimental colleagues to incorporate this seemingly consistent unitary procedure in their generators to allow for a proper theory-experiment comparison. In fact, having a reliable estimate of the resonances cross sections in the region of interest is probably the most urgent task.

The apparent signal coming from the LHC experiments has triggered a flurry of activity that has mostly concentrated on proposing specific models ranging from introducing resonances [26], the obvious possibility of excited or left-right symmetric W', Z' states to more exotic models [27,28,29]. Our proposal is somewhat different: It is not primarily aimed at advancing a definite *ad hoc* proposal but rather to help understand if the signal is there in the first place and at trying to elucidate the properties of the resonance (or resonances) that might be present in an extended electroweak symmetry breaking sector in WW scattering. We regard the restriction on some coefficients of the effective Lagrangian provided by unitarity considerations as nontrivial and, if confirmed, would undoubtedly play a relevant role in constraining the underlying model.

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[1] D. Espriu and B. Yencho, Phys. Rev. D **87**, 055017 (2013).
[2] D. Espriu, F. Mescia, and B. Yencho, Phys. Rev. D **88**, 055002 (2013).
[3] D. Espriu and F. Mescia, Phys. Rev. D **90**, 015035 (2014).
[4] R. Delgado, A. Dobado, and F. Llanes-Estrada, J. Phys. G **41**, 025002 (2014); J. High Energy Phys.02 (2014) 121; Phys. Rev. Lett. **114**, 221803 (2015).
[5] D. Barducci, H. Cai, S. De Curtis, F. Llanes-Estrada, and S. Moretti, Phys. Rev. D **91**, 095013 (2015).
[6] R. Delgado, A. Dobado, and F. Llanes-Estrada, Phys. Rev. D **91**, 075017 (2015).
[7] T. N. Truong, Phys. Rev. Lett. **61**, 2526 (1988); A. Dobado, M. J. Herrero, and T. N. Truong, Phys. Lett. B **235**, 134 (1990); A. Dobado and J. R. Pelaez, Phys. Rev. D **47**, 4883 (1993); **56**, 3057 (1997); J. A. Oller, E. Oset, and J. R. Pelaez, Phys. Rev. Lett. **80**, 3452 (1998); Phys. Rev. D **59**, 074001 (1999); **60**, 099906(E) (1999); F. Guerrero and J. A. Oller, Nucl. Phys. **B537**, 459 (1999); A. Dobado and J. R. Pelaez, Phys. Rev. D **65**, 077502 (2002); A. Filipuzzi, J. Portoles, and P. Ruiz-Femenia, J. High Energy Phys. 08 (2012) 080; *Proc. Sci.*, CD12 (2013) 053.
[8] J. M. Cornwall, D. N. Levin, and G. Tiktopoulos, Phys. Rev. D **10**, 1145 (1974); C. E. Vayonakis, Lett. Nuovo Cimento. **17**, 383 (1976); B. W. Lee, C. Quigg, and H. B. Thacker, Phys. Rev. D **16**, 1519 (1977); G. J. Gounaris, R. Kogerler, and H. Neufeld, Phys. Rev. D **34**, 3257 (1986); M. S. Chanowitz and M. K. Gaillard, Nucl. Phys. **B261**, 379 (1985); A. Dobado and J. R. Pelaez, Nucl. Phys. **B425**, 110 (1994); Phys. Lett. B **329**, 469 (1994); C. Grosse-Knetter and I. Kuss, Z. Phys. C **66**, 95 (1995); H. J. He, Y. P. Kuang, and X. Li, Phys. Lett. B **329**, 278 (1994).
[9] D. Espriu and J. Matias, Phys. Rev. D **52**, 6530 (1995).
[10] A. Adams, N. Arkani-Hamed, S. Dubovsky, A. Nicolis, and R. Rattazzi (CERN), J. High Energy Phys. 10 (2006) 014.
[11] C. Englert, P. Harris, M. Spannowsky, and M. Takeuchi, Phys. Rev. D **92**, 013003 (2015); G. Cacciapaglia and M. T. Frandsen, Phys. Rev. D **92**, 055035 (2015); T. Abe, R. Nagai, S. Okawa, and M. Tanabashi, Phys. Rev. D **92**, 055016 (2015).
[12] G. Aad *et al.* (ATLAS Collaboration), J. High Energy Phys. 12 (2015) 055.
[13] V. Khachatryan *et al.* (CMS Collaboration), J. High Energy Phys. 08 (2014) 173; R. Koegler, http://indico.cern.ch/event/388148/attachments/775833/.
[14] B. C. Allanach, B. Gripaios, and D. Sutherland, Phys. Rev. D **92**, 055003 (2015); J. A. Aguilar-Saavedra, J. High Energy Phys. 10 (2015) 099.
[15] A. Falkowski, S. Rychkov, and A. Urbano, J. High Energy Phys. 04 (2012) 073; A. Urbano, J. High Energy Phys. 06 (2014) 060.

[16] H. Georgi and M. Machacek, Nucl. Phys. **B262**, 463 (1985).
[17] G. Ecker, J. Gasser, A. Pich, and E. de Rafael, Nucl. Phys. **B321**, 311 (1989).
[18] B. Henning, X. Lu, and H. Murayama, J. High Energy Phys. 01 (2016) 023.
[19] A. Dobado, D. Espriu, and M. J. Herrero, Phys. Lett. B **255**, 405 (1991); D. Espriu and M. J. Herrero, Nucl. Phys. **B373**, 117 (1992); M. J. Herrero and E. Ruiz-Morales, Nucl. Phys. **B418**, 431 (1994); **B437**, 319 (1995); D. Espriu and J. Matias, Phys. Lett. B **341**, 332 (1995); R. Foadi, M. Jarvinen, and F. Sannino, Phys. Rev. D **79**, 035010 (2009).
[20] A. Dobado, M. J. Herrero, J. R. Peláez, and E. Ruiz-Morales, Phys. Rev. D **62**, 055011 (2000).
[21] G. F. Giudice, C. Grojean, A. Pomarol, and R. Rattazzi, J. High Energy Phys. 06 (2007) 045; G. Buchalla and O. Cata, J. High Energy Phys. 07 (2012) 101; R. Contino, M. Ghezzi, C. Grojean, M. Muhlleitner, and M. Spira, J. High Energy Phys. 07 (2013) 035; R. Alonso, M. B. Gavela, L. Merlo, S. Rigolin, and J. Yepes, Phys. Lett. B **722**, 330 (2013); R. Alonso, I. Brivio, B. Gavela, L. Merlo, and S. Rigolin, J. High Energy Phys. 12 (2014) 034; G. Buchalla, O. Catà, and C. Krause, Nucl. Phys. **B880**, 552 (2014); G. Cacciapaglia and F. Sannino, J. High Energy Phys. 04 (2014) 111.
[22] A. Falkowski, F. Riva, and A. Urbano, J. High Energy Phys. 11 (2013) 111.
[23] I. Brivio, T. Corbett, O. J. P. Eboli, M. B. Gavela, J. Gonzalez-Fraile, M. C. Gonzalez-Garcia, L. Merlo, and S. Rigolin, J. High Energy Phys. 03 (2014) 024.
[24] A. Pich, I. Rosell, and J. J. Sanz-Cillero, Phys. Rev. Lett. **110**, 181801 (2013); J. High Energy Phys. 01 (2014) 157.
[25] G. L. Kane, W. W. Repko, and W. B. Rolnick, Phys. Lett. **148B**, 367 (1984); S. Dawson, Phys. B **249**, 42 (1985); M. S. Chanowitz and M. K. Gaillard, Nucl. Phys. **B261**, 379 (1985).
[26] H. S. Fukano, M. Kurachi, S. Matsuzaki, K. Terashi, and K. Yamawaki, Phys. Lett. B **750**, 259 (2015); D. B. Franzosi, M. T. Frandsen, and F. Sannino, Phys. Rev. D **92**, 115005 (2015); A. Thamm, R. Torre, and A. Wulzer, Phys. Rev. Lett. **115**, 221802 (2015); T. Abe, T. Kitahara, and M. M. Nojiri, arXiv:1507.01681; A. Carmona, A. Delgado, M. Quiros, and J. Santiago, J. High Energy Phys. 09 (2015) 186; C. W. Chiang, H. Fukuda, K. Harigaya, M. Ibe, and T. T. Yanagida, J. High Energy Phys. 11 (2015) 015; L. Bian, D. Liu, and J. Shu, arXiv:1507.06018; M. Low, A. Tesi, and L. T. Wang, Phys. Rev. D **92**, 085019 (2015).
[27] A. Alves, A. Berlin, S. Profumo, and F. S. Queiroz, J. High Energy Phys. 10 (2015) 076; J. Brehmer, J. Hewett, J. Kopp, T. Rizzo, and J. Tattersall, J. High Energy Phys. 10 (2015) 182; Q. H. Cao, B. Yan, and D. M. Zhang, Phys. Rev. D **92**, 095025 (2015); B. A. Dobrescu and Z. Liu, J. High Energy Phys. 10 (2015) 118; G. Cacciapaglia, A. Deandrea, and M. Hashimoto, Phys. Rev. Lett. **115**, 171802 (2015); V. Sanz, arXiv:1507.03553; C. H. Chen and T. Nomura, Phys. Lett. B **749**, 464 (2015); M. E. Krauss and W. Porod, Phys. Rev. D **92**, 055019 (2015); Y. Omura, K. Tobe, and K. Tsumura, Phys. Rev. D **92**, 055015 (2015); W. Chao, Phys. Rev. B **753**, 117 (2016); L. A. Anchordoqui, I. Antoniadis, H. Goldberg, X. Huang, D. Lust, and T. R. Taylor, Phys. Lett. B **749**, 484 (2015); H. M. Lee, D. Kim, K. Kong, and S. C. Park, J. High Energy Phys. 11 (2015) 150.
[28] A. Alboteanu, W. Kilian, and J. Reuter, J. High Energy Phys. 11 (2008) 010; W. Kilian, T. Ohl, J. Reuter, and M. Sekulla, Phys. Rev. D **91**, 096007 (2015).
[29] J. Hisano, N. Nagata, and Y. Omura, Phys. Rev. D **92**, 055001 (2015); K. Cheung, W. Y. Keung, P. Y. Tseng, and T. C. Yuan, Phys. Lett. B **751**, 188 (2015); Y. Gao, T. Ghosh, K. Sinha, and J. H. Yu, Phys. Rev. D **92**, 055030 (2015); J. Yepes, arXiv:1507.03974; J. Yepes, R. Kunming, and J. Shu, arXiv:1507.04745.