LHC signatures of a generalized Georgi-Machacek model

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(Received 6 March 2025; accepted 1 August 2025; published 28 August 2025)

Vector-boson fusion productions of new Higgs bosons decaying into pairs of electroweak gauge bosons ($W^\pm W^\pm$, WZ, and ZZ) are smoking-gun signature of new Higgs bosons in general, and of the Georgi-Machacek (GM) model in particular. Notably, ATLAS has observed a 3.3σ excess in $W^\pm W^\pm$ at $\approx$450 GeV and a 2.8σ excess in the WZ channel at $\approx$375 GeV, while compact muon solenoid reported weaker-than-expected limits at these masses. In this article, we consider a generalized version of the GM model without the custodial $SU(2)_C$ symmetry in the scalar potential as a minimal model to explain these excesses. In the limit of small mixing among the Higgs bosons, the $W^\pm W^\pm$ and WZ signals can be explained by the doubly and singly charged Higgs bosons originating primarily from the $Y = 1$ triplet, while respecting the bounds from ZZ searches. Furthermore, the neutral Higgs boson mostly contained in the $Y = 0$ triplet can account for the ($\approx 4\sigma$) excess at $\approx$152 GeV in associated diphoton production. Furthermore, the model is consistent with the constraints from vacuum stability and the Standard Model Higgs signal strength measurements, such that it can provide a valid combined explanation for some of the most significant LHC excesses.

DOI: 10.1103/wzjg-3fqg

I. INTRODUCTION

The Standard Model (SM) of particle physics is the currently accepted theoretical description of the known fundamental constituents and interactions of matter. It has been successfully tested experimentally [1], culminating in the discovery of the Higgs boson at the Large Hadron Collider (LHC) [2,3]. Although the observed 125 GeV scalar has properties consistent with those predicted by the SM [4,5], this does not rule out the existence of additional scalar bosons. In fact, as long as their contribution to electroweak symmetry (EW) breaking is subleading and they approximately respect custodial symmetry [$\rho = m_W^2/(m_Z^2 c_W^2) \approx 1$, where c_W is the cosine of the Weinberg angle], they are in general compatible with the bounds from SM Higgs signal strength measurements [4,5] and electroweak (EW) precision data [6,7].

Interestingly, ATLAS observed excesses in vector boson fusion (VFB) topologies in the $W^\pm W^\pm$ [8] and WZ channels [9] (see Fig. 1), with local significances of 3.3σ at 450 GeV and 2.8σ at 375 GeV, respectively. Furthermore, compact muon solenoid (CMS) reported weaker-than-expected limits in these mass regions [10]. Combining these data, Ref. [11] finds a combined global significance of 4σ for the $W^\pm W^\pm$ and WZ excesses. As the $W^\pm W^\pm$ excess requires a doubly charged Higgs boson, this requires as a minimal-dimensional representation a triplet scalar with hypercharge $Y = -1$.[1] Since their VEVs break custodial symmetry, they must be small—of the order of a few GeV—to result in a W mass prediction compatible with experiments. Because the couplings of the scalars to a pair

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[1] The requirement of a vacuum expectation value (VEV) to generate VBF together with the masslessness of the photon demands that the new field has an electrically neutral component.

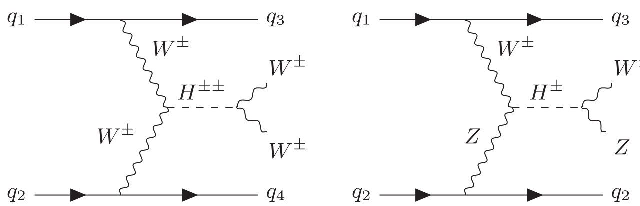

FIG. 1. Feynman diagrams showing the VBF production of singly and doubly charged Higgses decaying to EW gauge bosons.

of EW gauge bosons are proportional to their VEVs, their vector boson fusion (VBF) cross sections are also small. Consequently, despite the distinct experimental signatures, VBF processes in these models are beyond the LHC reach.

Georgi and Machacek (GM) proposed supplementing the SM with a $Y=0$ and a $Y=1$ scalar triplet. By imposing custodial symmetry, the triplets are arranged into an $SU(2)_C$ ($C \equiv$ custodial) multiplet [12,13] so that $\rho = 1$ at the tree level. This allows sizable triplet VEVs, leading to potentially measurable VBF production cross sections for new Higgses at the LHC. In particular, VBF production of $W^\pm W^\pm$ and WZ are smoking-gun signatures of the GM model [14–39], as they are absent in two-Higgs-doublet models (or loop-suppressed for WZ in the case of CP violation [40]).

However, these indications for new Higgs bosons cannot be (fully) explained within the canonical custodial-symmetric GM model [41], as it predicts the gaugephilic Higgs bosons to be mass degenerate. To overcome this limitation, we consider the generic GM model (gGMM) [36,42,43], where the requirement of custodial symmetry in the Higgs potential is relaxed. In this framework, the mass degeneracy among the Higgses is lifted, while the VEVs of the two triplet Higgses have to remain approximately equal in size, as required by EW precision data. Furthermore, we examine whether the indications for a new Higgs in associated diphoton production at 152 GeV [44–49] obtained from the sidebands of SM Higgs analyses [50,51][2] can be explained within the gGMM.[3]

II. THE MODEL

Like the canonical GM model, the gGMM is constructed by adding an $SU(2)_L$ triplet scalar χ with $Y=1$ and a triplet scalar ζ with $Y=0$ to the SM $SU(2)_L$ Higgs doublet ϕ. These Higgses decompose into their $SU(2)_L$ components as

$$\phi = \begin{pmatrix} \phi^+ \\ \phi^0 \end{pmatrix}, \quad \chi = \begin{pmatrix} \frac{\chi^+}{\sqrt{2}} & -\chi^{++} \\ \chi^0 & -\frac{\chi^+}{\sqrt{2}} \end{pmatrix}, \quad \zeta = \begin{pmatrix} \frac{\zeta^0}{\sqrt{2}} & -\zeta^+ \\ -\zeta^- & -\frac{\zeta^0}{\sqrt{2}} \end{pmatrix},$$

with

$$\phi^0 = \frac{\phi_r + i\phi_i + v_\phi}{\sqrt{2}}, \quad \chi^0 = \frac{\chi_r + i\chi_i}{\sqrt{2}} + v_\chi, \quad \zeta^0 = \zeta_r + v_\zeta,$$

where v_ϕ, v_χ, and v_ζ are the respective VEVs. With these conventions, the ρ parameter at tree level is given by

$$\rho = \frac{v^2}{v^2 + 4(v_\chi^2 - v_\zeta^2)}, \tag{1}$$

where $v^2 = v_\phi^2 + 4(v_\chi^2 + v_\zeta^2) \approx (246\ \text{GeV})^2$. From $\rho \approx 1$, it follows that $v_\chi \approx v_\zeta$. Furthermore, we know experimentally that the 125 GeV Higgs is SM-like, i.e., $h \approx \phi_r$, which implies $v \approx v_\phi$ or $v_\chi^2, v_\zeta^2 \ll v^2$.

Without the custodial symmetry of the canonical GM model, the scalar potential is given by [43]

$$\begin{aligned} V = &-m_\phi^2(\phi^\dagger\phi) - m_\chi^2 \mathrm{Tr}(\chi^\dagger\chi) - m_\zeta^2 \mathrm{Tr}(\zeta^2) + (\mu_{\phi\chi}\phi^\dagger\chi\tilde{\phi} + \text{H.c.}) + \mu_{\phi\zeta}\phi^\dagger\zeta\phi + \mu_{\chi\zeta}\mathrm{Tr}(\chi^\dagger\chi\zeta) \\ &+ \lambda(\phi^\dagger\phi)^2 + \rho_1[\mathrm{Tr}(\chi^\dagger\chi)]^2 + \rho_2\mathrm{Tr}(\chi^\dagger\chi\chi^\dagger\chi) + \rho_3[\mathrm{Tr}(\zeta^2)]^2 + \rho_4\mathrm{Tr}(\chi^\dagger\chi)\mathrm{Tr}(\zeta^2) + \rho_5\mathrm{Tr}(\chi^\dagger\zeta)\mathrm{Tr}(\zeta\chi) \\ &+ \sigma_1\mathrm{Tr}(\chi^\dagger\chi)\phi^\dagger\phi + \sigma_2\phi^\dagger\chi\chi^\dagger\phi + \sigma_3\mathrm{Tr}(\zeta^2)\phi^\dagger\phi + (\sigma_4\phi^\dagger\chi\zeta\tilde{\phi} + \text{H.c.}), \end{aligned} \tag{2}$$

where only $\mu_{\phi\chi}$ and σ_4 can be complex (their phases are related [43]), but are assumed to be real, albeit for simplicity, while the other parameters are real by construction.

As outlined in the introduction, we aim at a proof of principle that the $W^\pm W^\pm$ and WZ excesses can be explained in the gGMM while remaining consistent with existing bounds and accommodating the indication for a 152 GeV diphoton resonance. For this purpose, we consider a simplified setup in which the $Y=1$ and $Y=0$ triplets, χ and ζ, are decoupled from each other in a first approximation. This is achieved by setting $\mu_{\chi\zeta} \approx 0, \rho_4 \approx 0, \rho_5 \approx 0, \sigma_4 \approx 0$. In the limit of small mixing among the triplets, the model can be treated as a combination of the $Y=0$ (see, e.g., [49,59–65]) and the $Y=1$ (see, e.g., [66–71]) models, with the key difference that the

[2]The existence of a new scalar boson with a mass of around 150 GeV is a prediction [52–55] of the multilepton anomalies [56,57].

[3]An explanation of these excesses within the GM model, albeit without detailed calculation, was proposed in Ref. [58].

triplet VEVs can be sizable, as their contributions to the ρ parameter cancel (partially).

The particle spectrum contains a doubly charged Higgs $H^{\pm\pm} \equiv \chi^{\pm\pm}$, which originates purely from the $Y = 1$ triplet. For small mixing, the remaining mass eigenstates can be identified with the interaction ones as

$$h \approx \phi_r, \quad H^0_{\chi(\zeta)} \approx \chi_r(\zeta_r), \quad A^0 \approx \chi_i, \quad H^\pm_{\chi(\zeta)} \approx \chi(\zeta)^\pm, \quad (3)$$

while $\phi^\pm$ (ϕ_i) is approximately the charged (neutral) would-be Goldstone boson. Concerning their masses, the splitting between $m_{H^0_\zeta}$ and $m_{H^\pm_\zeta}$ is of the order of v_ζ^2/v (i.e., $m_{H^0_\zeta} \approx m_{H^\pm_\zeta}$), while the splitting among A^0/H^0_χ, $H^\pm_\chi$, and $H^{\pm\pm}$ is of order v. To be more specific, with modulus $v_{\chi,\zeta}$ corrections, we have

$$m^2_{H^{\pm\pm}} - m^2_{H^\pm_\chi} \approx m^2_{H^\pm_\chi} - m^2_{A^0,H^0_\chi}. \quad (4)$$

Via their VEVs, the new Higgses not only acquire diboson couplings but also induce fermionic couplings for the singly charged and neutral scalars through mixing with the SM Goldstone bosons. The doubly charged Higgs couples to $W^\pm W^\pm$ and $H^\pm_\chi W^\pm$. In the limit of small mixing,

$$H^{\mp\mp} W^\pm_\mu W^\pm_\nu:\ i2g^2 v_\chi g_{\mu\nu},$$
$$H^{\mp\mp} H^\pm_\chi W^\pm_\mu:\ \mp ig(-p_\mu^{H^\pm_\chi} + p_\mu^{H^{\mp\mp}}). \quad (5)$$

Similarly, the singly charged Higgs $H^\pm_\zeta$ ($H^\pm_\chi$) couples to WZ and $H^0_\zeta W$ ($H^0_\chi W$, $A^0 W$, and $H^{\pm\pm} W^\mp$), and their coupling with hW becomes relevant for sizable mixing between the neutral Higgses $H^0_{\chi,\zeta}$ and h, denoted as $\alpha_{\phi\chi}$ and $\alpha_{\phi\zeta}$. They also couple to fermions with a strength of $\approx 2v_{\chi,\zeta}/vY^f$. Likewise, $H^0_{\chi,\zeta}$ couples to fermions, with a strength proportional to $\alpha_{\phi\chi,\phi\zeta}$. Without mixing, H^0_χ couples to both WW and ZZ while H^0_ζ couples only to WW. These couplings, normalized to their SM counterparts, are $2\sqrt{2}v_\chi/v$, $4\sqrt{2}v_\chi/v$, and $4v_\zeta/v$, respectively. Also $H^0_{\chi,\zeta}$ has a coupling to hh proportional to $\mu_{\chi\phi} + (\sigma_1 + \sigma_2)v_\chi$ [$\mu_{\phi\zeta} - 2\sqrt{2}\sigma_3 v_\zeta$] which, via the minimization conditions, is of the order of $v_{\chi,\zeta}$. Finally, the CP-odd Higgs A^0 couples to fermions with a strength of $\approx 2\sqrt{2}v_\chi/vY^f$, and to hZ with a strength of $\approx g/c_W\sqrt{2}v_\chi/v$. Note that these couplings receive corrections due to mixing with the SM Higgs.

III. PHENOMENOLOGY

Because we assume that the $Y = 0$ and $Y = 1$ triplets are decoupled from each other and have only a small mixing with the SM Higgs, to a good approximation, the phenomenology can be studied by considering the effects of the $Y = 1$ and $Y = 0$ triplet models separately. However, v_χ and v_ζ are allowed to be sizable, as their effect on the W mass can cancel.[4] Importantly, VBF production of the new Higgs bosons is driven by their diboson couplings, which are proportional to the VEVs. The production cross sections, in the limit of vanishing mixing, are shown in Fig. 2 as a function of their mass for $v_\chi = v_\zeta = 1$ GeV and scale as v_χ^2 and v_ζ^2 for the $Y = 1$ and $Y = 0$ triplets, respectively.

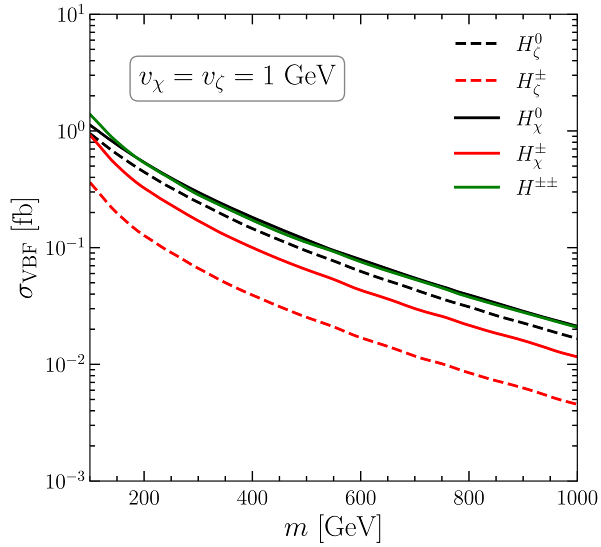

FIG. 2. VBF production cross sections for the new Higgs states for $v_\chi = v_\zeta = 1$ GeV at the 13 TeV LHC, including the next-to-next-to-leading order quantum chromodynamics corrections of Ref. [73].

The $W^\pm W^\pm$ excess can only be explained by the doubly charged Higgs of the $Y = 1$ triplet. Accordingly, we can fix $m_{H^{\pm\pm}} \approx 450$ GeV, for which $\sigma_{\mathrm{VBF}}(H^{\pm\pm}) \times \mathrm{Br}(H^{\pm\pm} \to W^\pm W^\pm) = (72 \pm 25)$ fb is preferred [8]. Note that this cross section is compatible with the weaker-than-expected 95% CL limit of 80 fb from CMS [10]. The best-fit point corresponds to $v_\chi \approx 23$ GeV (see the cyan band in Fig. 3) as long as the decay $H^{\pm\pm} \to H^\pm W^\pm$ is negligible, i.e., for $\mathrm{Br}(H^{\pm\pm} \to W^\pm W^\pm) \approx 100\%$, which holds for $\Delta m = m_{H^{\pm\pm}} - m_{H^\pm} \lessapprox 80$ GeV.

In fact, the WZ excess suggests $m_{H^\pm_\chi} \approx 375$ GeV, thereby ensuring $\mathrm{Br}(H^{\pm\pm} \to W^\pm W^\pm) \approx 100\%$. The best-fit point is $\sigma_{\mathrm{VBF}}(H^\pm_\chi) \times \mathrm{Br}(H^\pm_\chi \to W^\pm Z) \approx 60 \pm 20$ fb.[5] This is smaller than the 95% CL upper limit obtained by CMS (which is also a weaker-than-expected limit) [10]. The dominant decays of $H^\pm_\chi$ are WZ, $t\bar{b}$, and Wh. Since the

[4]The mass splitting among the χ-like Higgses significantly impacts the W mass prediction [72]. In contrast, $Y = 0$ triplet contribution is suppressed due to $\mathcal{O}$ (GeV) mass splitting, required by perturbative unitarity and vacuum stability.

[5]While ATLAS provides the preferred cross section for the $W^\pm W^\pm$ excess, they do not provide one for WZ. Therefore, we estimated the latter from the provided observed and expected limits at 375 GeV.

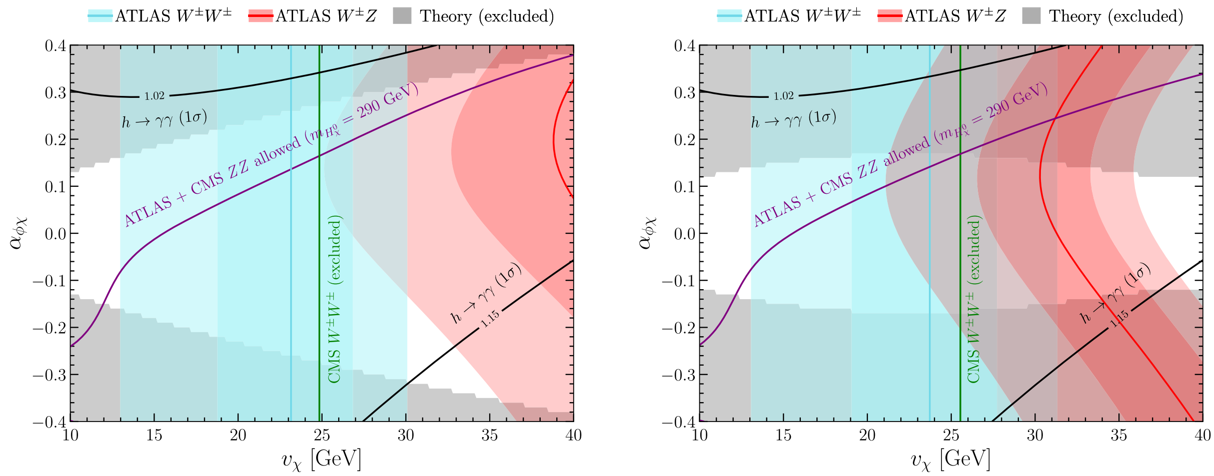

FIG. 3. Preferred (1σ and 2σ) and excluded (95% CL) regions in the v_χ-$\alpha_{\phi\chi}$ plane. The ATLAS measurement of $W^\pm W^\pm$ prefers the cyan region, while the region above the purple line is allowed by ZZ searches in VBF. Note that both limits are, to a good approximation, independent of the parameters of the $Y = 0$ triplet. However, the region (red) preferred by the ATLAS WZ search can be affected by the $Y = 0$ triplet parameters in the case of sizable mixing with the $Y = 1$ triplet. Also, the region (gray) excluded by theory constraints (vacuum stability and perturbativity) depends on the parameters of this sector. This can be seen by comparing the left plot to the right one. In the left plot, we fix $m_{H^0_\zeta} = 152$ GeV to explain the excess in associated diphoton production, while in the right plot, the charged Higgs states are close in mass (parametrized by $\kappa = 0.95$) and have sizable mixing among them, which maximize the effect in WZ (see text for details). The $h \to \gamma\gamma$ contour labels refer to the corresponding signal strength normalized to the SM.

latter depends on the mixing angle of H^0_χ with the SM Higgs, the red region explaining the WZ excess in Fig. 3 exhibits its maximal elongation toward small values of v_χ at $\alpha_{\phi\chi} \approx 0.15$ as this minimizes the $H^\pm W h$ coupling.

Given the masses of $H^{\pm\pm}$ and $H^\pm_\chi$, the mass of H^0_χ is approximately determined, i.e., $\approx 280 + O(v_{\chi,\zeta})$ GeV. In addition, v_χ fixes its production cross section, and its dominant decays (WW, ZZ and hh) depend, to a good approximation, only on $\alpha_{\phi\chi}$. The experimental bounds on $\sigma_{\rm VBF}(H^0_\chi \to ZZ)$ in this region of parameter space vary between 10 fb and 40 fb, depending on its exact mass [74,75]. For a mass of 290 GeV, the ATLAS and CMS limits coincide at $\approx$49 fb, and the resulting limit is shown in purple in Fig. 3 (the region above the purple lines is allowed). Note that this bound suggests a preference for small but positive values of $\alpha_{\phi\chi}$. Further, the constraint from $H^0_\chi \to hh \to 4b$ in VBF [76], which is approximately 1 pb, is easily satisfied. Finally, the CP-odd Higgs A^0—with a mass similar to that of H^0_χ—does not couple to WW, ZZ and is dominantly produced via gluon fusion with a cross section of $\approx 1.5\alpha^2_{\phi\chi}$ times that of a hypothetical SM Higgs with the same mass; it decays mainly to Zh and $b\bar{b}$. Therefore, the experimental limits are not constraining; even if the branching ratio to Zh were 100%, the bounds from the corresponding search [77] would not be violated, given the small production cross section.

We now discuss the impact of the $Y = 0$ triplet. If its quasi-mass-degenerate states H^0_ζ and $H^\pm_\zeta$ are much heavier or much lighter than $m_{H^\pm_\chi}$, the mixing is small, and the interplay with the bounds discussed so far is limited. In fact, the primary effect is on vacuum stability, which is sensitive to the full scalar potential. The resulting excluded regions are shown in gray in Fig. 3.

Only $Y = 0$ triplet masses in the range 180–200 GeV are excluded by multiletpton production at the LHC [78]. Thus, the $Y = 0$ triplet can explain the 152 GeV excess in associated diphoton production [46,47,49]. Given that the Drell-Yan production cross section $pp \to W^\pm \to H^0_\zeta H^\pm_\zeta$ and the dominant branching ratios of $H^\pm_\zeta$ are fixed, the excess prefers ${\rm Br}(H^0_\zeta \to \gamma\gamma) = (0.7 \pm 0.18)\%$ [49]. In Fig. 4, we show the region in parameter space that gives rise to this branching ratio (and thus explain the excess) for $\alpha_{\phi\chi} = 0.15$. Depending on v_ζ, H^0_ζ is also produced in VBF processes, and the resulting diphoton signatures have been studied for the SM Higgs [79,80]. From the sidebands, one can estimate that the cross section of a new Higgs with a mass of 152 GeV should not exceed 20% of that of the SM Higgs, excluding the region to the bottom right of the orange line in Fig. 4 (left).

We next project this scenario onto the v_χ-v_ζ plane in the right plot in Fig. 4. Here the green band illustrates the degree of cancellation in the W mass between the tree-level contribution from v_χ, v_ζ and the loop-induced contribution from the mass splitting among the $Y = 1$ components needed to respect the EW precision bounds. Contrary to the pure $Y = 0$ triplet model, where there are tensions between the preferred branching ratio from the di-photon excess, vacuum stability, and the SM Higgs signal strength measurements, the gGMM can resolve these issues.

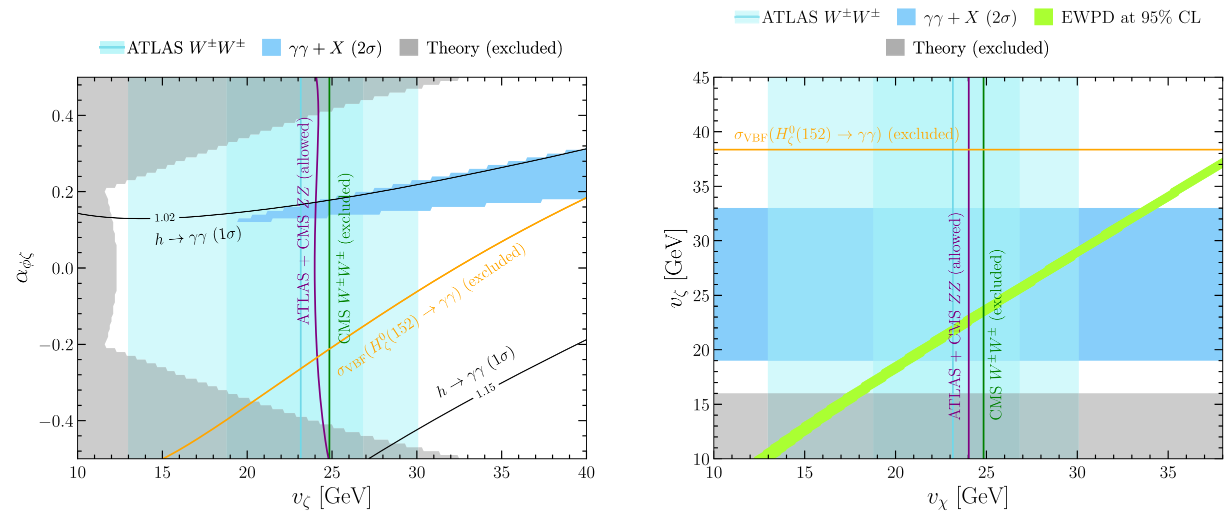

FIG. 4. Preferred (1σ and 2σ) and excluded (95% CL) regions for $m_{H^0_\zeta} = 152$ GeV in the v_ζ-$\alpha_{\phi\zeta}$ and v_χ-v_ζ planes for $\alpha_{\phi\chi} = 0.15$ (left) and $\alpha_{\phi\zeta} = 0.15$ (right), respectively. In the left plot, we assume $v_\chi \approx v_\zeta$, as justified by the green band from EW precision observables in the right plot. The $h \to \gamma\gamma$ contour labels refer to the corresponding signal strength normalized to the SM.

Finally, if one allows for similar masses of the two charged Higgses, one has sizable mixing among them due to v_χ and v_ζ. This mixing enhances $\mathrm{Br}(H^\pm_\chi \to WZ)$, thereby enlarging the overlap of the corresponding preferred region with that of $W^\pm W^\pm$. For this, we consider $m_{H^{\pm\pm}} = 450$ GeV, $m_{H^\pm_\chi} = 375$ GeV, $m_{H^\pm_\zeta} = \kappa \times \frac{v_\phi}{\sqrt{v_\phi^2 + 8v_\chi^2}} m_{H^\pm_\chi}$ GeV; $\kappa \to 1$ leads to maximal mixing. The resulting preferred and excluded regions for $\kappa = 0.95$ are shown in the right plot of Fig. 3.

IV. CONCLUSIONS AND OUTLOOK

In this article, we explored the gGMM as a minimal explanation of the excesses in VBF production in $W^\pm W^\pm$ and WZ at $\approx$450 GeV and $\approx$375 GeV with local significance of 3.3σ and 2.8σ, respectively. Working in the limit of small mixing angles, we fixed the masses of the $Y = 1$ triplet states to these values, which predicts its neutral components to be around 280 GeV. We showed that the resulting constraints from ZZ in VBF, SM Higgs signal strength, as well as the theoretical bounds from vacuum stability and perturbative unitarity, can be respected for VEVs capable of providing sufficiently high production cross sections (see Fig. 3). However, the regions preferred by the $W^\pm W^\pm$ and WZ channels only overlap at the 2σ level.

Furthermore, we showed that the neutral component of the $Y = 0$ triplet can account for the diphoton excess in associated production channels without violating the previously mentioned bounds or the $\gamma\gamma$ signal in VBF or EW precision data (see Fig. 4). In fact, the agreement with data is better than in the ΔSM (SM extended only with the $Y = 0$ triplet), as higher values of $\mathrm{Br}(H^0_\zeta \to \gamma\gamma)$ can be obtained without conflicting with vacuum stability or the SM Higgs signal strength measurements. Therefore, the gGGM can provide a combined explanation for some of the most significant LHC excesses.

ACKNOWLEDGMENTS

This research was supported by the Deutsche Forschungsgemeinschaft (DFG, German Research Foundation) under Grant No. 396021762—TRR 257. This research was supported in part by Grant No. NSF PHY-2309135 to the Kavli Institute for Theoretical Physics (KITP). The work of A. C. is supported by a professorship grant from the Swiss National Science Foundation (Grant. No. PP00P21_76884). S. P. M. would like to acknowledge the support of the Research Office of the University of the Witwatersrand.

DATA AVAILABILITY

The data are not publicly available. The data are available from the authors upon reasonable request.

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