Higgs couplings in a model with triplets

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Abstract

We study the couplings of a CP-even neutral Higgs boson $h$ in a model containing one scalar \(SU(2)_L\) doublet, one real triplet, and one complex triplet with hypercharge 1. Because the two triplets contribute to the $\rho$ parameter with opposite signs, the triplet vacuum expectation values can be sizable. We show that (i) the $hWW$ and $hZZ$ couplings can be larger than the corresponding values in the Standard Model, and (ii) the ratio of the $WW$ and $ZZ$ couplings of $h$ can be different than the corresponding ratio in the Standard Model. Neither of these results can occur in models containing only Higgs doublets. We also compute the rates for $gg \to h \to WW$ and $gg \to h \to ZZ$ and find that, for reasonable parameter values and $M_h \sim 140-180$ GeV, the hadron collider rate for $gg \to h \to WW$ ($ZZ$) can be up to 20% (5 times) larger than in the Standard Model. We discuss implications for Higgs coupling extraction at the LHC.

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I. INTRODUCTION

The primary purpose of the CERN Large Hadron Collider (LHC) is to shed light on the dynamics of electroweak symmetry breaking. In the Standard Model (SM), this involves the discovery of the Higgs boson and the measurement of its mass-generating couplings to fermions and gauge bosons. Most extensions of the SM contain one or more Higgs-like scalars that play a similar role in mass generation; the characterization of such models also involves measurement of the couplings of these scalars to the SM fermions and gauge bosons.

The LHC will be able to measure Higgs signal rates in a variety of production and decay modes \[1, 2\], especially if the Higgs mass lies in the range \(114\ \text{GeV} < M_h \lesssim 200\ \text{GeV}\), as suggested by direct searches \[3\] and electroweak precision data \[4\]. These signal rate measurements will provide access to the combinations of Higgs couplings that determine the production cross sections and decay branching ratios involved in each channel. By taking ratios of Higgs signal rates with either the production mode or the decay mode in common, ratios of Higgs couplings-squared can be found with no model-dependent assumptions \[5, 6\].

Unfolding the individual Higgs couplings from these rate measurements, however, is a major challenge. The main difficulty lies in the fact that the LHC will not provide an absolute Higgs production cross section measurement in any channel. Together with the possible presence of undetectable Higgs decay modes (e.g., \(h \to gg\) or decays into light quarks), this leads to a degeneracy in the Higgs coupling fit corresponding to increasing all production couplings by a common factor while simultaneously decreasing the detectable branching fractions by the same factor. This forces the adoption of theoretical assumptions in Higgs coupling fits. Analyses of LHC Higgs coupling extraction prospects have dealt with this difficulty either by assuming that the total Higgs width is dominated by the known SM decay modes \[5, 6\], or by assuming that the couplings of the Higgs to \(WW\) and \(ZZ\) are bounded from above by their SM values \[7–9\] (the latter analyses make the simultaneous assumption that the ratio of the \(hWW\) and \(hZZ\) couplings is the same as in the SM). The latter approach allows the total width of the Higgs to be bounded from above by a rate measurement in a channel with Higgs production and decay via the \(hWW\) or \(hZZ\) coupling (e.g., weak boson fusion with \(h \to WW\)). These two assumptions about the \(hWW\) and \(hZZ\) couplings are valid in any Higgs sector that contains only SU(2)\(_L\) doublets and/or singlets. Larger SU(2)\(_L\) multiplets are usually considered unimportant because their contribution to
the $hWW$ and $hZZ$ couplings are generally tightly constrained by the experimental limits on the $\rho$ parameter.

In this paper we explore the possibility that these assumptions about the $hWW$ and $hZZ$ couplings are violated. To do so, we need a model with at least one Higgs doublet (required to generate fermion masses), together with at least one larger SU(2)$_L$ representation (to give rise to group-theoretic coefficients in the $hWW$ and $hZZ$ couplings different from those of a doublet). Because the couplings of interest arise from quartic terms in the Lagrangian in which one scalar field is replaced by its vacuum expectation value (vev), we need the vev(s) of the larger SU(2)$_L$ multiplet(s) to be non-negligible.

To this end we study a model containing one Higgs doublet and two triplets first introduced by Georgi and Machacek [10]. One triplet is real with hypercharge zero while the other is complex with hypercharge 1. Because these two triplets contribute to the $\rho$ parameter with opposite sign, their vevs need not be small so long as they are chosen such that their contributions to $\rho$ cancel; this is achieved naturally by imposing a global “custodial” SU(2) symmetry [10–14]. (The required global symmetry is however broken by hypercharge, so it is no longer exact once radiative corrections are included [13].) The model contains three CP-even neutral scalars. We study one of these states, which we call $h$ and parameterize as a general mixture of the three weak eigenstates. We find that the $hWW$ and $hZZ$ couplings can be significantly larger than their SM values, and that the ratio of these couplings can be significantly different from the ratio in the SM.

This paper is organized as follows. In the next section we summarize the model and present general expressions for the couplings of interest. In Sec. III we study the range of $hWW$ and $hZZ$ couplings allowed in the model, subject to experimental constraints that set a lower bound on the doublet vev. We also consider the impact on Higgs production and decay in two of the most promising early Higgs discovery channels at the LHC, $gg \to h \to WW$ and $gg \to h \to ZZ$. In particular, we find that the $gg \to h \to ZZ$ rate can be significantly enhanced compared to that in the SM, due to a suppression of the $h \to WW$ branching fraction in some parts of parameter space without a simultaneous suppression of $h \to ZZ$. In Sec. IV we discuss ways to experimentally determine whether the $hZZ$ and $hWW$ couplings have the same ratio as in the SM and whether these couplings are larger than their SM values. In Sec. V we summarize our conclusions. Some computational details

1 We use the convention that the hypercharge of the SM Higgs doublet is 1/2.
II. THE MODEL

The Georgi-Machacek model \[10\] contains a complex SU(2)\textsubscript{L} scalar doublet \(\Phi\) with hypercharge \(Y = Q - T^3 = 1/2\), a real triplet \(\Xi\) with \(Y = 0\), and a complex triplet \(X\) with \(Y = 1\), with components and vevs denoted by

\[
\Phi = \begin{pmatrix}
\phi^+
\\
\frac{1}{\sqrt{2}}(v_\phi + \phi^0_r + i\phi^0_i)
\end{pmatrix}, \quad
\Xi = \begin{pmatrix}
\xi^+
\\
v_\xi + \xi^0
\\
\xi^-
\end{pmatrix}, \quad
X = \begin{pmatrix}
\chi^{++}
\\
\frac{1}{\sqrt{2}}(v_\chi + \chi^0_r + i\chi^0_i)
\end{pmatrix}.
\]

Starting from the canonically normalized covariant derivative terms in the scalar Lagrangian, we obtain the \(W\) and \(Z\) boson masses and their couplings to the CP-even neutral scalars as follows:

\[
\mathcal{L} \supset |D_\mu \Phi|^2 + \frac{1}{2} |D_\mu \Xi|^2 + |D_\mu X|^2
\supset (v_\phi + \phi^0_r)^2 \left[ \frac{g^2}{4} W^+_{\mu} W^-_{\mu} + \frac{g^2 + g'^2}{8} Z_{\mu} Z^{\mu} \right]
+ (v_\xi + \xi^0)^2 \left[ \frac{4 g^2}{4} W^+_{\mu} W^-_{\mu} \right]
+ (v_\chi + \chi^0_r)^2 \left[ \frac{2 g^2}{4} W^+_{\mu} W^-_{\mu} + 4 \frac{g^2 + g'^2}{8} Z_{\mu} Z^{\mu} \right],
\]

where \(g\) and \(g'\) are the SU(2)\textsubscript{L} and U(1)\textsubscript{Y} gauge couplings of the SM, respectively. The covariant derivative and the SU(2) generators for the triplet representation are given in Appendix A. The squared masses of the \(W\) and \(Z\) bosons are then given by

\[
M^2_W = \frac{g^2}{4} (v_\phi^2 + 4v_\xi^2 + 2v_\chi^2) \equiv \frac{g^2}{4} v^2_{\text{SM}}, \quad
M^2_Z = \frac{g^2 + g'^2}{4} (v_\phi^2 + 4v_\chi^2),
\]

where we define \(v^2_{\text{SM}} = v_\phi^2 + 4v_\xi^2 + 2v_\chi^2 \simeq (246\ \text{GeV})^2\).

The different form of the triplet contributions to the \(W\) and \(Z\) masses results in a tree-level \(\rho\) parameter different from 1,

\[
\rho \equiv \frac{M^2_W}{M^2_Z \cos^2 \theta_W} = \frac{v_\phi^2 + 4v_\xi^2 + 2v_\chi^2}{v_\phi^2 + 4v_\chi^2} = 1 + \frac{4v_\xi^2 - 2v_\chi^2}{v_\phi^2 + 4v_\chi^2} \equiv 1 + \Delta \rho,
\]

where we used the definition \(\cos \theta_W = g/\sqrt{g^2 + g'^2}\). The \(\rho\) parameter also receives radiative corrections in the SM from isospin violation, depending mostly on the top quark and SM
Higgs masses. After separating out the SM contributions, the constraint on new sources of SU(2) breaking is given at the 2σ level by \[ \Delta \rho = 4^{+27}_{-7} \times 10^{-4} \] (2σ constraint). \[ (5) \]

This leads to the requirement \( v_\xi \simeq v_\chi / \sqrt{2} \) to high precision. This relation can be enforced by imposing an SU(2)\(_R\) global symmetry on the Higgs potential, resulting in a residual global “custodial” SU(2) symmetry—the diagonal subgroup of SU(2)\(_L \times SU(2)\(_R\)—after electroweak symmetry breaking \[10–12\]. The SU(2)\(_R\) global symmetry is broken by hypercharge (for the scalars, hypercharge corresponds to the \( T_3 \) generator of SU(2)\(_R\)), so that SU(2)\(_R\)-violating counterterms must be introduced at one-loop level \[13\]. Nevertheless, if the scale where SU(2)\(_R\) is a good symmetry is not too high, the approximate global symmetry protects the \( \rho \) parameter and has been used, e.g., to control the effects of triplets on electroweak precision observables in little Higgs models \[16\].

The imposition of custodial SU(2) on the Higgs potential results in Higgs mass eigenstates after electroweak symmetry breaking that lie in multiplets of the custodial SU(2) with common masses \[10\]. In particular, the physical states form two singlets \( H_0^0 \) and \( H_0^{0'} \), a 3-plet \( H_3^{+,0,-} \), and a 5-plet \( H_5^{+++,+,0,-,-} \) under custodial SU(2), with the CP-even neutral states given by \[10\]

\[
H_1^0 = \phi^{0,r}, \quad H_1^{0'} = \frac{1}{\sqrt{3}}(\xi^0 + \sqrt{2}\chi^{0,r}), \quad H_5^0 = \frac{1}{\sqrt{3}}(\sqrt{2}\xi^0 - \chi^{0,r}). \] \[ (6) \]

If the custodial SU(2) is exact, the two singlets \( H_1^0 \) and \( H_1^{0'} \) can mix, but \( H_5^0 \) must be a mass eigenstate.

In this framework, the couplings of the CP-even neutral states to \( W, Z \), and fermion pairs are given by \[10, 12\]

\[
\bar{g}_{H_1^0 WW} = c_H, \quad \bar{g}_{H_1^0 ZZ} = c_H, \quad \bar{g}_{H_1^0 ff} = \frac{1}{c_H},
\]

\[
\bar{g}_{H_1^{0'} WW} = \frac{2\sqrt{2}}{3} s_H, \quad \bar{g}_{H_1^{0'} ZZ} = \frac{2\sqrt{2}}{3} s_H, \quad \bar{g}_{H_1^{0'} ff} = 0,
\]

\[
\bar{g}_{H_5^0 WW} = \frac{1}{3} s_H, \quad \bar{g}_{H_5^0 ZZ} = \frac{2}{3} s_H, \quad \bar{g}_{H_5^0 ff} = 0, \] \[ (7) \]

where we define \( \cos \theta_H \equiv c_H = v_\phi / v_{SM} \), \( s_H = \sqrt{1 - c_H^2} \), and the barred couplings are normalized to the corresponding SM Higgs couplings, \( \bar{g}_{Hxx} \equiv g_{Hxx} / g_{HSMxx} \).

In this paper we relax the assumption of custodial SU(2) in the Higgs potential. We still require that \( v_\xi \simeq v_\chi / \sqrt{2} \) in order to obey the constraint on the \( \rho \) parameter, but we allow
arbitrary mixing of the three CP-even neutral Higgs states. This is a fine-tuned situation, but it allows us to illustrate the full range of couplings allowed by experimental constraints. We will also present results for custodial SU(2)-preserving mixing. The most general gauge-invariant Higgs potential (without custodial SU(2)) is given in Appendix B where we show that enough parameter freedom exists to allow arbitrary mixing even after the vevs are fixed as above.

The doublet and triplet vevs can be parameterized in terms of $v_{\text{SM}}$, $c_H$, and $\Delta \rho$ according to

\begin{align*}
    v_\phi^2 &= c_H^2 v_{\text{SM}}^2, \\
    v_\xi^2 &= \frac{1}{8} \left[ s_H^2 + \frac{\Delta \rho}{1 + \Delta \rho} \right] v_{\text{SM}}^2 \approx \frac{1}{8} s_H^2 v_{\text{SM}}^2, \\
    v_\chi^2 &= \frac{1}{4} \left[ s_H^2 - \frac{\Delta \rho}{1 + \Delta \rho} \right] v_{\text{SM}}^2 \approx \frac{1}{4} s_H^2 v_{\text{SM}}^2. \quad (8)
\end{align*}

We consider two parameterizations of the mixing among the three neutral states. First, we parameterize the CP-even neutral Higgs state $h$ of interest as

\begin{equation}
    h = a \phi^0 + b \xi^0 + c \chi^0, \quad (9)
\end{equation}

where $a^2 + b^2 + c^2 = 1$. The couplings of $h$ are then given by

\begin{align*}
    \bar{g}_{hW} &= \frac{av_\phi + 4bv_\xi + 2cv_\chi}{v_{\text{SM}}}, \quad \bar{g}_{hZZ} = \frac{av_\phi + 4cv_\chi}{v_{\text{SM}}}, \quad \bar{g}_{hff} = \frac{av_{\text{SM}}}{v_\phi} = \frac{a}{c_H}. \quad (10)
\end{align*}

In the limit $\Delta \rho \to 0$, the $hWW$ and $hZZ$ couplings can be written as

\begin{equation}
    \bar{g}_{hWW} = a c_H + (\sqrt{2}b + c)s_H, \quad \bar{g}_{hZZ} = a c_H + 2cs_H. \quad (11)
\end{equation}

As a second parameterization we can write

\begin{equation}
    h = \cos \phi \left( \cos \theta H^0_1 + \sin \theta H^0_1' \right) + \sin \phi H^0_5, \quad (12)
\end{equation}

where $\theta$ parameterizes the mixing between the two custodial SU(2) singlets and $\phi$ parameterizes the mixing between the custodial SU(2) singlets and the 5-plet. These mixing angles are related to $a$, $b$, and $c$ above by

\begin{equation}
    a = \cos \phi \cos \theta, \quad b = \frac{1}{\sqrt{3}} \cos \phi \sin \theta + \sqrt{\frac{2}{3}} \sin \phi, \quad c = \sqrt{\frac{2}{3}} \cos \phi \sin \theta - \frac{1}{\sqrt{3}} \sin \phi. \quad (13)
\end{equation}

In this notation, $H^0_1$, $H^0_1'$, and $H^0_5$ correspond to $(a, b, c) = (1, 0, 0)$, $(0, \sqrt{1/3}, \sqrt{2/3})$, and $(0, \sqrt{2/3}, -\sqrt{1/3})$, respectively.
In the limit $\Delta \rho = 0$ we have for the couplings,

\[
\bar{g}_{hWW} = \cos \phi \cos \theta c_H + \cos \phi \sin \theta \frac{2\sqrt{2}}{\sqrt{3}} s_H + \sin \phi \frac{1}{\sqrt{3}} s_H,
\]

\[
\bar{g}_{hZZ} = \cos \phi \cos \theta c_H + \cos \phi \sin \theta \frac{2\sqrt{2}}{\sqrt{3}} s_H - \sin \phi \frac{2}{\sqrt{3}} s_H,
\]

\[
\bar{g}_{hf} = \cos \phi \cos \theta \frac{1}{c_H}.
\]

Though less compact, this second parameterization will be more useful for our analysis. We note in particular that when $\sin \phi = 0$, $h$ is a mixture of custodial SU(2) singlets only, and $\bar{g}_{hWW} = \bar{g}_{hZZ} \equiv \bar{g}_{hVV}$. For $\sin \phi = 0$ we have,

\[
\bar{g}_{hVV} = \cos \theta c_H + \sin \theta \frac{2\sqrt{2}}{\sqrt{3}} s_H.
\]

This coupling reaches a maximum value of

\[
\bar{g}_{hVV}^{\text{max}} = \sqrt{\frac{8 - 5c_H^2}{3}} \quad \text{when} \quad \sin \theta^{\text{max}} = \sqrt{\frac{8s_H^2}{8 - 5c_H^2}}.
\]

### III. NUMERICAL RESULTS

Imposing $\Delta \rho = 0$, the largest $hWW$ and $hZZ$ couplings occur when $v_\phi = 0$ and $v_\chi = \sqrt{2} v_\xi = v_{\text{SM}}/2$. This yields coupling maxima of $\bar{g}_{hWW} = \sqrt{3}$ for $h = \sqrt{\frac{2}{3}} \xi^0 + \frac{1}{\sqrt{3}} \chi^{0,r}$, for which $\bar{g}_{hZZ} = 2/\sqrt{3}$; and $\bar{g}_{hZZ} = 2$ for $h = \chi^{0,r}$, for which $\bar{g}_{hWW} = 1$. This situation is clearly unrealistic, however, because a nonzero value of $v_\phi$ is needed to generate fermion masses. This leads to a lower bound on $c_H$ (equivalently $v_\phi$) from processes that constrain the top quark Yukawa coupling in multi-Higgs-doublet models. In particular, the $\bar{b}b$ fraction in hadronic $Z$ decays, $R_b$, was studied in the Georgi-Machacek model in Ref. [17]. For a custodial SU(2)-symmetric spectrum, $H^\pm_3$ contributes to $R_b$ at one-loop level and leads to the constraint $c_H \geq 0.9$ (0.6) for $M_3 = 100$ (1000) GeV. If mixing between $H^\pm_3$ and $H^\pm_5$ is allowed, new contributions arise that can have either sign. In what follows we plot results for $c_H = 0.9$ and 0.95 as reasonable values. We also show the $h$ couplings to $WW$ and $ZZ$ with $c_H = 0.6$ as an extreme case.

#### A. Higgs couplings

When plotted against each other for a fixed value of $c_H$, the expressions for $\bar{g}_{hWW}$ and $\bar{g}_{hZZ}$ given in Eqs. (11) and (14) describe a filled ellipse centered at the origin as shown in...
FIG. 1. Normalized Higgs couplings $\bar{g}_{hZZ}$ versus $\bar{g}_{hWW}$ for a mixed state $h$ in the Georgi-Machacek model, for $c_H = 0.6$, 0.9, and 0.95 (left to right). We sampled both signs for $a$, $b$, and $c$ in Eq. (11).

Points corresponding to a mixture of $H_1^0$ and $H_1^{0\nu}$ lie along the diagonal $\bar{g}_{hWW} = \bar{g}_{hZZ}$. $H_5^0$ is indicated by the crossed circle on the lower right boundary of the ellipse.

Fig. 1. The ellipse is bounded by a curve corresponding to $\theta = \theta^{\text{max}}$ as defined in Eq. (16) and given parametrically by

$$
\begin{align*}
\bar{g}_{hWW} &= \cos \phi \sqrt{\frac{8 - 5c_H^2}{3}} + \sin \phi \frac{s_H}{\sqrt{3}}, \\
\bar{g}_{hZZ} &= \cos \phi \sqrt{\frac{8 - 5c_H^2}{3}} - \sin \phi \frac{2s_H}{\sqrt{3}}.
\end{align*}
$$

The bounding ellipse can be conveniently expressed by defining $g_\pm \equiv (\bar{g}_{hZZ} \pm \bar{g}_{hWW})/\sqrt{2}$; we then obtain

$$
1 = \frac{3}{2(8 - 5c_H^2)} g_+^2 - \frac{1}{(8 - 5c_H^2)} g_+ g_- + \frac{(11 - 7c_H^2)}{2s_H^2(8 - 5c_H^2)} g_-^2.
$$

Note that this ellipse is not oriented along the diagonal $\bar{g}_{hWW} = \bar{g}_{hZZ}$, but instead has its semimajor axis tilted at a small angle $\varphi$ above the diagonal, where

$$
\tan 2\varphi = \frac{s_H^2}{2(1 + s_H^2)}.
$$

In the custodial SU(2)-symmetric case, $H_5^0$ must be a mass eigenstate but $H_1^0$ and $H_1^{0\nu}$ can mix. Points corresponding to $h = \cos \theta H_1^0 + \sin \theta H_1^{0\nu}$ fall along the diagonal, $\bar{g}_{hWW} = \bar{g}_{hZZ}$. The normalized couplings of $H_5^0$ are given in Table II. In particular, $\bar{g}_{H_5^0 ZZ} = -2\bar{g}_{H_5^0 WW}$. The $h = H_5^0$ point falls on the boundary of the ellipses in Fig. 1, marked with a crossed circle.
TABLE I. Normalized couplings of $H_0^0$ to $WW$ and $ZZ$ for various values of $c_H$.

| $c_H$ | $g_{H_{WW}}^2$ | $g_{H_{ZZ}}^2$ |
|-------|----------------|----------------|
| 0.6   | 0.462          | -0.924         |
| 0.9   | 0.252          | -0.503         |
| 0.95  | 0.180          | -0.361         |

B. Higgs production and decay rates

We now compute the rates for $gg \rightarrow h \rightarrow WW$ and $gg \rightarrow h \rightarrow ZZ$, normalized to their SM values. In the SM, these processes constitute the most promising discovery modes for an intermediate-mass Higgs. Our calculation is valid for any hadron collider, because dependence on the initial-state parton densities and collider center-of-mass energy cancel in the ratio with the SM rate. We assume throughout that the narrow-width approximation can be applied. For $VV = WW$ or $ZZ$ we have,

$$\frac{\sigma(gg \rightarrow h \rightarrow VV)}{\sigma_{\text{SM}}(gg \rightarrow H \rightarrow VV)} = \frac{\sigma(gg \rightarrow h)}{\sigma_{\text{SM}}(gg \rightarrow H)} \frac{\text{BR}(h \rightarrow VV)}{\text{BR}_{\text{SM}}(H \rightarrow VV)}.$$  \hspace{1cm} (20)

The $gg \rightarrow h$ production process proceeds through the coupling of $h$ to a fermion loop and is dominated in the SM by the top quark contribution. Because all fermion couplings are scaled by the same factor $\bar{g}_{hff}$ in the Georgi-Machacek model, we have\(^3\)

$$\frac{\sigma(gg \rightarrow h)}{\sigma_{\text{SM}}(gg \rightarrow H)} = \bar{g}_{hff}^2.$$  \hspace{1cm} (21)

To compute the decay branching ratios in the Georgi-Machacek model, we start with the partial widths of the SM Higgs of the same mass, computed using the public FORTRAN code HDECAY version 3.531 \cite{19}. The corresponding partial widths of $h$ are computed as follows:

$$\Gamma(h \rightarrow WW) = \bar{g}_{h_{WW}}^2 \Gamma_{\text{SM}}(H \rightarrow WW), \quad \Gamma(h \rightarrow ZZ) = \bar{g}_{h_{ZZ}}^2 \Gamma_{\text{SM}}(H \rightarrow ZZ)$$

$$\Gamma(h \rightarrow f\bar{f}) = \bar{g}_{h_{ff}}^2 \Gamma_{\text{SM}}(H \rightarrow f\bar{f}), \quad \Gamma(h \rightarrow gg) = \bar{g}_{h_{ff}}^2 \Gamma_{\text{SM}}(H \rightarrow gg),$$  \hspace{1cm} (22)

\(^3\) At two-loop order, there are electroweak contributions to the $gg \rightarrow h$ amplitude in which $h$ couples to $WW$ or $ZZ$ inside the loop, which should instead be scaled by $\bar{g}_{h_{WW}}$ or $\bar{g}_{h_{ZZ}}$, respectively. In the SM, these contributions result in a 5–6% positive correction to the cross section \cite{18}. Ignoring their different coupling dependence thus results in an error of order 0.06 $\bar{g}_{h_{ff}}(\bar{g}_{h_{ff}} - \bar{g}_{h_{WW},ZZ})$ on the right-hand side of Eq. (21).
FIG. 2. The rate for $gg \to h \to WW$ normalized to its SM value plotted as a function of $\sigma_{hWW}^2$, for $M_h = 140, 160$ and $180$ GeV (left to right), with $c_H = 0.9$. The solid line shows points for which $h = \cos \theta H_1^0 + \sin \theta H_1^0$ and the crossed circle near the origin indicates the point corresponding to $H_5^0$. The dashed horizontal line shows the Tevatron upper limit on the rate for $gg \to h \to WW$ for the corresponding Higgs mass from Ref. [20].

where $f f'$ refers to any fermion pair. The remaining partial widths, $\Gamma(h \to \gamma\gamma)$ and $\Gamma(h \to \gamma Z)$, constitute a very small portion of the $h$ total width; for simplicity we ignore the corrections to these modes and instead take them equal to the corresponding SM Higgs widths. The branching ratios of $h$ are then given by taking the ratio of the relevant partial width to the $h$ total width.

We plot $\sigma(gg \to h \to WW)/\sigma_{SM}$ in Figs. [2][3] and $\sigma(gg \to h \to ZZ)/\sigma_{SM}$ in Figs. [6][9]. We show results for $M_h = 140, 160$, and $180$ GeV and $c_H = 0.9$ and $0.95$. The points correspond to a general mixed state $h$, while the solid line shows the custodial SU(2)-preserving mixture $h = \cos \theta H_1^0 + \sin \theta H_1^0$. For $h = H_5^0$, the $gg \to h$ cross section is zero because $H_5^0$ does not couple to fermions; this is shown in the plots by a crossed circle.

The rate for $gg \to h \to WW$ can be enhanced by up to about 20% (10%) compared to the corresponding SM rate for $c_H = 0.9$ (0.95), as shown in Table [11]. The enhancement is due to an interplay between the production and decay couplings. For $c_H = 0.9$ (0.95), the $gg \to h$ production cross section is about $1.23a^2$ ($1.11a^2$) times its SM value, for $a$ defined in Eq. (9). The branching ratio for $h \to WW$ depends mainly on competition between the $h \to WW$ and $h \to bb$ partial widths, especially for lower Higgs masses.

The error introduced by this assumption in the $h$ branching ratio calculation is at the level of $10^{-3}(1 - \tilde{g}_{ff,hWW}^2)$. 
FIG. 3. As in Fig. 2 but for $c_H = 0.95$.

FIG. 4. The rate for $gg \to h \to WW$ normalized to its SM value plotted as a function of $g_{HZZ}^2$, for $M_h = 140$, 160 and 180 GeV (left to right), with $c_H = 0.9$. The solid line shows points for which $h = \cos \theta H_1^0 + \sin \theta H_1^0$ and the crossed circle at the lower edge of the allowed region indicates the point corresponding to $H_1^0$. The dashed horizontal line shows the Tevatron upper limit on the rate for $gg \to h \to WW$ for the corresponding Higgs mass from Ref. [20].

FIG. 5. As in Fig. 4 but for $c_H = 0.95$. 
FIG. 6. The rate for $gg \rightarrow h \rightarrow ZZ$ normalized to its SM value plotted as a function of $g_{hWW}^2$, for $M_h = 140, 160$, and $180$ GeV (left to right), with $c_H = 0.9$. The solid line shows points for which $h = \cos \theta_{H_1} + \sin \theta_{H_1}$ and the crossed circle near the origin indicates the point corresponding to $H_5^0$. Points excluded by the Tevatron in Figs. 2 and 4 are not shown, leading to the gap in the solid line for $M_h = 160$ GeV.

FIG. 7. As in Fig. 6 but for $c_H = 0.95$.

For $c_H = 0.9$ and $M_h = 160$ GeV, the rate for $gg \rightarrow h \rightarrow WW$ can be enhanced enough to exceed the current limit on this cross section from the Tevatron $^{20}$. Our calculation of the upper limit is given in Table III. Points excluded by the Tevatron are not shown in the plots of $gg \rightarrow h \rightarrow ZZ$ in Figs. 6 and 7.

The rate for $gg \rightarrow h \rightarrow ZZ$ can be enhanced much more dramatically, by more than a factor of 5 (3) for $M_h = 180$ (160) GeV and $c_H = 0.9$. Details are given in Table IV. This

Note that the Tevatron analysis in Ref. $^{20}$ selected specifically for the $gg \rightarrow H \rightarrow WW$ process, and is not a combination of multiple SM Higgs search channels. The resulting cross section limits are thus directly applicable to the $gg \rightarrow h \rightarrow WW$ rates found here.
FIG. 8. The rate for $gg \to h \to ZZ$ normalized to its SM value plotted as a function of $\bar{g}^2_{hZZ}$, for $M_h = 140, 160, \text{and } 180 \text{ GeV}$ (left to right), with $c_H = 0.9$. The solid line shows points for which $h = \cos \theta H^0_1 + \sin \theta H^{0*}_1$ and the crossed circle at the lower edge of the allowed region indicates the point corresponding to $H^0_5$. Points excluded by the Tevatron in Figs. 2 and 4 are not shown, leading to the gap in the solid line for $M_h = 160 \text{ GeV}$.

FIG. 9. As in Fig. 8 but for $c_H = 0.95$.

large enhancement occurs when the $hWW$ coupling is near zero but the $hZZ$ coupling is non-negligible, as can be seen by comparing Figs. 6 and 7 and Figs. 8 and 9. This suppresses the dominant Higgs partial width to $WW$, resulting in $BR(h \to ZZ)$ close to 1 for the heavier masses. Clearly, this can only happen if $\bar{g}_{hWW} \neq \bar{g}_{hZZ}$, so it requires mixing between $H^0_5$ and the custodial SU(2)-singlet states. This would lead to an early LHC discovery of such a Higgs in the “golden mode,” $h \to ZZ \to 4$ leptons, while the $h \to WW$ signal (which would provide the first discovery in the SM) would be absent or very suppressed.

If $h = \cos \theta H^0_1 + \sin \theta H^{0*}_1$, the enhancement in the $ZZ$ channel is limited to about 20%
Maximum of $\sigma(gg \to h \to WW)/\sigma_{SM}$

| $c_H$ | $M_h$ [GeV] | Overall | $h \sim H_1^0, H_1^{0'}$ |
|-------|-------------|---------|-------------------------|
| 0.9   | 140         | 1.1581  | 1.1352                  |
|       | 160         | 1.2085  | 1.2056                  |
|       | 180         | 1.2327  | 1.2296                  |
| 0.95  | 140         | 1.0777  | 1.0656                  |
|       | 160         | 1.0971  | 1.0962                  |
|       | 180         | 1.1069  | 1.1061                  |

TABLE II. Maximum values of $\sigma(gg \to h \to WW)/\sigma_{SM}$ for $c_H = 0.9$ and 0.95 and $M_h = 140$, 160, and 180 GeV. The third column gives the maximum cross section enhancement obtainable when $h$ is any linear combination of $H_1^0$, $H_1^{0'}$, and $H_0^5$, while the fourth column gives the maximum when $h$ is a linear combination of only $H_1^0$ and $H_1^{0'}$.

| $M_h$ [GeV] | $\sigma_{SM}(gg \to H)$ [18] | $\sigma_{SM}(gg \to H)/\sigma_{SM}$ | $\text{BR}_{SM}(H \to WW)$ [19] | $\text{BR}_{SM}(H \to WW)/\sigma_{SM}$ | Tevatron upper limit on $\sigma(gg \to h \to WW)$ [20] | $\text{Upper limit on } \sigma(gg \to h \to WW)/\sigma_{SM}$ |
|-------------|-------------------------------|-------------------------------------|----------------------------------|--------------------------------|----------------------|----------------------------------|
| 140         | 0.680 pb                      | 1.29 pb                            | 0.4916                           | 0.29                                  | 3.86                 | 1.20                              |
| 160         | 0.434 pb                      | 0.47 pb                            | 0.9048                           | 0.47                                  | 1.20                 | 1.58                              |
| 180         | 0.279 pb                      | 0.41 pb                            | 0.9325                           | 0.41                                  | 1.58                 |                                   |

TABLE III. The $gg \to H$ cross section and $H \to WW$ branching fraction in the SM, Tevatron upper limit on the $gg \to h \to WW$ cross section, and resulting upper limit on the cross section normalized to its SM value, for $M_h = 140$, 160, and 180 GeV. The uncertainty on $\sigma_{SM}(gg \to H)$ from the parton density functions is $\pm8$–10% [18]; here we have taken the cross section central value for our limit.

(10%) for $c_H = 0.9$ (0.95). Note that when $h = \cos \theta H_1^0 + \sin \theta H_1^{0'}$, the enhancement in the $WW$ and $ZZ$ channels is identical, i.e., $\sigma(gg \to h \to WW)/\sigma_{SM} = \sigma(gg \to h \to ZZ)/\sigma_{SM}$. Such an enhancement would lead to a marginally earlier discovery of $h$ in both channels at the LHC. For $c_H = 0.95$, however, the enhancement is about the same size as the current uncertainty in $\sigma_{SM}(gg \to H)$ from the parton density functions [18].
Maximum of $\sigma(gg \rightarrow h \rightarrow ZZ)/\sigma_{SM}$

| $c_H$ | $M_h$ [GeV] | Overall | $h \sim H_1^0, H_1^{0\prime}$ |
|-------|-------------|---------|-------------------------------|
| 140   | 1.5068      | 1.1356  |
| 0.9   | 160         | 3.6616  | 1.1978                        |
| 180   | 5.7291      | 1.2297  |
| 140   | 1.2507      | 1.0659  |
| 0.95  | 160         | 1.9716  | 1.0962                        |
| 180   | 3.0156      | 1.1061  |

TABLE IV. As in Table III but for $\sigma(gg \rightarrow h \rightarrow ZZ)/\sigma_{SM}$.

### IV. HIGGS COUPLING EXTRACTION AND THE TRIPLET NATURE OF $h$

In models containing only Higgs doublets and singlets, the $hVV$ couplings satisfy

$$\bar{g}_{hWW} = \bar{g}_{hZZ} \equiv \bar{g}_{hVV}, \quad |\bar{g}_{hVV}| \leq 1.$$  \hspace{1cm} (23)

These relations are assumed to hold in the Higgs coupling fits of Refs. [7–9]; however, both of them are violated in the Georgi-Machacek model.

The relation $\bar{g}_{hWW} = \bar{g}_{hZZ}$ is easy to test by taking the ratio of any two Higgs signal rates with the same production mechanism and decays to $WW$ and $ZZ$, respectively:

$$\frac{\sigma(X \rightarrow h \rightarrow WW)/\sigma_{SM}(X \rightarrow h \rightarrow WW)}{\sigma(X \rightarrow h \rightarrow ZZ)/\sigma_{SM}(X \rightarrow h \rightarrow ZZ)} = \frac{\bar{g}_{hWW}^2}{\bar{g}_{hZZ}^2}.$$  \hspace{1cm} (24)

Determining whether $\bar{g}_{hVV}$ is greater than 1 is more complicated, and is not always possible at the LHC. The approach is as follows. One of the key ingredients in the Higgs coupling fits is the measurement of the rate for Higgs production in vector boson fusion ($WW, ZZ \rightarrow h$) followed by decay to $WW$. This rate can be written as

$$\sigma(VBF \rightarrow h \rightarrow WW) = \sigma(VBF \rightarrow h) \cdot BR(h \rightarrow WW).$$  \hspace{1cm} (25)

The usual Higgs coupling fit involves assuming that $\sigma(VBF \rightarrow h)$ is no larger than its SM value, i.e., $\sigma_{\text{max}}(VBF \rightarrow h) = \sigma_{SM}(VBF \rightarrow H)$. Together with the rate measurement, this sets a lower bound on $BR(h \rightarrow WW)$:

$$BR_{\text{min}}(h \rightarrow WW) = \frac{\sigma(VBF \rightarrow h \rightarrow WW)}{\sigma_{\text{max}}(VBF \rightarrow h)}.$$  \hspace{1cm} (26)
If the assumption that $\sigma(\text{VBF} \to h) \leq \sigma_{\text{SM}}(\text{VBF} \to H)$ is false, the resulting value of $\text{BR}^\text{min}(h \to WW)$ can be greater than one, which is clearly unphysical. In particular, for $c_H = 0.9 \ (0.95)$, we find that $\sigma(\text{VBF} \to h)$ can be as much as 31% \ (16%) larger than in the SM. An unphysical $\text{BR}^\text{min}(h \to WW)$ would then be obtained using Eq. (26) if the true value of $\text{BR}(h \to WW)$ is greater than 0.76 \ (0.86), which is true in the SM for $M_H$ in the range 154–192 GeV \ (159–184 GeV).

This range can be extended by considering Higgs decays to other final states as follows. If the Higgs is detectable in the $WW$ final state and some other final state ($ZZ$, $\tau\tau$, etc.) via a common production mechanism, we can write

$$\frac{\text{BR}(h \to ZZ)}{\text{BR}(h \to WW)} = \frac{\sigma(X \to h \to ZZ)}{\sigma(X \to h \to WW)}, \quad \frac{\text{BR}(h \to \tau\tau)}{\text{BR}(h \to WW)} = \frac{\sigma(X \to h \to \tau\tau)}{\sigma(X \to h \to WW)},$$

etc. The total branching fraction of the Higgs to detected modes is then given by

$$\text{BR}(h \to \text{detected}) = \text{BR}(h \to WW) \left[ 1 + \frac{\text{BR}(h \to ZZ)}{\text{BR}(h \to WW)} + \frac{\text{BR}(h \to \tau\tau)}{\text{BR}(h \to WW)} + \cdots \right].$$

If inserting $\text{BR}^\text{min}(h \to WW)$ from Eq. (26) on the right-hand side yields a value of $\text{BR}(h \to \text{detected})$ greater than 1, then we can conclude that the assumption that $\sigma(\text{VBF} \to h) \leq \sigma_{\text{SM}}(\text{VBF} \to H)$ is false.

If the detectable branching fraction of the Higgs is sufficiently small, the technique discussed above will fail. In that case, detection of the triplet nature of $h$ at the LHC would rely on the discovery of additional Higgs states. Search prospects at the LHC for single production of $\chi^{++}$ with decays to $W^+W^+$ were studied in Refs. [21, 22], and for pair production of $\chi^{++}\chi^{-+}$ again with decays to like-sign $W$ bosons in Ref. [23]. Single production of $H^+_5 \equiv (\chi^+ - \xi^+)/\sqrt{2}$ through $W^+Z$ fusion with decays back to $W^+Z$ was also studied in Refs. [22, 24].

We note that the technique discussed here can yield only a lower bound on $\sigma(\text{VBF} \to h)$. Furthermore, removing the SM assumption that $\sigma(\text{VBF} \to h) \leq \sigma_{\text{SM}}(\text{VBF} \to H)$ reopens the parameter degeneracy in the Higgs coupling fit discussed in the introduction. While ratios of Higgs couplings would still be measurable at the LHC, knowledge of absolute Higgs couplings would have to wait for direct model-independent measurements at an $e^+e^-$ collider [25].

Reference [23] also considered $\chi^{++}$ decays to like-sign charged leptons via lepton-number-violating Yukawa couplings that could be responsible for neutrino masses. For the triplet vevs that we consider here, $v_{\chi} = 54 \ (38)$ GeV for $c_H = 0.9 \ (0.95)$, such Yukawa couplings would be of order $10^{-11}$ and dilepton decays would be totally negligible.
V. CONCLUSIONS

We studied the couplings of a CP-even neutral Higgs state $h$ in the Georgi-Machacek model with Higgs triplets. We found that the effect of the triplet components of $h$ on its couplings to $WW$ and $ZZ$ can be significant, even for relatively small triplet vevs consistent with the experimental upper bounds on the top quark Yukawa coupling.

If $h$ is a mixture of only custodial SU(2) singlets, the $hWW$ and $hZZ$ couplings can be enhanced by up to 15% (8%) for $c_H = 0.9$ (0.95), but the ratio of these couplings is the same as in the SM. If $h$ is a generic mixture of the CP-even neutral states, the ratio of its $WW$ and $ZZ$ couplings can be very different than in the SM; in fact, one of these couplings can be zero while the other is finite. Such a suppression of the $hWW$ coupling while the $hZZ$ coupling is nonzero can lead to large enhancements in the rate for $gg \to h \to ZZ$, especially for $M_h \gtrsim 160$ GeV where the SM Higgs width is dominated by decays to $WW$.

While a ratio of $hWW$ and $hZZ$ couplings different than in the SM would be easy to measure by taking ratios of rates in these channels, an overall enhancement in the $hVV$ coupling could only be detected at the LHC if it leads to a rate in vector boson fusion with decays to $WW$ too large to be consistent with the SM assumption for this coupling.

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Appendix A: Covariant derivative and SU(2) generators for triplets

The gauge-covariant derivative in Eq. (2) is given by

$$D_\mu = \partial_\mu + igT^a W^a_\mu + ig' Y B_\mu.$$  \hspace{1cm} (A1)

The SU(2) generators for a doublet representation are given by $T^a = \sigma^a/2$, where $\sigma^a$ are the $2 \times 2$ Pauli matrices. The SU(2) generators for a triplet representation are

$$T^1 = \frac{1}{\sqrt{2}} \begin{pmatrix} 0 & 1 & 0 \\ 1 & 0 & 1 \\ 0 & 1 & 0 \end{pmatrix}, \quad T^2 = \frac{1}{\sqrt{2}} \begin{pmatrix} 0 & -i & 0 \\ i & 0 & -i \\ 0 & i & 0 \end{pmatrix}, \quad T^3 = \begin{pmatrix} 1 & 0 & 0 \\ 0 & 0 & 0 \\ 0 & 0 & -1 \end{pmatrix}.$$  \hspace{1cm} (A2)
Appendix B: Scalar potential and mass matrix

Reference [13] presented the most general gauge-invariant scalar potential for this model containing only terms with an even number of fields. While not completely general, this restricted potential captures most of the physics [11]. In our notation it is,

\[
V_{\text{even}} = \mu_1^2 \Phi^\dagger \Phi + \mu_2^2 X^\dagger X + \mu_3^2 \Xi^\dagger \Xi \\
+ \lambda_1 (\Phi^\dagger \Phi)^2 + \lambda_2 |X^T C X|^2 + \lambda_3 (\Phi^\dagger T^a \Phi)(X^\dagger T^a X) \\
+ \lambda_4 \left[ (\Phi^\dagger T^a \Phi^c)(\Xi^\dagger T^a X) + \text{h.c.} \right] + \lambda_5 (\Phi^\dagger \Phi)(X^\dagger X) + \lambda_6 (\Phi^\dagger \Phi)(\Xi^\dagger \Xi) \\
+ \lambda_7 (X^\dagger X)^2 + \lambda_8 (\Xi^\dagger \Xi)^2 + \lambda_9 |X^\dagger \Xi|^2 + \lambda_10 (\Xi^\dagger \Xi)(X^\dagger X),
\]

(B1)

where \( \lambda_i \) are dimensionless and \( \mu_i^2 \) have dimension mass squared. Here \( \Phi^c = i\sigma_2 \Phi^* \) and

\[
C = \begin{pmatrix}
0 & 0 & 1 \\
0 & -1 & 0 \\
1 & 0 & 0
\end{pmatrix}
\]

(B2)
is the conjugation operator for the triplet fields such that \( X^c = CX^* \) transforms as a triplet of SU(2).

We add the most general gauge-invariant set of trilinear terms,

\[
V_{\text{trilinear}} = \kappa_1 [(\Phi^\dagger T^a \Phi^c)X^a + \text{h.c.}] + \kappa_2 (\Phi^\dagger T^a \Phi)\Xi^a + \kappa_3 X^{*a}(i\epsilon^{abc} X^b \Xi^c),
\]

(B3)

where \( \kappa_i \) have dimensions of mass and \( \epsilon^{abc} \) is the totally antisymmetric tensor.

Minimizing the potential allows us to eliminate \( \mu_i^2 \) in favor of the scalar vevs:

\[
\mu_1^2 = -\kappa_1 v_\chi + \kappa_2 \frac{v_\xi}{2} - \lambda_1 v_\phi^2 - \left( \frac{\lambda_3}{4} + \frac{\lambda_5}{2} \right) v_\phi^2 - \lambda_4 v_\chi v_\xi - \lambda_6 v_\xi^2 \\
\mu_2^2 = -\kappa_1 \frac{v_\phi^2}{2 v_\chi} + \kappa_3 v_\xi - \left( \frac{\lambda_3}{4} + \frac{\lambda_5}{2} \right) v_\phi^2 - \lambda_4 \frac{v_\phi^2 v_\xi}{2 v_\chi} - \lambda_7 v_\chi^2 - \lambda_10 v_\xi^2 \\
\mu_3^2 = \kappa_2 \frac{v_\phi^2}{8 v_\xi} + \kappa_3 \frac{v_\chi^2}{4 v_\xi} - \lambda_4 \frac{v_\phi^2 v_\chi}{4 v_\xi} - \lambda_6 \frac{v_\phi^2}{2} - 2\lambda_8 v_\xi^2 - \lambda_10 \frac{v_\chi^2}{2}.
\]

(B4)

Applying these minimization conditions, the quadratic terms in the potential involving CP-even neutral fields are,

\[
V_{m^2 \text{ terms}} = (\phi^{0,r})^2 \left[ \lambda_1 v_\phi^2 \right] + (\chi^{0,r})^2 \left[ -\kappa_1 \frac{v_\phi^2}{4 v_\chi} - \lambda_4 \frac{v_\phi^2 v_\xi}{4 v_\chi} + \lambda_7 v_\xi^2 \right]
\]

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The last two terms, representing mixing between \( H_H \) and \( H_R \) to be invariant under a global SU(2) violate the custodial SU(2). Imposing the conditions given in Ref. [13] for the scalar potential \( \rho \) for the vevs, and setting \( \Delta \rho = 0 \), we obtain

\[
V_{\text{in\_terms}} = v_{\text{SM}}^2 \left\{ (H_1^0)^2 \lambda_1 c_H^2 + \frac{(H_1^0)^2}{2} - \frac{(H_2^0)^2}{6} \right\}
\]

Rewriting the fields in the \( H_1^0, H_2^0, H_3^0 \) basis given in Eq. (6), using the notation of Eq. (8) for the vevs, and setting \( \Delta \rho = 0 \), we obtain

\[
V_{\text{in\_terms}} = v_{\text{SM}}^2 \left\{ (H_1^0)^2 \lambda_1 c_H^2 + \frac{(H_1^0)^2}{6} - \frac{(H_2^0)^2}{2} - \frac{(H_3^0)^2}{2} \right\}
\]

The last two terms, representing mixing between \( H_1^0 \) and \( H_3^0 \) and \( H_2^0 \) and \( H_3^0 \), respectively, violate the custodial SU(2). Imposing the conditions given in Ref. [13] for the scalar potential to be invariant under a global SU(2)\(_R\) causes the \( \lambda_i \) contributions of these last two terms to vanish. When \( \kappa_2 = -\sqrt{2} \kappa_1 \) the remaining contribution vanishes as well.

Even if the custodial SU(2)–violating couplings are small, the mixing between the singlets and \( H_3^0 \) can be large if their masses are degenerate enough.

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\(^7\) The relevant conditions are \( \lambda_3 = \sqrt{2} \lambda_4, \lambda_5 = 2 \lambda_6, 2 \lambda_7 = \lambda_9 + 2 \lambda_{10}, \) and \( 4 \lambda_8 = \lambda_9 + \lambda_{10} \).

\(^8\) The near-degenerate case was studied in a model with one doublet and one triplet in Ref. 26.
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