

Have we observed the Higgs boson (imposter)?Ian Low,^{1,2} Joseph Lykken,³ and Gabe Shaughnessy⁴¹*High Energy Physics Division, Argonne National Laboratory, Argonne, Illinois 60439, USA*²*Department of Physics and Astronomy, Northwestern University, Evanston, Illinois 60208, USA*³*Fermi National Accelerator Laboratory, P.O. Box 500, Batavia, Illinois 60510, USA*⁴*Department of Physics, University of Wisconsin, Madison, Wisconsin 53706, USA*

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We interpret the new particle at the Large Hadron Collider as a CP -even scalar and investigate its electroweak quantum number. Assuming an unbroken custodial invariance as suggested by precision electroweak measurements, only four possibilities are allowed if the scalar decays to pairs of gauge bosons, as exemplified by a dilaton/radion, a nondilaton electroweak singlet scalar, an electroweak doublet scalar, and electroweak triplet scalars. We show that current LHC data already strongly disfavor both the “plain-vanilla” dilaton and nondilaton singlet imposters. On the other hand, a generic Higgs doublet gives excellent fits to the measured event rates of the newly observed scalar resonance, while the Standard Model Higgs boson gives a slightly worse overall fit due to the lack of a signal in the $\tau\tau$ channel. The triplet imposter exhibits some tension with the data. The global fit indicates that the enhancement in the diphoton channel could be attributed to an enhanced partial decay width, while the production rates are consistent with the Standard Model expectations. We emphasize that more precise measurements of the ratio of event rates in the WW over ZZ channels, as well as the event rates in $b\bar{b}$ and $\tau\tau$ channels, are needed to further distinguish the Higgs doublet from the triplet imposter.

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I. INTRODUCTION

The new resonance discovered [1] by the ATLAS and CMS experiments at the CERN Large Hadron Collider (LHC) could be the long-sought Higgs boson of the Standard Model (SM) [2]. This is only the beginning of a challenging program of “Higgs identification” to rigorously establish the quantum numbers and couplings of the new particle, and to reveal its relationship, if any, to electroweak symmetry breaking and fermion mass generation.

To confirm the identity of the new particle, we should first establish what it is not. For example, the diphoton decay mode shows not only that the new state is a massive neutral boson but also that it does not have spin-one, which would violate the Landau-Yang theorem [3]. By studying angular correlations in the decays to four-lepton final states, it should be possible to distinguish whether the boson is CP -even, CP -odd, or a mixture [4,5], and eventually rule out the possibility that the boson has spin-two rather than spin-zero [5].

Here we will assume that the new particle is a CP -even scalar, and address the question of determining its electroweak quantum numbers. A Higgs boson is the CP -even neutral component of a complex weak doublet with unit hypercharge, with the other three states comprising the Goldstone bosons that become the longitudinal components of the W^\pm and Z bosons. Together these four states also transform as a $(\mathbf{2}_L, \mathbf{2}_R)$ under the accidental $SU(2)_L \times SU(2)_R$ global symmetry of the SM Lagrangian. After electroweak symmetry breaking, SM interactions still respect an approximate diagonal symmetry called the

custodial symmetry $SU(2)_C$ [6], as evidenced by precision electroweak measurements of the ρ parameter [7].

As shown in Ref. [8], we can classify the leading-order couplings of any neutral CP -even scalar to W and Z bosons according to its properties under custodial symmetry. There are five possibilities that could apply to the resonance discovered by ATLAS and CMS:

- (1) The scalar is an electroweak singlet (and thus also a custodial singlet), but has dimension-four couplings to W and Z . The latter property implies that the Higgs imposter is a dilaton [9,10] or radion [11] resulting from new electroweak symmetry breaking dynamics in a strongly interacting conformal sector or a warped extra dimension, the two being related by AdS/CFT duality [12]. The conformal dynamics couples the “dilaton imposter” χ to SM fermions, and to photons and gluons through operators of dimension five.
- (2) The scalar is an electroweak singlet with dimension-five couplings to W and Z . This is the electroweak singlet imposter s discussed in Ref. [13]. The dimension-five couplings arise from integrating out other charged exotics, which also generically produce dimension-five couplings to photons and gluons, and higher-dimension couplings to SM fermions. This “singlet imposter” could be related to electroweak symmetry breaking indirectly, either through an extended Higgs sector or as a “techni-axion” in technicolor models.
- (3) The scalar is not an electroweak singlet, but is nevertheless a custodial singlet. This could be the

Higgs boson h , which is the custodial singlet component in the decomposition $(\mathbf{2}_L, \mathbf{2}_R) = \mathbf{1} \oplus \mathbf{3}$. We will refer to this possibility simply as the ‘‘Higgs boson,’’ although it could very well be a custodial singlet in a more exotic representation of $SU(2)_L \times SU(2)_R$.

- (4) The scalar is the neutral member of a custodial 5-plet. This imposter could belong to an electroweak triplet in an extended Higgs sector [14], and will be referred to as the ‘‘triplet imposter’’ h_5 .
- (5) Mixtures of the above are possible. However, note that, to the extent that mixtures (and thus mass eigenstates) respect custodial symmetry, the only plausible possibility that cuts across cases is a mixture of (1) and (3) [15] or a mixture of (2) and (3) [16], as might indeed occur in an extended Higgs sector. Mixtures that do not respect the custodial symmetry have been studied in Ref. [17].

There have been earlier works on fitting the couplings of a Higgs boson using the LHC 2011 data [18]. In this work we wish to focus on understanding the electroweak property of the observed resonance. For simplicity we consider only the pure cases (1)–(4), and demonstrate that a ‘‘plain-vanilla’’ dilaton imposter in case (1), where all the SM gauge bosons including gluons and the electroweak gauge bosons are part of the conformal dynamics, as well as the singlet imposter in case (2) are already strongly disfavored by LHC data probing scalar couplings with pairs of SM gauge bosons $V_1 V_2 = \{WW, ZZ, Z\gamma, \gamma\gamma, gg\}$. We will show that the custodial singlet Higgs in case (3) gives the best fit to current data and an SM Higgs boson, for which all couplings are fixed to the SM value, gives a slightly worse fit. The triplet imposter in case (4) exhibits some tension with the data, mainly due to the excess in $b\bar{b}$ and $\tau\tau$ channels.

II. SCALAR COUPLINGS TO $V_1 V_2$

As seen in Ref. [8], tree-level couplings to W and Z bosons of a scalar charged under electroweak symmetry can be classified using the quantum number of the scalar under the custodial symmetry $SU(2)_C$, which is the diagonal subgroup, after electroweak symmetry breaking, of an accidental $SU(2)_L \times SU(2)_R$ global symmetry of the SM Lagrangian. The approximate custodial invariance implies $\rho \equiv m_W^2/(m_Z^2 c_w^2) = 1$, where c_w is the cosine of the Weinberg angle, which was verified by the precision electroweak measurements to be true at the percent level [7].

The $SU(2)_L$ and the $U(1)_Y$ subgroup of $SU(2)_R$ are gauged in the SM, which implies that the weak isospin gauge bosons W_μ^a and the hypercharge gauge boson B_μ transform as a triplet and the T^3 component of a triplet, respectively, under $SU(2)_C$. Using the familiar rule for the addition of angular momentum in quantum mechanics, it is immediately clear that a pair of W/Z bosons can only couple to a CP -even neutral scalar that is either a custodial

singlet h or a custodial 5-plet h_5 [here both h and h_5 are charged under $SU(2)_L \times U(1)_Y$]. Any $(\mathbf{N}_L, \mathbf{N}_R)$ representation of $SU(2)_L \times SU(2)_R$ contains a custodial singlet for $\mathbf{N} \geq 2$ and also a custodial 5-plet for $\mathbf{N} \geq 3$. The usual Higgs doublet scalar is a $(\mathbf{2}_L, \mathbf{2}_R)$ representation, while the $(\mathbf{3}_L, \mathbf{3}_L) = \mathbf{1} \oplus \mathbf{3} \oplus \mathbf{5}$ representation contains a real triplet scalar with $Y = 2$ and a complex triplet scalar with $Y = 0$.

We parametrize effective couplings of h and h_5 to $V_1 V_2$ as

$$\begin{aligned} \mathcal{L}_{hV_1V_2} = & c_V \left(\frac{2m_W^2}{v} h W_\mu^+ W^{-\mu} + \frac{m_Z^2}{v} h Z_\mu Z^\mu \right) \\ & + c_g \frac{\alpha_s}{12\pi v} h G_{\mu\nu}^a G^{a\mu\nu} + c_\gamma \frac{\alpha}{8\pi v} h F_{\mu\nu} F^{\mu\nu} \\ & + c_{Z\gamma} \frac{\alpha}{8\pi v s_w} h F_{\mu\nu} Z^{\mu\nu}, \end{aligned} \quad (1)$$

$$\begin{aligned} \mathcal{L}_{h_5V_1V_2} = & c_{5V} \left(-\frac{m_W^2}{v} h_5 W_\mu^+ W^{-\mu} + \frac{m_Z^2}{v} h_5 Z_\mu Z^\mu \right) \\ & + c_{5g} \frac{\alpha_s}{12\pi v} h_5 G_{\mu\nu}^a G^{a\mu\nu} + c_{5\gamma} \frac{\alpha}{8\pi v} h_5 F_{\mu\nu} F^{\mu\nu} \\ & + c_{5Z\gamma} \frac{\alpha}{8\pi v s_w} h_5 F_{\mu\nu} Z^{\mu\nu}, \end{aligned} \quad (2)$$

where $v \approx 246$ GeV. The first lines in Eqs. (1) and (2) contain couplings to pairs of massive electroweak gauge bosons, which could arise at the tree level, while the second lines include couplings to massless gauge bosons (including the $Z\gamma$ channel), which only occur at the one-loop level. Notice that the ratios of couplings to WW over ZZ for the custodial singlet Higgs and the triplet imposter are different [8]:

$$\frac{g_{hWW}}{g_{hZZ}} = \frac{m_W^2}{m_Z^2} = c_w^2, \quad \frac{g_{h_5WW}}{g_{h_5ZZ}} = -\frac{m_W^2}{2m_Z^2} = -\frac{c_w^2}{2}. \quad (3)$$

Otherwise, they have similar coupling structure to $V_1 V_2$.

In the SM c_g , c_γ , and $c_{Z\gamma}$ are form factors which depend on the Higgs mass m_h , the top quark mass m_t , and the W boson mass m_W . More explicitly,

$$c_g^{(\text{SM})} = \frac{3}{4} A_{1/2}(\tau_t), \quad (4)$$

$$c_\gamma^{(\text{SM})} = A_1(\tau_W) + N_c Q_t^2 A_{1/2}(\tau_t), \quad (5)$$

$$\begin{aligned} c_{Z\gamma}^{(\text{SM})} = & c_w A_1(\tau_W, \lambda_W) \\ & + N_c \frac{Q_t (2T_3^{(t)} - 4Q_t s_w^2)}{c_w} A_{1/2}(\tau_W, \lambda_W), \end{aligned} \quad (6)$$

where $N_c = 3$ is the number of colors, Q_t is the top quark electric charge in units of $|e|$, $\tau_i = 4m_i^2/m_h^2$, and $\lambda_i = 4m_i^2/m_Z^2$. We use the same definitions of loop functions as in Ref. [19]. At 125 GeV, the numerical values are

$$c_g^{(\text{SM})}(125 \text{ GeV}) = 1, \quad c_\gamma^{(\text{SM})}(125 \text{ GeV}) = -6.48, \\ c_{Z\gamma}^{(\text{SM})}(125 \text{ GeV}) = 5.48. \quad (7)$$

More generally, these coefficients would depend on the masses of new particles contributing to the decay widths. However, for on-shell production of the Higgs at a fixed mass, it is a good approximation to regard these coefficients as constant.

As mentioned in the Introduction, it is possible to have a custodial singlet scalar that is also an electroweak singlet scalar, contrary to the Higgs boson h , which is charged under electroweak symmetry. For this possibility, the dilaton imposter χ turns out to have effective couplings to $V_1 V_2$ that are identical to the ordinary Higgs boson [10]. So we have

$$\mathcal{L}_{\chi V_1 V_2} = c_{\chi V} \left(\frac{2m_W^2}{v} \chi W_\mu^+ W^{-\mu} + \frac{m_Z^2}{v} \chi Z_\mu Z^\mu \right) \\ + c_{\chi s} \frac{\alpha_s}{12\pi v} \chi G_{\mu\nu}^a G^{a\mu\nu} + c_{\chi\gamma} \frac{\alpha}{8\pi v} \chi F_{\mu\nu} F^{\mu\nu} \\ + c_{\chi Z\gamma} \frac{\alpha}{8\pi v s_w} \chi F_{\mu\nu} Z^{\mu\nu}. \quad (8)$$

Moreover, in the plain-vanilla scenario where all of the SM gauge bosons are part of the conformal dynamics, the dilaton coupling to gauge bosons are determined entirely by the one-loop beta functions [10], which would then predict a dilaton-gluon-gluon coupling that is much enhanced over the SM values. We focus on the plain-vanilla dilaton in this work. In the other scenario, the singlet imposter s discussed in case (2) in the Introduction, leading-order couplings to all possible pairs of $V_1 V_2$ come from dimension-five operators and are induced only at the loop level. Three, and only three, gauge-invariant operators could be generated at this order:

$$\kappa_g \frac{\alpha_s}{4\pi} \frac{s}{4m_s} G_{\mu\nu}^a G^{a\mu\nu} + \kappa_W \frac{\alpha}{4\pi s_w^2} \frac{s}{4m_s} W_{\mu\nu}^a W^{a\mu\nu} \\ + \kappa_B \frac{\alpha}{4\pi c_w^2} \frac{s}{4m_s} B_{\mu\nu} B^{\mu\nu}. \quad (9)$$

At leading order these three operators determine the singlet coupling to all five pairs of SM gauge bosons, massive or not. In terms of mass eigenstates, the effective Lagrangian for couplings of a singlet imposter to SM gauge bosons is

$$\mathcal{L}_{s V_1 V_2} = \kappa_W \frac{\alpha}{8\pi m_s s_w^2} s W_{\mu\nu}^+ W^{-\mu\nu} + \left(\kappa_W \frac{c_w^2}{s_w^2} + \kappa_B \frac{s_w^2}{c_w^2} \right) \\ \times \frac{\alpha}{16\pi m_s} s Z_{\mu\nu} Z^{\mu\nu} + \kappa_g \frac{\alpha_s}{16\pi m_s} s G_{\mu\nu}^a G^{a\mu\nu} \\ + (\kappa_W + \kappa_B) \frac{\alpha}{16\pi m_s} s F_{\mu\nu} F^{\mu\nu} \\ + \left(\kappa_W \frac{c_w}{s_w} - \kappa_B \frac{s_w}{c_w} \right) \frac{\alpha}{8\pi m_s} s F_{\mu\nu} Z^{\mu\nu}, \quad (10)$$

from which we see that, generically, couplings to the massive and massless gauge bosons are of the same order of magnitude, unlike other cases we have considered so far where couplings to massive gauge bosons are tree-level and are the dominant decay channels. Expressions for the partial decay widths of the singlet scalar into SM gauge bosons can be found in Ref. [8]. From Eq. (10) it is also clear that, if there is any change in the decay width in the diphoton channel, the partial width in the $Z\gamma$ channel would be modified as well [19].¹

As already pointed out in Ref. [13], the democratic nature of a singlet imposter coupling to pairs of SM gauge bosons has important implications for phenomenology. First of all, the phase-space factor now plays an important role in its decay patterns. For example, the phase-space factor in the gg channel is a factor of 8 larger than that in the diphoton channel because of color. Below kinematic thresholds decays into massive gauge bosons like WW and ZZ are generically suppressed, which is the case for the mass range in which we are interested. Moreover, decays into all four pairs of electroweak gauge bosons, $\{WW, ZZ, \gamma\gamma, Z\gamma\}$, are correlated with one another, as they are controlled by only two parameters, κ_W and κ_B from Eq. (10). In sharp contrast, decays of h , h_5 , or χ into $\gamma\gamma$ and $Z\gamma$ are controlled by two free parameters in Eqs. (1) and (2), respectively, and are independent of the decays into WW and ZZ .

III. INTERPRETING THE DATA

So far data collected at the LHC show the greatest sensitivities and significances in decay channels into $V_1 V_2$, while there are also strong hints from decays into $b\bar{b}$ [20] and, to a lesser extent, $\tau\tau$ final states. Before we present our analyses, it is worth recalling that what is being measured experimentally is the event rate $B\sigma_X(Y)$ for a particular production mechanism X of the scalar $S = \{\chi, s, h, h_5\}$, which subsequently decays into final states Y :

$$B\sigma_X(Y) \equiv \sigma(X \rightarrow S) \frac{\Gamma(S \rightarrow Y)}{\Gamma_{\text{tot}}}, \quad (11)$$

where Γ_{tot} is the total width of S . For $V_1 V_2$ channels at the LHC, two different production mechanisms are considered in the current data—the gluon fusion $X = gg$ and the vector boson fusion (VBF) $X = \text{VBF}$ —while three decay channels to gauge bosons are measured: $\{WW, ZZ, \gamma\gamma\}$. We will denote inclusive production of the scalar by $X = pp$. The Tevatron $b\bar{b}$ result comes from the associated production of the Higgs with W/Z , $X = VH$. Experimental collaborations present their $B\sigma_X(Y)$ in units of the SM signal strength $B\sigma_X^{(\text{SM})}(Y)$ by defining a best-fit signal strength $\mu = B\sigma_X(Y)/B\sigma_X^{(\text{SM})}(Y)$. Given these notations,

¹This statement is true generically, regardless of the electroweak quantum number of the scalar.

we consider the following results from the most recent LHC and Tevatron announcements as well as the 2011 LHC data:

- (I) Inclusive channels
- (a) $B\sigma_{pp}(WW)$: $1.4^{+0.5}_{-0.5}$ (ATLAS) [21], $0.3^{+1.1}_{-0.3}$ (Tevatron) [20].
 - (b) $B\sigma_{pp}(ZZ)$: $1.1^{+0.6}_{-0.4}$ (ATLAS 7 TeV) [22], $0.7^{+0.5}_{-0.4}$ (CMS) [23].
 - (c) $B\sigma_{pp}(\gamma\gamma)$: $2.2^{+0.7}_{-0.8}$ (ATLAS 7 TeV) [22], $1.8^{+0.5}_{-0.8}$ (ATLAS 8 TeV) [24], $3.6^{+3.0}_{-2.5}$ (Tevatron) [20].
 - (d) $B\sigma_{pp}(\tau\tau)$: $0.5^{+1.6}_{-2.1}$ (ATLAS 7 TeV) [22].
- (II) Exclusive channels
- (a) $B\sigma_{\text{Non-VBF}}(\gamma\gamma)$: $1.7^{+1.1}_{-1.1}$ (ATLAS) [25], $1.4^{+0.6}_{-0.6}$ (CMS) [23].
 - (b) $B\sigma_{\text{VBF}}(\gamma\gamma)$: $2.8^{+3}_{-2.3}$ [25], $2.2^{+1.3}_{-1.1}$ (CMS) [23].
 - (c) $B\sigma_{\text{Non-VBF}}(WW)$: $0.7^{+0.5}_{-0.5}$ (CMS) [23].
 - (d) $B\sigma_{\text{VBF}}(WW)$: $0.3^{+1.5}_{-1.6}$ (CMS) [23].
 - (e) $B\sigma_{\text{Non-VBF}}(\tau\tau)$: $1.3^{+1.1}_{-1.1}$ (CMS) [26].
 - (f) $B\sigma_{\text{VBF}}(\tau\tau)$: $-1.8^{+1.0}_{-1.0}$ (CMS) [26].
 - (g) $B\sigma_{VH}(b\bar{b})$: $0.5^{+2.1}_{-2.2}$ (ATLAS 7 TeV) [22], $0.5^{+0.8}_{-0.8}$ (CMS) [23], $2.0^{+0.7}_{-0.7}$ (Tevatron) [20].

Unless otherwise stated, the LHC results assume combinations of the $\sqrt{s} = 7$ and 8 TeV data sets. ATLAS only provides results in inclusive channels, with the exception of the $b\bar{b}$ channel. While CMS provides both inclusive and exclusive results, we only use the exclusive results in the fit so as to avoid double counting.

In the absence of any information on the total width of the resonance, we could proceed in a model-independent fashion by taking the ratios of event rates, so that the total width cancels in the ratio. On the other hand, if we make assumptions on the total width of the scalar, it is possible to fit the event rate itself, although the outcome is clearly model-dependent.

Taking ratios of event rates has the advantage that some of the common uncertainties, such as systematics and theoretical error in production cross section, should cancel [27]. In addition, modifications in properties of the scalar that are universal in all decay channels would drop out in the ratio. Two examples are i) mixing with other scalars that have not been observed to date, and ii) higher-dimensional operators giving additional contributions to the scalar kinetic term and resulting in a finite wave function renormalization of the scalar.² The drawback of taking the ratio, on the other hand, is that we may not have information on the overall normalization of the parameters in the effective Lagrangian.

A. Model-independent fits in V_1V_2 channels

We focus on taking the ratios of event rates in diboson channels, since these provide useful discriminators among

²This is the effect of c_H in the strongly interacting light Higgs Lagrangian [28].

different Higgs imposters. Two classes of ratios could be taken:

- (i) Ratios of event rates with the same production mechanism but different decay channels. In this class we consider

$$D_{W/Z} \equiv \frac{B\sigma_{gg}(WW)}{B\sigma_{gg}(ZZ)} = \frac{\Gamma(S \rightarrow WW)}{\Gamma(S \rightarrow ZZ)}, \quad (12)$$

$$D_{\gamma/Z} \equiv \frac{B\sigma_{gg}(\gamma\gamma)}{B\sigma_{gg}(ZZ)} = \frac{\Gamma(S \rightarrow \gamma\gamma)}{\Gamma(S \rightarrow ZZ)}, \quad (13)$$

$$D_{Z\gamma/Z} \equiv \frac{B\sigma_{gg}(Z\gamma)}{B\sigma_{gg}(ZZ)} = \frac{\Gamma(S \rightarrow Z\gamma)}{\Gamma(S \rightarrow ZZ)}. \quad (14)$$

The first two ratios can be extracted from existing data, while the $Z\gamma$ decay channel has been suggested [29], but not reported. Since ATLAS did not report the exclusive channel, we use the number from the inclusive channel as an approximation. It is well-known that in the SM the inclusive rate is dominated by the gg channel, with the VBF channel making up only about 7% of the inclusive rate [30]: we include the relative weights of the gg and VBF production mechanisms when considering inclusive rates.

- (ii) Ratios of event rates with different production mechanisms but the same decay channel. Since at the LHC the dominant production mechanisms are the gg channel and, to a much lesser extent, the VBF channel, we only consider one ratio in this class:

$$P_{g/V} \equiv \frac{B\sigma_{gg}(\gamma\gamma)}{B\sigma_{\text{VBF}}(\gamma\gamma)} = \frac{\sigma(gg \rightarrow S)}{\sigma(\text{VBF} \rightarrow S)}. \quad (15)$$

When more data becomes available it will also be useful to form this ratio for the other three diboson channels.

For a 125 GeV singlet imposter the decays into all four pairs of electroweak gauge bosons are controlled by only two free parameters, κ_W and κ_B in Eq. (10). Therefore the three ratios in the first class depend only on one number: κ_W/κ_B . In Fig. 1(a) we show the χ^2 of using one parameter, κ_W/κ_B , to fit the measured $D_{\gamma/Z}$ and $D_{W/Z}$ from ATLAS and CMS at the same time. We see that the best-fit value is

$$\frac{\kappa_W}{\kappa_B} \approx -1, \quad (16)$$

and the absolute χ^2 is below the 95% C.L. limit. Using the above value, the predicted ratio of $D_{Z\gamma/Z}$ is

$$D_{Z\gamma/Z} \sim 500, \quad (17)$$

which would be a spectacular signal. Although a dedicated search for a resonance in the $Z\gamma$ channel has not been reported, measurements for SM diboson production in the $Z\gamma$ channel have been made. Resonance decays in the $Z\gamma$

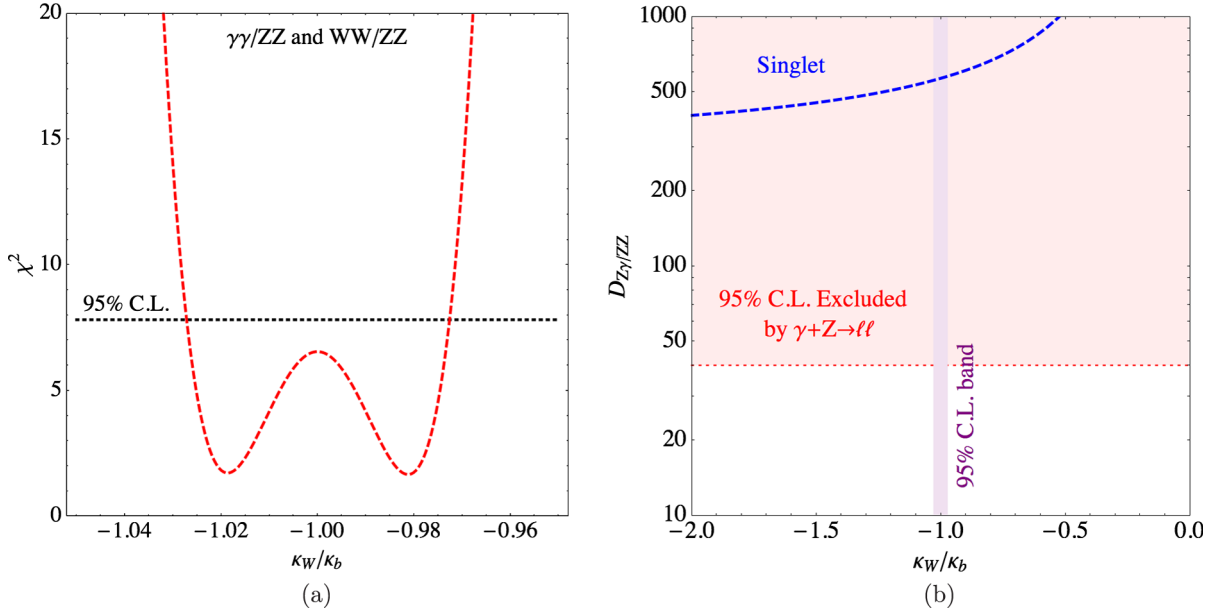


FIG. 1 (color online). (a) χ^2 from fitting $D_{\gamma/Z}$ and $D_{W/Z}$ using one single parameter κ_W/κ_B , which is above the 95% C.L. limit. (b) The predicted $D_{Z\gamma/ZZ}$ using current data. The 95% C.L. exclusion limit is derived from measurements of SM diboson production in the $Z\gamma$ channel, while the 95% C.L. band for κ_W/κ_B is derived from comparing $\Delta\chi^2$ with the best-fit value in (a).

channel with a much enhanced rate certainly would contribute to this set of measurements as well. In Refs. [31,32] the event rates of $\sigma(pp \rightarrow Z\gamma + X) \times \text{Br}(Z \rightarrow \ell^+ \ell^-)$ are measured to be consistent with that expected from the SM prediction:

$$\begin{aligned} \text{ATLAS: } & 6.5 \pm 1.2(\text{stat}) \pm 1.7(\text{syst}) \pm 0.2(\text{lumi}) \text{ pb,} \\ & \text{Theory: } 6.9 \pm 0.5 \text{ pb,} \\ \text{CMS: } & 9.4 \pm 1.0(\text{stat}) \pm 0.6(\text{syst}) \pm 0.4(\text{lumi}) \text{ pb,} \\ & \text{Theory: } 9.6 \pm 0.4 \text{ pb.} \end{aligned}$$

The different values for ATLAS and CMS result from different selection cuts. On the other hand, using the best-fit signal strength for $B\sigma_{pp}(ZZ)$ at the LHC, we see that the predicted $\sigma(pp \rightarrow s \rightarrow Z\gamma + X) \times \text{Br}(Z \rightarrow \ell^+ \ell^-) \sim 15$ pb. Although we have not simulated the selection efficiency of the resonance decays into $Z\gamma$ for the cuts imposed in Refs. [31,32], it is worth noting that the p_T distribution of the photon from resonance decays is peaked at $m_s^2 - m_Z^2/(2m_s) \approx 30$ GeV, while that from the SM diboson production is peaked at $p_T = 0$ [29]. Therefore, we expect a significant amount of the events from the resonance decay to pass the photon p_T cut. In the end, we see that an event rate of the order of 15 pb in the resonance decays into $Z + \gamma \rightarrow \ell^+ \ell^- + \gamma$ is strongly disfavored. Using these arguments we derive an estimate of the 95% C.L. limit on $D_{Z\gamma/Z}$, using the measured $B\sigma_{pp}(ZZ)$, which is shown in Fig. 1(b). We see that the predicted $D_{Z\gamma/Z}$ from a singlet imposter is an order of magnitude larger than the 95% C.L. limit. Therefore,

Fig. 1 shows that a singlet imposter is excluded at 95% C.L. as the interpretation of the excess at the LHC.

It is possible to understand why the partial width in the $Z\gamma$ channel is enhanced by so much for the singlet imposter. As mentioned in the end of Sec. II, its couplings to gauge bosons are democratic and the partial width is largely determined by phase factors and kinematics. Therefore at 125 GeV, partial widths of s has the following generic feature [13]:

$$\Gamma_{gg} \gtrsim \Gamma_{\gamma\gamma} \gtrsim \Gamma_{Z\gamma} \gtrsim \Gamma_{WW} \gtrsim \Gamma_{ZZ}. \quad (18)$$

However, in the SM we have

$$\Gamma_{WW}^{(\text{SM})} > \Gamma_{gg}^{(\text{SM})} > \Gamma_{ZZ}^{(\text{SM})} > \Gamma_{\gamma\gamma}^{(\text{SM})} > \Gamma_{Z\gamma}^{(\text{SM})}, \quad (19)$$

and current measurements suggest a diphoton partial width that is still smaller than those in the WW and ZZ channels. Therefore the diphoton decay width of a singlet imposter should be suppressed from its generic expectation in order to fit the measured event rate. In Eq. (10) the s - γ - γ coupling is controlled by $\kappa_W + \kappa_B$, which explains why the best-fit value is $\kappa_W/\kappa_B \approx -1$. In this region we see from Eq. (10) that there is also a partial cancellation in the s - Z - Z coupling, while the s - Z - γ coupling is enhanced. Together with the fact that at 125 GeV the ZZ final state is below the kinetic threshold, it is not surprising that the predicted $Z\gamma$ partial width is much larger than the ZZ partial width.

For a 125 GeV custodial singlet and 5-plet, $D_{W/Z}$ is completely fixed to be

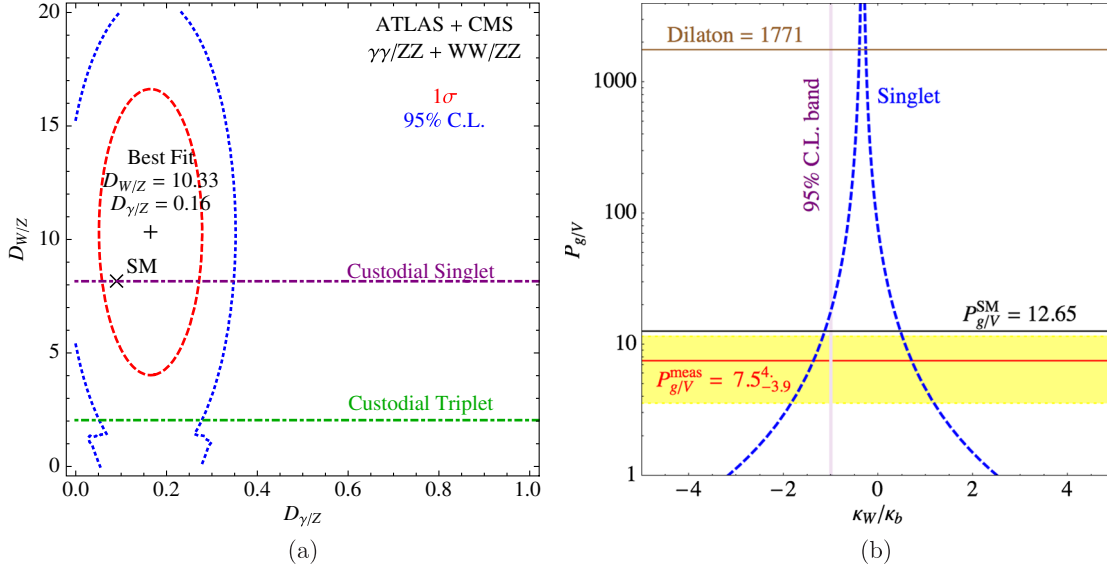


FIG. 2 (color online). (a) Predictions of the Higgs boson and the triplet imposter on the $D_{W/Z}$ - $D_{\gamma/Z}$ plane. The custodial singlet is within the 1σ range while the custodial triplet is within the 2σ range of the measured value. The feature at the bottom of the 2σ contour is due to asymmetric uncertainties. (b) The predicted $P_{g/V}$ for the dilaton and the singlet imposters. The dilaton imposter predicts a $P_{g/V}$ that is strongly disfavored.

$$D_{W/Z}^{(h)} = 8.16, \quad D_{W/Z}^{(h_5)} = \frac{1}{4} D_{W/Z}^{(h)} = 2.04. \quad (20)$$

Thus a large deviation of $D_{W/Z}$ from these two values would disfavor the custodial singlet Higgs and the triplet imposter as the interpretation of the excess.³ The ratio $D_{\gamma/Z}$ also allows for an estimate of the ratios c_γ/c_V and $c_{5\gamma}/c_{5V}$. There is no prediction on the $D_{Z\gamma/Z}$ in these two scenarios, although simultaneous measurements of $D_{\gamma/Z}$ and $D_{Z\gamma/Z}$ may shed light on electroweak properties of new light degrees of freedom mediating scalar decays in the $\gamma\gamma$ and $Z\gamma$ channel [19]. In Fig. 2(a) we show the ratios extracted from the LHC data on the $D_{W/Z}$ - $D_{\gamma/Z}$ plane, as well as the expectations for the custodial singlet Higgs and the triplet imposter. We see that the custodial singlet Higgs is consistent with data within the 1σ contour, while the triplet imposter is consistent within the 95% C.L. limit.

It should be emphasized that $\{WW, ZZ, \gamma\gamma\}$ are the three channels with the most sensitivity to Higgs boson searches. So the uncertainties in Fig. 2(a) could be reduced significantly in the future, which would then allow for better discrimination between the Higgs boson and the triplet imposter.

In Fig. 2(b) we show the extracted $P_{g/V}$ from data using only diphoton final states, which have the best precision, as

³One could include higher-dimensional operators which break custodial invariance to shift $D_{W/Z}$ away from the SM value for a Higgs boson [33]. However, a potentially large effect is needed, implying a low cutoff for the higher-dimensional operators and new light degrees of freedom at the electroweak scale, which may be in tension with null results from direct searches.

well as the predicted ratio for the dilaton and the singlet imposters. The SM expectations at 7 and 8 TeV are [34]

$$P_{g/V}^{(\text{SM})}(7 \text{ TeV}) = 12.65, \quad P_{g/V}^{(\text{SM})}(8 \text{ TeV}) = 12.52. \quad (21)$$

However, while ATLAS presented best-fit signal strengths in gluon fusion and VBF production channels in Ref. [25], CMS only presented best-fit signal strengths in the ‘‘VBF-tag’’ category, which is expected to have some $gg \rightarrow h$ contamination with the additional two jets arising from higher-order QCD effects. To account for this contamination, we include a 25% contribution from gluon fusion for the CMS VBF-tag signal strength:

$$\sigma(pp \rightarrow S + 2j) = \epsilon \sigma(gg \rightarrow S + 2j) + (1 - \epsilon) \sigma(\text{VBF} \rightarrow S + 2j), \quad (22)$$

where CMS assumes $\epsilon \sim 0.25$ [35]. We can then relate the CMS measured value of $P_{g/V}^{\text{VBF-tag}}$ to the true $P_{g/V}$:

$$P_{g/V}^{\text{VBF-tag}} = \frac{P_{g/V}}{1 + \epsilon(P_{g/V} - 1)}, \quad (23)$$

which can be used to obtain the true $P_{g/V}$ from the CMS measurements. Such a conversion is not needed for the ATLAS results since the numbers are presented in terms of production channels, not selection categories. In the end we find

$$P_{g/V}^{\text{meas}} = 7.5_{-3.9}^{+4.0}, \quad (24)$$

which is the combined value for the ATLAS and CMS diphoton measurements. We see that the SM value in

Eq. (21) is consistent with the measured value. The observation that

$$P_{g/V} \sim P_{g/V}^{(\text{SM})} \quad (25)$$

suggests that the enhancements in $B\sigma_{gg}(\gamma\gamma)$ and $B\sigma_{\text{VBF}}(\gamma\gamma)$ could be explained simultaneously with an enhanced diphoton partial width resulting from an increased c_γ or $c_{5\gamma}$. We will see that this is indeed the case when fitting the event rates directly.

It turns out that $P_{g/V}$ could be used as a discriminator for the plain-vanilla dilaton imposter [36], which at 125 GeV gives

$$P_{g/V}^{(\text{D})} = 140 \times P_{g/V}^{(\text{SM})} \sim 1700, \quad (26)$$

which is clearly disfavored strongly by current measurements. This prediction of the ratio holds for the radion in the Randall-Sundrum model as well [37]. Essentially, a dilaton imposter is ruled out as soon as one can establish the presence of the VBF production channel. In Fig. 2(b) we show the ratio $P_{g/V}$ for the dilaton and singlet imposters, as well as the SM expectation and the value extracted from current data.

B. Model-dependent fits in all channels

Since many significant cross-section measurements have been made by the LHC and the Tevatron, we can fit the parameters of the model. Since the dilaton and the singlet imposters can not fit the model-independent ratios considered in the previous subsection, we only consider the Higgs boson and the triplet imposter when fitting all channels. In order to include data in the $b\bar{b}$ and $\tau\tau$ channels, we need to introduce the Higgs couplings to $b\bar{b}$ and $\tau\tau$:

$$\mathcal{L}_{hff} = c_b \frac{m_b}{v} h\bar{b}b + c_\tau \frac{m_\tau}{v} h\bar{\tau}\tau, \quad (27)$$

where $c_b^{(\text{SM})} = c_\tau^{(\text{SM})} = 1$. On the other hand, the triplet imposter does not have renormalizable couplings to SM fermions, so we simply set

$$\mathcal{L}_{h_s f f} = 0. \quad (28)$$

For the total width, we parametrize it as

$$\Gamma_{\text{tot}}^h = \sum_{V_1 V_2} \Gamma(h \rightarrow V_1 V_2) + \sum_f \Gamma(h \rightarrow f\bar{f}), \quad (29)$$

$$\Gamma_{\text{tot}}^{h_5} = \sum_{V_1 V_2} \Gamma(h_5 \rightarrow V_1 V_2). \quad (30)$$

Therefore the total width depends on all the c coefficients in the effective couplings during the fit. In principle one could introduce an extra free parameter in the total width to incorporate the possibility that the scalar could decay into other channels that have not been observed. In the end, we fit five parameters, $\{c_g, c_V, c_\gamma, c_b, c_\tau\}$, for the Higgs boson and three parameters, $\{c_{5g}, c_{5V}, c_{5\gamma}\}$, for the triplet imposter.

For the fitting procedure, we assume Gaussian uncertainties since a full treatment of the experimental uncertainties is beyond the scope of this work. We then fit the event rate measurements by minimizing the χ^2 :

$$\chi^2 = \sum_i \frac{\left(\tilde{\sigma}^i - \frac{\tilde{\Gamma}_{\text{prod}}^i \tilde{\Gamma}_{\text{decay}}^i}{\tilde{\Gamma}_{\text{total}}^i} \right)^2}{(\delta \tilde{\sigma}^i)^2}, \quad (31)$$

where $\tilde{\sigma}$ and $\tilde{\Gamma}$ are the signal cross section and decay width scaled with respect to the SM expectation, respectively. The measurement uncertainty on the cross section is given by $\delta \tilde{\sigma}$ and the asymmetric errors are retained.

The outcome of the fits is summarized in Table I, where we showed the χ^2 per degree-of-freedom for an SM Higgs boson with all the effective couplings fixed at the SM value, a generic Higgs boson with free varying effective couplings, and a triplet imposter. We see that a generic Higgs boson gives the best fit among the three to the current data with a p -value of 0.27, while the SM Higgs and a triplet imposter give increasingly worse fits (p -values of 0.70 and 0.84, respectively). We can also estimate the parameters of the Higgs boson and the triplet imposter at the 1σ level by determining the interval about which $\Delta\chi^2 \leq 1$, also shown in Table I. In the generic Higgs case, both c_g and c_V have best-fit values very close to the SM expectations, while c_γ is significantly enhanced over the SM expectation. The best-fit $c_\tau = 0$ is driven by the lack of excess in the CMS $\tau\tau$ measurement. In the triplet case, we generally find lower best-fit values of c_g , c_V and c_γ .

TABLE I. Comparison of fits for an SM Higgs, a generic Higgs boson, and a triplet imposter. One-dimensional parameter estimates in the custodial singlet (h_1) and triplet (h_5) models are under the total width assumptions. Uncertainties indicate the 1σ range. The SM Higgs boson is encapsulated in the custodial singlet scenario with $\Gamma_{\text{tot}} = \Gamma_{\text{tot}}^{h_{\text{SM}}}$, with $c_g = c_V = 1$ and $c_\gamma = 6.48$.

	χ^2/ν	p -value	c_g	c_V	c_γ	c_b	c_τ
SM Higgs	1.08	0.63	1	1	6.48	1	1
Higgs Boson	0.74	0.27	$0.92^{+0.30}_{-0.19}$	$1.07^{+0.15}_{-0.17}$	$9.7^{+1.9}_{-1.8}$	$1.1^{+0.5}_{-0.4}$	<0.73
Triplet Imposter	1.34	0.84	$0.37^{+0.08}_{-0.06}$	$0.45^{+0.10}_{-0.09}$	$3.8^{+0.5}_{-0.6}$

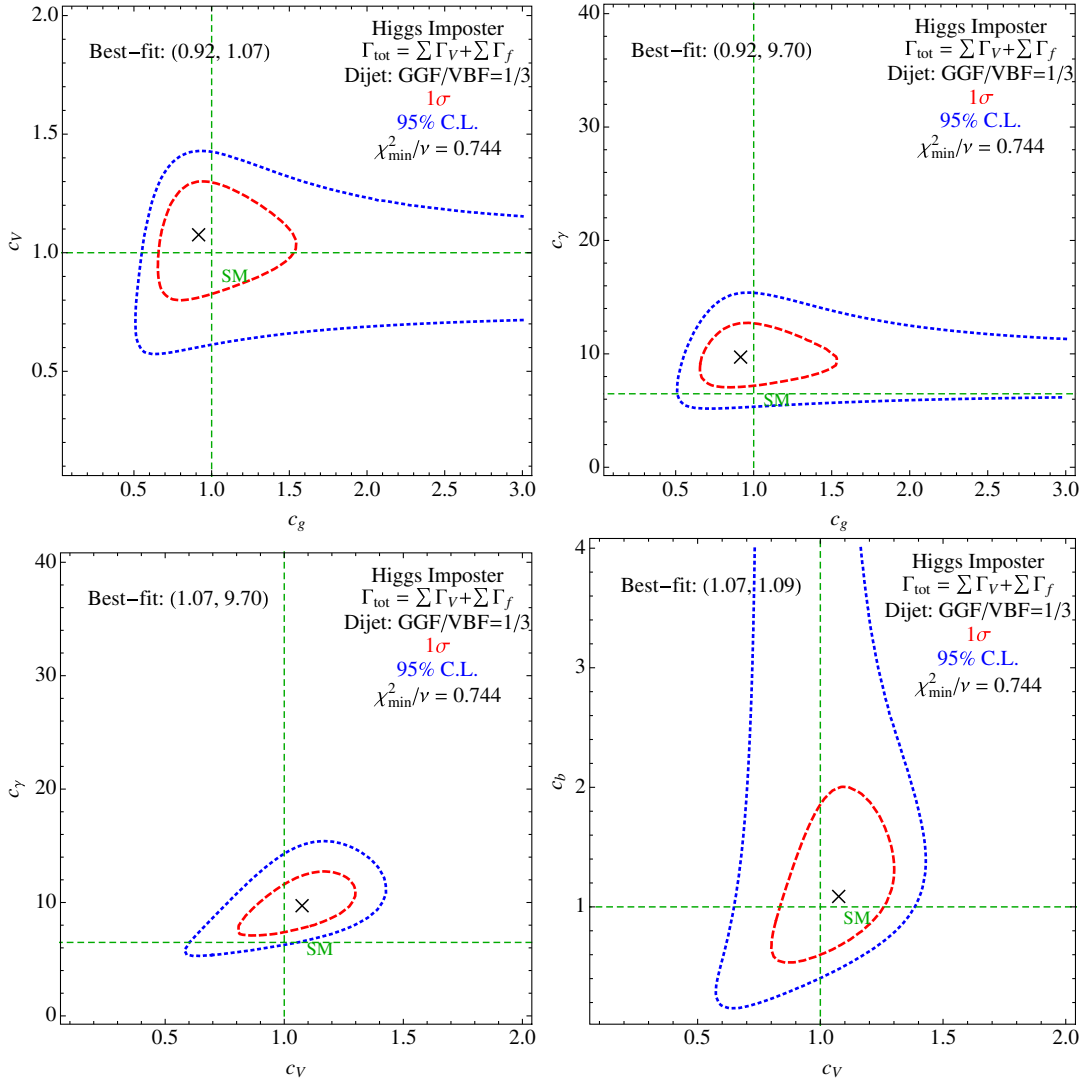


FIG. 3 (color online). Two-dimensional contours for four pairs of effective couplings. (c_V, c_g) enters into the decays into WW and ZZ from gluon fusion production. (c_γ, c_g) enters into the decays into diphotons from gluon fusion production. (c_γ, c_V) enters into the decays into diphotons from vector boson fusion production. (c_b, c_V) enters into the decays into $b\bar{b}$ from associated production with W/Z .

This is expected as we assume the triplet does not decay into fermions and the corresponding total width is therefore smaller than in the Higgs case, which gives rise to larger branching fractions and lower production cross sections.

In Fig. 3, we show the joint probability map in the plane of two model parameters.⁴ We choose to show the contours for the following four pairs of coefficients which enter into the channels with significant excesses: (c_V, c_g) for $gg \rightarrow h \rightarrow WW/ZZ$, (c_γ, c_g) for $gg \rightarrow h \rightarrow \gamma\gamma$, (c_γ, c_V) for $VBF \rightarrow h \rightarrow \gamma\gamma$, and (c_b, c_V) for $Vh \rightarrow V + b\bar{b}$. Since

⁴Note that since these are joint two-dimensional distributions, the 1 σ region may lie outside the one-parameter confidence intervals shown in Table I.

production primarily occurs through gg fusion and VBF, we expect to see the values of c_g and c_V having a strong upper bound, while the value of c_γ is allowed to rise well beyond the SM value of $c_\gamma^{\text{SM}} = 6.48$. However, since gg fusion can contaminate the dijet channel, the value of c_g can rise to compensate for a lower c_V . Indeed, at the 2 σ level, the value of c_g can be quite large. We generally find good agreement with the SM expectation, with the exception of c_γ . Indeed, the one-dimensional parameter fit of c_γ is nearly 2 σ away from the SM value. Overall, to fit the data, we require an enhancement to $\gamma\gamma$.

The corresponding two-dimensional contours for the triplet imposter are shown in Fig. 4. In this scenario the resonance decays into vector bosons with no appreciable decay into $f\bar{f}$. This is immediately at odds with the

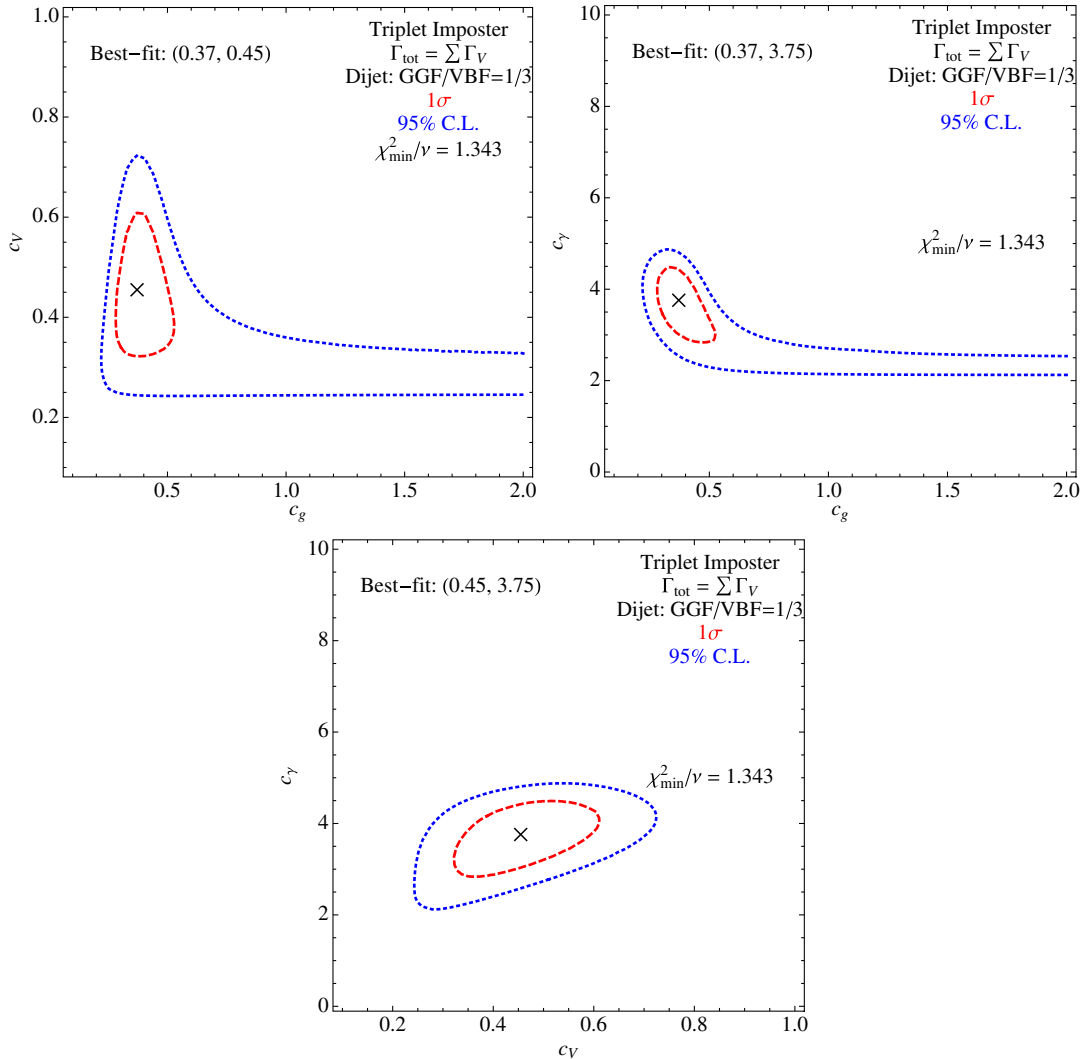


FIG. 4 (color online). Two-dimensional χ^2 contours for the triplet imposter. There are only three pairs of effective couplings, which enter into the event rates in the V_1V_2 channel. The triplet imposter does not decay into $b\bar{b}$ and $\tau\tau$ final states by assumptions.

Tevatron $Vh \rightarrow Vb\bar{b}$ result. However, the absence of a signal in the CMS measurements in the $VBF \rightarrow h \rightarrow \tau\tau$ channel supports this possibility. Due to the absence of the fermionic decay modes, we expect the total width to be smaller than in the singlet case, which is consistent with the fits shown in Fig. 4. The value of c_g is substantially lower than what is expected in the SM, meaning the total production of the scalar is suppressed. However, the decay branching fraction to $\gamma\gamma$ and WW/ZZ is increased due to the lower total width.

IV. CONCLUSIONS

Under the assumption that the new resonance discovered at the LHC is a CP -even scalar particle with a mass of 125 GeV, we have performed a general analysis of its possible electroweak quantum numbers. We have used a naive combination of the latest data from ATLAS, CMS,

and the Tevatron experiments, focusing on the four possible decays into pairs of electroweak gauge bosons, $\{WW, ZZ, \gamma\gamma, Z\gamma\}$, but also taking into account the two most important decay channels into pairs of fermions, $b\bar{b}$ and $\tau\tau$.

We have seen that interpreting the new particle as an electroweak singlet with loop-induced couplings to W and Z is strongly disfavored by current data. A plain-vanilla dilaton arising from scenarios where the SM gauge bosons are part of the conformal dynamics is disfavored as well. It will be important for the LHC experiments to quantify this statement, both by better constraints on decays to $Z\gamma$, and by more accurate measurements of the VBF production modes. In the latter regard we note the critical importance of having reliable estimates of the contamination of VBF analyses by gg fusion-initiated signal events.

Using chi-squared fits to the relevant free parameters, we have compared the compatibility of current data between

an SM Higgs boson, a more general custodial singlet boson, and a custodial 5-plet boson as would arise from an electroweak scalar triplet. All of the fits show some tension with the data, but the differences in the fit quality are not large. Thus, for example, one can not yet exclude the possibility that the new particle is the neutral member of electroweak triplets, provided that one is willing to discount the Tevatron excess in $b\bar{b}$. Similarly, one cannot greatly prefer an SM Higgs over a more general custodial singlet scalar, especially if one takes seriously the lack of a $\tau\tau$ excess in the CMS data. As we have seen, precise measurements of the ratios $D_{W/Z}$ and $D_{\gamma/Z}$ offer a clean way of distinguishing a triplet imposter from a Higgs

boson, but currently the uncertainties in these quantities are too large, and the central values actually favor the triplet imposter.

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- [1] S. Chatrchyan *et al.* (CMS Collaboration), *Phys. Lett. B* **716**, 30 (2012); G. Aad *et al.* (ATLAS Collaboration), *Phys. Lett. B* **716**, 1 (2012).
 - [2] S. Weinberg, *Phys. Rev. Lett.* **19**, 1264 (1967).
 - [3] L.D. Landau, *Dokl. Akad. Nauk SSSR* **60**, 207 (1948); C.N. Yang, *Phys. Rev.* **77**, 242 (1950).
 - [4] Q.H. Cao, C.B. Jackson, W.Y. Keung, I. Low, and J. Shu, *Phys. Rev. D* **81**, 015010 (2010).
 - [5] Y. Gao, A.V. Gritsan, Z. Guo, K. Melnikov, M. Schulze, and N.V. Tran, *Phys. Rev. D* **81**, 075022 (2010); A. De Rujula, J. Lykken, M. Pierini, C. Rogan, and M. Spiropulu, *Phys. Rev. D* **82**, 013003 (2010).
 - [6] P. Sikivie, L. Susskind, M.B. Voloshin, and V.I. Zakharov, *Nucl. Phys.* **B173**, 189 (1980).
 - [7] K. Nakamura *et al.* (Particle Data Group Collaboration), *J. Phys. G* **37**, 075021 (2010).
 - [8] I. Low and J. Lykken, *J. High Energy Phys.* **10** (2010) 053.
 - [9] E. Gildener and S. Weinberg, *Phys. Rev. D* **13**, 3333 (1976).
 - [10] C. Csaki, J. Hubisz, and S.J. Lee, *Phys. Rev. D* **76**, 125015 (2007); W.D. Goldberger, B. Grinstein, and W. Skiba, *Phys. Rev. Lett.* **100**, 111802 (2008); J. Fan, W.D. Goldberger, A. Ross, and W. Skiba, *Phys. Rev. D* **79**, 035017 (2009); Y. Eshel, S.J. Lee, G. Perez, and Y. Soreq, *J. High Energy Phys.* **10** (2011) 015; K. Cheung and T.-C. Yuan, *Phys. Rev. Lett.* **108**, 141602 (2012).
 - [11] W.D. Goldberger and M.B. Wise, *Phys. Rev. Lett.* **83**, 4922 (1999).
 - [12] J.M. Maldacena, *Adv. Theor. Math. Phys.* **2**, 231 (1998).
 - [13] I. Low, J. Lykken, and G. Shaughnessy, *Phys. Rev. D* **84**, 035027 (2011).
 - [14] H. Georgi and M. Machacek, *Nucl. Phys.* **B262**, 463 (1985); R.S. Chivukula and H. Georgi, *Phys. Lett. B* **182**, 181 (1986); M.S. Chanowitz and M. Golden, *Phys. Lett.* **165B**, 105 (1985); J.F. Gunion, R. Vega, and J. Wudka, *Phys. Rev. D* **42**, 1673 (1990); S. Chang, C.A. Newby, N. Raj, and C. Wanotayaroj, [arXiv:1207.0493](https://arxiv.org/abs/1207.0493) [*Phys. Rev. D* (to be published)].
 - [15] G.F. Giudice, R. Rattazzi, and J.D. Wells, *Nucl. Phys.* **B595**, 250 (2001).
 - [16] V. Barger, P. Langacker, M. McCaskey, M.J. Ramsey-Musolf, and G. Shaughnessy, *Phys. Rev. D* **77**, 035005 (2008).
 - [17] H.E. Logan and M.-A. Roy, *Phys. Rev. D* **82**, 115011 (2010).
 - [18] D. Carmi, A. Falkowski, E. Kuflik, and T. Volansky, *J. High Energy Phys.* **07** (2012) 136; A. Azatov, R. Contino, and J. Galloway, *J. High Energy Phys.* **04** (2012) 127; J.R. Espinosa, C. Grojean, M. Muhlleitner, and M. Trott, *J. High Energy Phys.* **05** (2012) 097; P.P. Giardino, K. Kannike, M. Raidal, and A. Strumia, *J. High Energy Phys.* **06** (2012) 117; J. Ellis and T. You, *J. High Energy Phys.* **06** (2012) 140; M. Klute, R. Lafaye, T. Plehn, M. Rauch, and D. Zerwas, *Phys. Rev. Lett.* **109**, 101801 (2012).
 - [19] M. Carena, I. Low, and C.E.M. Wagner, *J. High Energy Phys.* **08** (2012) 060.
 - [20] Tevatron New Physics Higgs Working Group and CDF and D0 Collaborations, [arXiv:1207.0449](https://arxiv.org/abs/1207.0449).
 - [21] ATLAS Collaboration, Report No. ATLAS-CONF-2012-098.
 - [22] G. Aad *et al.* (ATLAS Collaboration), *Phys. Rev. D* **86**, 032003 (2012).
 - [23] CMS Collaboration, CMS, Report No. PAG HIG-12-020.
 - [24] ATLAS Collaboration, Report No. ATLAS-CONF-2012-091.
 - [25] G. Aad *et al.* (The ATLAS Collaboration), *Phys. Lett. B* **716**, 1 (2012).
 - [26] CMS Collaboration, CMS, Report No. PAG HIG-12-018.
 - [27] D. Zeppenfeld, R. Kinnunen, A. Nikitenko, and E. Richter-Was, *Phys. Rev. D* **62**, 013009 (2000).
 - [28] G.F. Giudice, C. Grojean, A. Pomarol, and R. Rattazzi, *J. High Energy Phys.* **06** (2007) 045.
 - [29] J.S. Gainer, W.-Y. Keung, I. Low, and P. Schwaller, *Phys. Rev. D* **86**, 033010 (2012).
 - [30] A. Djouadi, *Phys. Rep.* **457**, 1 (2008).
 - [31] G. Aad *et al.* (ATLAS Collaboration), *J. High Energy Phys.* **09** (2011) 072.

- [32] S. Chatrchyan *et al.* (CMS Collaboration), *Phys. Lett. B* **701**, 535 (2011).
- [33] M. Farina, C. Grojean, and E. Salvioni, *J. High Energy Phys.* **07** (2012) 012.
- [34] <https://twiki.cern.ch/twiki/bin/view/LHCPhysics/CrossSections>.
- [35] CMS Collaboration, CMS, Report No. PAG HIG-12-015.
- [36] B. Coleppa, T. Gregoire, and H. E. Logan, *Phys. Rev. D* **85**, 055001 (2012); H. Logan, in *The Next Stretch of the Higgs Magnificent Mile, Chicago, 2012* (to be published).
- [37] V. Barger, M. Ishida, and W.-Y. Keung, *Phys. Rev. Lett.* **108**, 101802 (2012).