A scalar hint from the diboson excess?

Giacomo Cacciapaglia\textsuperscript{a},\textsuperscript{b}, Aldo Deandrea\textsuperscript{a},\textsuperscript{b}

\textsuperscript{a}Université de Lyon, France; Université Lyon 1, Villeurbanne, France; CNRS/IN2P3, UMR5822, IPNL F-69622 Villeurbanne Cedex, France

\textsuperscript{b}Institut Universitaire de France, 103 boulevard Saint-Michel, 75005 Paris, France

Michio Hashimoto

Chubu University, 1200 Matsumoto-cho, Kasugai-shi, Aichi, 487-8501, Japan

(Dated: October 22, 2015)

PACS numbers: 11.15.Ex, 12.60.Rc, 14.80.Va

In view of the recent diboson resonant excesses reported by both ATLAS and CMS Collaborations, we suggest that a new weak singlet pseudo-scalar particle, $\eta_{WZ}$, may decay into two weak bosons while being produced in gluon fusion at the LHC. The couplings to the gauge bosons can arise from a Wess-Zumino-Witten anomaly term and thus we study an effective model based on the anomaly term as a well motivated phenomenological model. In models where the pseudo-scalar arises as a composite state, the coefficients of the anomalous couplings can be related to the fermion components of the underlying dynamics. We provide an example to test the feasibility of the idea.

Introduction.— Recently, the ATLAS search for resonant diboson (WW/ZZ/WZ) productions in fully hadronic final states has been reported and a discrepancy with the background-only model having 3σ significance is observed around 2 TeV \cite{1}. A similar analysis has been also performed by the CMS Collaboration \cite{2}, where moderate excesses, less than 2σ significance, are found in the same mass range $\sim$ 2TeV. In the semi-leptonic decay channel of the diboson resonances, however, there seems no excess \cite{3} \cite{4}. Also, in the fully leptonic decay channel of the $WZ$ resonance, no significant deviation from the Standard Model (SM) prediction is observed \cite{5} \cite{6}. Even if the excesses in the fully hadronic decay channels are not just a fluctuation, there is a possibility that a new resonance around 2 TeV contributes to some of the channels, while others are populated by misidentification of the boson-tagged jet: for instance, one may have a neutral resonance that only couples to the $WW$ and $ZZ$ channels, so that the excess in the other $WZ$ channel is just a contamination owing to uncertainties of the tagging selections. Indeed the situation is still unclear, but it is worth considering the possibility for the explanation of the 3σ discrepancy by the ATLAS, which is statistically most significant, with some new physics effects.

This kind of exploratory exercise allows us to consider different new physics hypotheses and to evaluate the theoretical motivations for various models behind such an experimental hint. Several authors suggest that these excesses may be due to the existence of a new vector resonance, such as a composite $p_T$ or a weakly coupled $Z'$ and $W'$, and/or some other effect \cite{7} \cite{22}. Recently, unitarity bounds for the picture describing the excess with new vector resonances were also discussed \cite{23} \cite{24}. A detailed discussion of the population of all diboson channels via misidentified jets can be found in Ref. \cite{18}. We want to point out a novel possibility related to the existence of a spin-0 resonance whose couplings match the observed excesses while other couplings are naturally absent, thus giving a well-motivated phenomenological model. The case of a scalar coupling to gauge bosons via higher dimensional operators is discussed in \cite{22}. We suggest instead that a new weak singlet pseudo-scalar particle $\eta_{WZ}$, which couples to gauge bosons only via a Lagrangian term inspired by the Wess-Zumino-Witten (WZW) anomaly \cite{26} \cite{28}, can decay into two weak bosons (WW/ZZ channels) while being produced via gluon fusion at the LHC. This case is theoretically well motivated especially in scenarios of multi-TeV scale strong dynamics, where such states arise as a massive scalar associated to an anomalous global symmetry of the confining dynamics. In the following, we will use it as a guiding line, allowing the precise values of the couplings to vary in order to explore a larger class of models. We will rely on the composite scenario to establish to what extent the couplings can be considered “natural”, in the sense of the order of magnitude in the effective theory when particular assumptions on the specific underlying model are considered. We assume that this pseudo-scalar $\eta_{WZ}$ is fermiophobic, i.e., its couplings to the SM fermions, in particular to the top quark, are vanishing or tiny. In detailed composite models this is a realistic possibility and such particles with these properties are typically present \cite{29}. Moreover, new scalar resonances are typically expected to be lighter than their vector counterparts\footnote{It is then easier to evade the S-parameter \cite{30} constraint, because one can push up the scale of a new strong dynamics.} so it is quite reasonable to see such a pseudo-scalar particle at a lower mass than new vectors, as in QCD. It is, however, required to construct an appropriate composite model beyond an adaptation from existing...
models, because, for example, familiar dynamical models do not have the $WW$ coupling \[31\].

An obvious question is how to enhance the production cross section of $\eta_{WZ}$, which is typically expected to be tiny. Notice that the color factor $N_c = 3$ counts strongly in $\pi^0 \to \gamma\gamma$. Similarly, such an enhancement factor can arise from the degrees of freedom of the constituent fermions of $\eta_{WZ}$, which emerges via the underlying strong dynamics. It is therefore interesting to consider this hypothesis and use data from the diboson excess to have an idea of the coefficients required to explain such an excess, which in turn can hint at the more fundamental structure beyond the effective model.

We therefore assume that $\eta_{WZ}$ couples to gluons and the weak bosons, as in the case of the anomaly, and does couple not (or only very weakly) to the SM quarks and leptons. Before considering detailed numbers, it is useful to put rough numbers on the model: the excess points to an effective diboson cross section of about 10 fb. This implies that the production cross section $\sigma(gg \to \eta_{WZ})$ should be around 100 fb in order to explain the diboson excesses as the Branching Ratio to the $WW$ channel is roughly $\text{Br}(\eta_{WZ} \to W^+W^-) \sim 2(N_c\alpha^2)/80^2 \sim 10\%$ when considering an anomaly induced coupling. For the $ZZ$ mode, we expect a half of it. Thus the desired situation, $\text{Br}(\eta_{WZ} \to W^+W^-/ZZ) \sim 10\%$, can in principle be achieved. In this scenario, we regard the $WZ$ channel maximizes a likelihood function in terms of the truth signal in the $WW/ZZ/WZ$ channels \[18\]. The total width should not be so large, which is constrained less than, say, 100 GeV, owing to the one-loop effects essentially.

One might also worry about the constraint from the diphoton resonance searches, because the pseudo-scalar can also decay into a pair of photons, unlike the vector resonance. The constraints of the diphoton channel for a spin 0 resonance have been studied in the mass ranges from 150 GeV to 850 GeV by the CMS Collaboration \[32\]. A similar analysis was also performed by the ATLAS collaboration \[33\]. For the high-mass diphoton resonances, the CMS Collaborations found the constraint of the production cross section times branching ratio less than 0.3 fb for the $2\text{TeV RS graviton}$ \[34\]. The ATLAS Collaborations also performed a similar analysis and the expected $+2\sigma$ variation limit is 0.5 fb \[35\]. Even if we take the same bound for the spin 0 particle, our suggested explanation is fairly safe against this constraint in any case, because $\Gamma(\eta_{WZ} \to \gamma\gamma)/\Gamma(\eta_{WZ} \to W^+W^-) \sim \alpha^2/2\alpha^2 \simeq 0.03$. The decay channel of $\eta_{WZ} \to Zh$ is potentially dangerous \[36\] \[37\], if one might expect a similar situation to the two Higgs doublet model, in which there appears $A \to Zh$ at the tree level. In our case however, we can safely assume that the mixing between the singlet and the Higgs doublet is absent. The possible constraint from final states with Higgs boson(s) is thus easily avoided.

In the rest of the Letter, we will present a general effective description of the model, and a specific dynamical model that may give rise to the desired couplings. Finally we will study in a model independent way the constraint on the couplings, necessary to reproduce the observed diboson excess.

**Effective Lagrangian.** — The action for a weak singlet pseudo-scalar $\eta_{WZ}$ with no hypercharge is

$$ S_\eta = \int d^4x \frac{1}{2} (\partial \mu \eta_{WZ} \partial^\mu \eta_{WZ} - M_\eta^2 \eta_{WZ}^2) + \Gamma_{WZW} , \quad (1) $$

where $M_\eta$ is the mass of the pseudo-scalar singlet and the WZW term contains the effective Lagrangian for the diboson decay, $\Gamma_{WZW} \supset \int d^4x \mathcal{L}_{\eta V V}$, with

$$ \mathcal{L}_{\eta g g} = \kappa_\eta^2 \frac{g_3^2}{32\pi^2} \frac{\eta_{WZ}}{F_\eta} \epsilon^{\mu \nu \rho \sigma} C^a_{\mu \nu} G^a_{\rho \sigma} , \quad (2) $$

$$ \mathcal{L}_{\eta W W} = \kappa_W^4 \frac{g_2^2}{32\pi^2} \frac{\eta_{WZ}}{F_\eta} \epsilon^{\mu \nu \rho \sigma} W^i_{\mu \nu} W^i_{\rho \sigma} , \quad (3) $$

$$ \mathcal{L}_{\eta B B} = \kappa_B^4 \frac{g_V^2}{32\pi^2} \frac{\eta_{WZ}}{F_\eta} \epsilon^{\mu \nu \rho \sigma} B_{\mu \nu} B_{\rho \sigma} , \quad (4) $$

where $F_\eta$ denotes the decay constant of $\eta_{WZ}$, and the couplings $\kappa_\eta^2$, $\kappa_W^4$, and $\kappa_B^4$ are arbitrary prefactors in the effective description, but they can be calculated in specific realizations when the content of the loop terms is calculated, as we shall see in the next section. The couplings $g_3$, $g_2$, and $g_V$ are, respectively, the gauge coupling constants of the strong, weak, and hypercharge groups.

From the previous Lagrangian, the partial widths in the various channels can be easily calculated:

$$ \Gamma(\eta_{WZ} \to gg) = \frac{g_3^4(\kappa_\eta^2)^2 M_\eta^3}{128 F_\eta^2 \pi^5} , \quad (5) $$

$$ \Gamma(\eta_{WZ} \to WW) = \frac{g_2^4(\kappa_W^4)^2 (M_W^2 - 4 M_Z^2)^2}{512 F_\eta^2 \pi^5} , \quad (6) $$

$$ \Gamma(\eta_{WZ} \to ZZ) = \frac{g_V^4 (\kappa_W^4 + \kappa_B^4)^2 (M_Z^2 - 4 M_Z^2)^2}{1024 F_\eta^2 \pi^5} , \quad (7) $$

$$ \Gamma(\eta_{WZ} \to Zh) = \frac{e^2 g_2^2 (\kappa_W - \kappa_B^2)^2 M_Z^3}{512 F_\eta^2 \pi^5 M_\eta^3} , \quad (8) $$

$$ \Gamma(\eta_{WZ} \to \gamma\gamma) = \frac{e^4 (\kappa_W^2 + \kappa_B^2)^2 M_\eta^3}{1024 F_\eta^2 \pi^5} , \quad (9) $$

$^2$ A term in the form $\mathcal{L}_{\eta W B} = \kappa_{W B}^4 \frac{g_{23}^2}{32\pi^2} \frac{\eta_{WZ}}{F_\eta} \epsilon^{\mu \nu \rho \sigma} W^i_{\mu \nu} B_{\rho \sigma}$ can appear through the EWSB effects, however its coefficient is expected to be suppressed by a $v^2/F_\eta^2$ factor as it violates gauge invariance.
where \( c_W \equiv \cos \theta_W \), \( t_W \equiv \tan \theta_W \), \( e = g_2 \sin \theta_W \) with \( \theta_W \) being the weak mixing angle, and we define \( \kappa^q = \kappa^q_W + \kappa^q_B \) for future reference. A naive counting of the coupling constants and of the numerical prefactors immediately shows that the production and decay to gluons will be the dominant channel, but the values of the couplings \( \kappa^q \) and \( \kappa^q_W \) will play a major role in the phenomenological results.

This effective model allows us to easily calculate the diboson rates at the LHC, and check other constraints on the model. Before showing the numerical results, in the next section we will introduce a simple model of underlying dynamics that may lead to the required phenomenology.

A Vector-like Model. — In order to discuss the expected phenomenology, we investigate in more detail the origin of the couplings \( \kappa^q \), \( \kappa_W^q \) and \( \kappa_B^q \). In the following we take a simple hypothesis of a vector-like model by giving the factors counting the fundamental particles in the anomaly loops. We do not discuss here the origin of the electroweak symmetry breaking, so that we assume that the SM-like Higgs boson with the mass being 125 GeV emerges as a composite object of the dynamics, or we incorporate it as an elementary particle.

Let us study as an example the vector-like model shown in Table I where \( SU(N) \) represents a strongly interacting gauge group. Such a dynamical model with the higher representations of the gauge group has been studied, for example, in Ref. [38]. Of course, we may take the fundamental representation as usual, if we allow arbitrary large \( N \). We introduce vector-like weak doublets \( Q \) and \( L \), with multiplicity \( n_Q \) and \( n_L \), respectively. The vector-like fermion \( N_{L,R} \) is a weak singlet. The total number of flavors is then \( N_f = 2N_c n_Q + 2n_L + 1 \), where \( N_c = 3 \) denotes the number of ordinary QCD colors. A large number of \( N_f \) is inappropriate, because the gauge theory loses asymptotic freedom when the fermion multiplicity is too large. At the one-loop level, asking for a negative coefficient of the \( \beta \) function [39], we obtain \( N_f < 11N/(4T(R)) \), where \( T(R) \) is the trace normaliza-
Taking ratios of the above bounds, we can extract direct bounds on the ratios of Branching Ratios:

\[
\frac{\text{Br}(\eta_{WZ} \rightarrow W^+W^-)}{\text{Br}(\eta_{WZ} \rightarrow gg)} > \frac{10}{200} = 0.05, \quad (15)
\]

\[
\frac{\text{Br}(\eta_{WZ} \rightarrow \gamma\gamma)}{\text{Br}(\eta_{WZ} \rightarrow W^+W^-)} < \frac{0.5}{10} = 0.05, \quad (16)
\]

which are easily satisfied in this model.

These simplified results clearly show that the fermiophobic pseudo-scalar with the anomalous interactions can explain the diboson excesses without conflict with the other experimental bounds we discussed. One has to keep in mind, however, that a detailed model built along these lines may require further scrutiny concerning other bounds, but such a detailed study is worth pursuing only if the present excess will be confirmed by the ongoing LHC run.

**Numerical results and discussion.**— In order to have more detailed numbers we have created a FeynRules \([43, 44]\) model and evaluated the cross sections, branching ratios and decay widths numerically using Madgraph \([45]\).

Using the following numerical values, \(n_Q = 1, n_L = 1, N = 2, N_e = 3\), which correspond to \(\kappa_f^i = 2\) and \(\kappa_g^i = \kappa_W^i = 4\), and \(F_\eta = 1\ TeV\), the production cross section of the \(\eta_{WZ}\) particle is 0.615 fb and its total width 1.12 GeV at LHC with 8 TeV of center of mass energy for a \(\eta_{WZ}\) particle of 2 TeV of mass.

Using instead \(N = 5\) and all the other same parameters as in the previous example, increases the couplings by a factor of 10: \(\kappa_f^i = 20\) and \(\kappa_g^i = \kappa_W^i = 40\), while the production cross section and width of the \(\eta_{WZ}\) particle are a factor of 100 larger as expected (production cross section of 61.5 fb and total width of 112 GeV). The results for the branching fractions are given in Table II. These number are just indications based on a particular choice of parameters. One can see easily from the previous results that increasing \(N\) (or decreasing \(F_\eta\)) will increase the cross section and allow reaching a value compatible with the excess.

We consider in the following the parameters \(\kappa_f^i\) in order to describe and bound the model in an effective way without reference to a particular underlying model. First, we can impose bounds on the couplings by taking ratios of

| decay mode | BR |
|------------|----|
| \(gg\)     | 83%|
| \(WW\)     | 11.2%|
| \(ZZ\)     | 3.2%|
| \(Z\gamma\) | 2% |
| \(\gamma\gamma\) | 0.4% |

**TABLE II.** Decay modes and branching fraction of the \(\eta_{WZ}\) particle of 2 TeV with \(\kappa_f^i/\kappa_g^i = 2\).

# FIG. 1. Cross section times branching ratios on the \(\kappa_f^i-\kappa_W^i/\kappa_g^i\) plane for \(F_\eta = 1\ TeV\) and \(\kappa_B^i = 0\). The shaded region in the right upper area is excluded owing to \(\sigma(gg \rightarrow \eta_{WZ}) \cdot \text{Br}(\eta_{WZ} \rightarrow \gamma\gamma) > 0.5\ fb\). The numbers \(N = 4, 5, 6\) represent the corresponding values for the vector-like model with \(n_Q = n_L = 1\).

Branching Ratios and compare them with the bounds detailed in the previous section on the diboson, dijet, and diphoton resonant cross sections. Taking ratios of formulas (5)–(9), we can eliminate the dependency on the cross section, and derive bounds on the couplings \(\kappa_f^i\):

\[
\frac{(\kappa_f^\eta)^2}{(\kappa_g^\eta)^2} > \frac{1}{5} \frac{g_3^2}{g_2^2} \sim 1.45, \quad (17)
\]

\[
\frac{(\kappa_f^\gamma)^2}{(\kappa_W^\gamma)^2} < \frac{1}{e^4} \frac{g_2^4}{\sin^2 \theta_W} \sim 1.86, \quad (18)
\]

where \(g_3 = 1.033, g_2 = 0.628, \) and \(\sin^2 \theta_W = 0.2319\) at an energy of 2 TeV.

To compute constraints in the \(\kappa_f^\eta-\kappa_W^\eta/\kappa_g^\eta\) plane, we need an expression for the cross section:

\[
\sigma(gg \rightarrow \eta_{WZ}) = \frac{(\kappa_g^\eta)^2}{2} \frac{(1\ TeV)^2}{F_\eta^2} \cdot 0.615 \ fb \quad (19)
\]

which can be estimated by rescaling our numerical results. For the \(\text{Br}(\eta_{WZ} \rightarrow gg)\) and \(\text{Br}(\eta_{WZ} \rightarrow WW)\), by using Eqs. (5)–(9) for \(\kappa_B^i = 0\), we have

\[
\text{Br}(\eta_{WZ} \rightarrow gg) \simeq \frac{8g_2^4(\kappa_g^\eta)^2}{8g_2^4(\kappa_g^\eta)^2 + 3g_2^4(\kappa_W^\eta)^2}, \quad (20)
\]

and

\[
\text{Br}(\eta_{WZ} \rightarrow WW) \simeq \frac{2g_2^4(\kappa_W^\eta)^2}{8g_2^4(\kappa_g^\eta)^2 + 3g_2^4(\kappa_W^\eta)^2}, \quad (21)
\]

respectively. These estimates can change if we introduce \(\kappa_B^i, \kappa_{WW}^i \neq 0\) in general. In Fig. 1 we show the dijet \(\sigma_{jj}\) and diboson \(\sigma_{WW}\) cross sections for \(F_\eta = 1\ TeV\) as a function of the \(\kappa_f^\eta\) and the ratio \(\kappa_W^\eta/\kappa_g^\eta\) (for \(\kappa_B^i = 0\)). We also show the model predictions for \(N = 4, 5, 6\). In
the case of \( N = 7 \), \( \sigma_{jjj} > 200 \text{ fb} \). The shaded region in Fig. [1] corresponds to \( \sigma(gg \rightarrow \eta_{WZ}) \cdot BR(\eta_{WZ} \rightarrow \gamma\gamma) > 0.5 \text{ fb} \). We then find that the model with \( N = 4 \) cannot explain the diboson excess, on the other hand, the case with \( N = 5 \) does. It is fairly safe for \( N = 6 \), although the diphoton production is slightly large. The Branching Ratios are depicted in Fig. [2].

In dynamical models, there usually appear many pseudos other than a singlet. How about the constraint of the color-octet pseudos? We estimate the difference \( \Delta M^2 \) of the mass squared by rescaling the electromagnetic mass splitting in the \( \pi^\pm - \pi^0 \) system [10].

\[
\frac{\Delta M^2}{m_{\pi^+}^2 - m_{\pi^0}^2} = \left( \frac{F_\pi}{f_\pi} \right)^2 \frac{\alpha_3(F_\eta)}{\alpha} \frac{3}{1},
\]

where the factor 3 is the color Casimir for the octets. We then find the mass of the color-octet pseudos as 3 TeV with \( F_\pi = 1 \text{ TeV} \), which is consistent with the lower mass bound of 2.7 TeV (2.5 TeV) by the ATLAS (CMS) Collaborations [42].

These results show that the possibility of a fermiophobic pseudo–scalar singlet coupling with anomaly type couplings to the gauge bosons can give an explanation of the diboson excess without requiring the artificial suppression of other channels such as \( \eta_{WZ} \rightarrow Zh \). The interpretation in terms of a more fundamental model, due to the large \( \kappa_\eta^0 \) and \( \kappa_W^0 \) couplings, requires relatively large representations as, for example, indicated in the vector-like model discussed in the previous section. This is not a problem in itself but the detailed model building requires some care in order to avoid other bounds from, for example, electroweak precision tests or the presence of other states which may be in the same mass range as the singlet \( \eta_{WZ} \). Although there is no \( W^+B \) mixing from the model construction, there may appear large contributions to the \( W^+W^- \) coupling such as \( W^+W^-\gamma \), for example. The motivation for a further more detailed analysis will depend on the confirmation or not of the present diboson excess in the near future.

Acknowledgements. — We thank the France-Japan Particle Physics Lab (TYL/FJPPL) for partial support. A.D. is partially supported by the “Institut Universitaire de France.” We also acknowledge partial support from the DéfiInphyNiTi-projet structurant TLF: the Labex-LIO (Lyon Institute of Origins) under Grant No. ANR-10-LABX-66 and FRAMA (FR3127, Fédération de Recherche “André Marie Ampère”).

[1] G. Aad et al. [ATLAS Collaboration], arXiv:1506.00962 [hep-ex].
[2] V. Khachatryan et al. [CMS Collaboration], JHEP 1408, 173 (2014).
[3] G. Aad et al. [ATLAS Collaboration], Eur. Phys. J. C 75, no. 5, 209 (2015).
[4] V. Khachatryan et al. [CMS Collaboration], JHEP 1408, 174 (2014).
[5] G. Aad et al. [ATLAS Collaboration], Phys. Lett. B 737, 223 (2014).
[6] V. Khachatryan et al. [CMS Collaboration], Phys. Lett. B 740, 83 (2015).
[7] H. S. Fukano, M. Kurachi, S. Matsuzaki, K. Terashi and K. Yamawaki, arXiv:1506.03751 [hep-ph].
[8] J. Hisano, N. Nagata and Y. Omura, arXiv:1506.03931 [hep-ph].
[9] D. B. Franzosi, M. T. Frandsen and F. Sannino, arXiv:1506.04392 [hep-ph].
[10] K. Cheung, W. Y. Keung, P. Y. Tseng and T. C. Yuan, arXiv:1506.06064 [hep-ph].
[11] B. A. Dobrescu and Z. Liu, arXiv:1506.06736 [hep-ph].
[12] J. A. Aguilar-Saavedra, arXiv:1506.06739 [hep-ph].
[13] A. Alves, A. Berlin, S. Profumo and F. S. Queiroz, arXiv:1506.06767 [hep-ph].
[14] Y. Gao, T. Ghosh, K. Sinha and J. H. Yu, arXiv:1506.07511 [hep-ph].
[15] A. Thamm, R. Torre and A. Wulzer, arXiv:1506.08688 [hep-ph].
[16] J. Brehmer, J. Hewett, J. Kopp, T. Rizzo and J. Tatter-sall, arXiv:1507.00013 [hep-ph].
[17] Q. H. Cao, B. Yan and D. M. Zhang, arXiv:1507.00268 [hep-ph].
[18] B. C. Allanach, B. Gripaios and D. Sutherland, arXiv:1507.01638 [hep-ph].
[19] T. Abe, T. Kitahara and M. M. Nojiri, arXiv:1507.01681 [hep-ph].
[20] A. Carmona, A. Delgado, M. Quiros and J. Santiago, arXiv:1507.01914 [hep-ph].
[21] B. A. Dobrescu and Z. Liu, arXiv:1507.01923 [hep-ph].
[22] C. W. Chiang, H. Fukuda, K. Harigaya, M. Ibe and T. T. Yanagida, arXiv:1507.02483 [hep-ph].
[23] C. Englert, P. Harris, M. Spannowsky and M. Takeuchi, Phys. Rev. D 92 (2015) 013003.
[24] G. Cacciapaglia and M. T. Frandsen, arXiv:1507.00900 [hep-ph].
[25] T. Abe, R. Nagai, S. Okawa and M. Tanabashi,
[26] J. Wess and B. Zumino, Phys. Lett. B 37, 95 (1971).
[27] E. Witten, Nucl. Phys. B 223, 422 (1983).
[28] For an effective approach, see also, T. Fujiwara, T. Kugo, H. Terao, S. Uehara and K. Yamawaki, Prog. Theor. Phys. 73, 926 (1985); M. Hashimoto, Phys. Lett. B 381, 465 (1996); Phys. Rev. D 54, 5611 (1996).
[29] A. Arbey, G. Cacciapaglia, H. Cai, A. Deandrea, S. Le Corre and F. Sannino, arXiv:1502.04718 [hep-ph].
[30] M. E. Peskin and T. Takeuchi, Phys. Rev. Lett. 65, 964 (1990); Phys. Rev. D 46, 381 (1992).
[31] S. Dimopoulos, S. Raby and G. L. Kane, Nucl. Phys. B 182, 77 (1981); J. R. Ellis, M. K. Gaillard, D. V. Nanopoulos and P. Sikivie, Nucl. Phys. B 182, 529 (1981); For a recent work, see, e.g., R. S. Chivukula, P. Ittisamai, E. H. Simmons and J. Ren, Phys. Rev. D 84, 115025 (2011) [Phys. Rev. D 85, 119903 (2012)].
[32] V. Khachatryan et al. [CMS Collaboration], arXiv:1506.02301 [hep-ex].
[33] G. Aad et al. [ATLAS Collaboration], Phys. Rev. Lett. 113, no. 17, 171801 (2014).
[34] CMS Collaboration [CMS Collaboration], CMS-PAS-EXO-12-045.
[35] G. Aad et al. [ATLAS Collaboration], arXiv:1504.05511 [hep-ex].
[36] G. Aad et al. [ATLAS Collaboration], Eur. Phys. J. C 75, no. 6, 263 (2015).
[37] V. Khachatryan et al. [CMS Collaboration], arXiv:1506.01443 [hep-ex].
[38] D. K. Hong, S. D. H. Hsu and F. Sannino, Phys. Lett. B 597, 89 (2004).
[39] W. E. Caswell, Phys. Rev. Lett. 33, 244 (1974).
[40] D. D. Dietrich and F. Sannino, Phys. Rev. D 75 (2007) 085018.
[41] L. Del Debbio, PoS Lattice 2010 (2010) 004; T. Degrand, Y. Shamir and B. Svetitsky, Phys. Rev. D 88, no. 5, 054505 (2013).
[42] V. Khachatryan et al. [CMS Collaboration], Phys. Rev. D 91, no. 5, 052009 (2015). For a similar analysis by the ATLASC collaboration, see, G. Aad et al. [ATLAS Collaboration], Phys. Rev. D 91, no. 5, 052007 (2015).
[43] N. D. Christensen and C. Duhr, Comput. Phys. Commun. 180, 1614 (2009).
[44] A. Alloul, N. D. Christensen, C. Degrande, C. Duhr and B. Fuks, Comput. Phys. Commun. 185, 2250 (2014).
[45] J. Alwall, R. Frederix, S. Frixione, V. Hirschi, F. Maltoni, O. Mattelaer, H.-S. Shao and T. Stelzer et al., JHEP 1407, 079 (2014).
[46] E. Farhi and L. Susskind, Phys. Rept. 74, 277 (1981).