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THE SCALAR SECTOR IN THE GEORGI–MACHACEK MODEL[1]

The Georgi–Machacek (GM) model extends the Standard Model (SM) Higgs sector by adding one complex and one real scalar triplet, while preserving a custodial $SU(2)_V$ symmetry. This setup predicts a rich scalar spectrum including a quintet, a triplet, and two CP-even singlets. In this article, we review the model structure, its mass spectrum, and the theoretical and experimental constraints from perturbativity, vacuum stability, electroweak precision tests, Higgs measurements, and direct searches at colliders in both cases where the SM-like Higgs should be the light or heavy CP-even eigenstate. We also discuss the model's capability to explain the 95 GeV excess reported in $\gamma\gamma$, $\tau\tau$ and $b\bar{b}$ channels, by assuming two degenerate (CP-even and CP-odd) resonances. We demonstrate that the nature of the 95 GeV scalar resonance candidate can be probed via the properties of its di-τ decay.

Keywords: GM model, Higgs, 95 GeV excess, strength modifiers.

1. Introduction

The discovery of the Higgs boson at the Large Hadron Collider (LHC) with a mass of about 125 GeV marked a major milestone in particle physics, confirming the mechanism of electroweak symmetry breaking (EWSB) in the Standard Model (SM) [1, 2]. Nevertheless, several questions remain open, particularly regarding the nature of the EWSB sector and the possible existence of additional scalar states. Experimental searches for new resonances, both lighter and heavier than the observed Higgs boson, are actively pursued at the LHC, as they could signal physics beyond the SM (BSM) [3, 4].

Interestingly, several excesses in the data have been reported around 95 GeV in different final states: a diphoton excess with a local significance of about 3.2σ [5, 6], an excess in the $\tau\tau$ channel [9], and an excess in the $b\bar{b}$ channel from LEP data [10]. These persistent anomalies, though not yet conclusive, motivate study of BSM scenarios that can accommodate a light scalar resonance while remaining consistent with all existing theoretical and experimental constraints.

A well-motivated extension of the SM Higgs sector is the Georgi–Machacek (GM) model [11], which introduces one complex and one real $SU(2)_L$ triplet while preserving a custodial $SU(2)_V$ symmetry at tree level. This structure leads to a rich scalar spectrum comprising a quintet, a triplet, and two CP-even singlets. The model has been extensively studied in the literature [12–16], and its phenomenology, including Higgs couplings, electroweak precision tests, vacuum stability, and collider signatures, has been carefully analyzed. In particular, the possibility that the observed 125 GeV Higgs boson could be either the lighter or the heavier CP-even eigenstate has been explored [17, 18].

In this work, we review the scalar sector of the GM model and examine its capability to explain the 95 GeV excesses in the $\gamma\gamma$, $\tau\tau$, and $b\bar{b}$ channels. We consider both the "single peak" scenario, where only one CP-even scalar η is responsible for the signal, and the "twin peak" scenario, where two nearly degenerate resonances, the CP-even η and the CP-odd H_3^0, contribute coherently. We show that the latter case can

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naturally accommodate the observed pattern of signal strengths, especially in the di-tau channel, which could serve as a probe of the CP structure of the resonance(s). We also discuss how future measurements of the acoplanarity angle in $\tau^+\tau^-$ decays can discriminate between the two possibilities.

In Section 2, the model is presented where the mass spectrum is defined and the different constraints are discussed. The 95 GeV excess is addressed in Section 3, where the suggested explanation within the GM model is discussed, and in Section 4, the numerical results are presented and discussed. Our conclusions are addressed in Section 5.

2. Model, Mass Spectrum & Constraints

The scalar sector of the GM model consists of a scalar doublet $(\phi^+, \phi^0)^T$ with hypercharge $Y = 1$; and two triplet representations with hypercharges $Y = 2{,}0$, respectively. These representations can be written as

$$\Phi = \begin{pmatrix} \phi^{0*} & \phi^+ \\ -\phi^{+*} & \phi^0 \end{pmatrix}, \quad \Delta = \begin{pmatrix} \chi^{0*} & \xi^+ & \chi^{++} \\ -\chi^{+*} & \xi^0 & \chi^+ \\ \chi^{++*} & -\xi^{+*} & \chi^0 \end{pmatrix}, \quad (1)$$

with $\phi^- = \phi^+$, $\xi^- = \xi^+$, $\chi^{--} = \chi^{++}$, $\chi^- = \chi^+$. The neutral components in (1) can be expressed as

$$\begin{aligned} \phi^0 &= \frac{1}{\sqrt{2}}(\upsilon_\phi + h_\phi + ia_\phi),\ \chi^0 = \\ &= \frac{1}{\sqrt{2}}(\upsilon_\chi + h_\chi + ia_\chi), \\ \xi^0 &= \upsilon_\xi + h_\xi, \end{aligned} \quad (2)$$

where υ_ϕ, υ_χ and υ_ξ are the VEVs for ϕ^0, χ^0 and ξ^0, respectively. Here, we have three CP-even scalar degrees of freedom (dofs) $\{h_\phi, h_\chi, h_\xi\}$, two CP-odd dof's $\{a_\phi, a_\chi\}$, six singly charged dof's $\{\phi^\pm, \chi^\pm, \xi^\pm\}$ and two doubly charged dof's $\chi^{\pm\pm}$. The most general scalar potential invariant under the global symmetry $SU(2)_L \times SU(2)_R \times U(1)_Y$ is given by

$$\begin{aligned} V(\Phi, \Delta) &= \frac{m_1^2}{2}\mathrm{Tr}[\Phi^\dagger\Phi] + \frac{m_2^2}{2}\mathrm{Tr}[\Delta^\dagger\Delta] + \lambda_1(\mathrm{Tr}[\Phi^\dagger\Phi])^2 + \\ &+ \lambda_2\mathrm{Tr}[\Phi^\dagger\Phi]\mathrm{Tr}[\Delta^\dagger\Delta] + \lambda_3\mathrm{Tr}[(\Delta^\dagger\Delta)^2] + \lambda_4(\mathrm{Tr}[\Delta^\dagger\Delta])^2 - \\ &- \lambda_5\mathrm{Tr}\left[\Phi^\dagger\frac{\sigma^a}{2}\Phi\frac{\sigma^b}{2}\right]\mathrm{Tr}[\Delta^\dagger T^a \Delta T^b] - \\ &- \mu_1\mathrm{Tr}\left[\Phi^\dagger\frac{\sigma^a}{2}\Phi\frac{\sigma^b}{2}\right] \times \\ &\times (U\Delta U^\dagger)_{ab} - \mu_2\mathrm{Tr}[\Delta^\dagger T^a \Delta T^b](U\Delta U^\dagger)_{ab}, \end{aligned} \quad (3)$$

where $\sigma^{1,2,3}$ are the Pauli matrices and the $T^{1,2,3}$ are the generators of the $SU(2)$ triplet representation. The matrix U can be found in [11]. The custodial symmetry condition at tree level $m_W^2 = m_Z^2 \cos^2\theta_W$ implies $\upsilon_\chi = \sqrt{2}\upsilon_\xi$ and $\upsilon_\phi^2 + 8\upsilon_\xi^2 \equiv \upsilon^2 = (246.22\ \mathrm{GeV})^2$, where m_W, m_Z and θ_W are the gauge boson masses and the Weinberg mixing angle. It is useful to introduce the parameter $t_\beta \equiv \tan\beta = 2\sqrt{2}\upsilon_\xi/\upsilon_\phi$ to describe the relations between the VEVs. By using the tadpole conditions, one can eliminate the parameters $m_{1,2}^2$. After the EWSB, the Goldstone bosons are eaten by the massive W and Z bosons, and we are left with the following mass eigenstates: three CP-even eigenstates $\{h, \eta, H_5^0\}$, one CP-odd eigenstate H_3^0, two singly charged scalars $\{H_3^\pm, H_5^\pm\}$, and one doubly charged scalar $H_5^{\pm\pm}$:

$$\begin{aligned} h &= c_\alpha h_\phi - \frac{s_\alpha}{\sqrt{3}}(\sqrt{2}h_\chi + h_\xi), \\ \eta &= s_\alpha h_\phi + \frac{c_\alpha}{\sqrt{3}}(\sqrt{2}h_\chi + h_\xi), \\ H_5^0 &= \sqrt{\frac{2}{3}}h_\xi - \sqrt{\frac{1}{3}}h_\chi, \quad H_3^0 = -s_\beta a_\phi + c_\beta a_\chi, \\ H_3^\pm &= -s_\beta\phi^\pm + c_\beta\frac{1}{\sqrt{2}}(\chi^\pm + \xi^\pm), \\ H_5^\pm &= \frac{1}{\sqrt{2}}(\chi^\pm - \xi^\pm), \quad H_5^{\pm\pm} = \chi^{\pm\pm}. \end{aligned} \quad (4)$$

The mixing angle α of the CP-even sector is defined by $\tan 2\alpha = 2M_{12}^2/(M_{22}^2 - M_{11}^2)$, where M^2 is the mass-squared matrix in the basis $\left\{h_\phi, \sqrt{\frac{2}{3}}h_\chi + \frac{1}{\sqrt{3}}h_\xi\right\}$ that is given in [18]. This allows us to write the SM-like Higgs boson and the heavy scalar (η) eigenmasses as $m_{h,\eta}^2 = \frac{1}{2}[M_{11}^2 + M_{22}^2 \mp \sqrt{(M_{11}^2 - M_{22}^2)^2 + 4(M_{12}^2)^2}]$. The other eigenmasses are

$$\begin{aligned} m_{H_3^0}^2 &= m_{H_3^\pm}^2 = m_3^2\left(\frac{\mu_1}{\sqrt{2}s_\beta\upsilon} + \frac{\lambda_5}{2}\right)\upsilon^2, \\ m_{H_5^0}^2 &= m_{H_5^\pm}^2 = m_{H_5^{\pm\pm}}^2 = m_5^2 = \\ &= \frac{\mu_1}{\sqrt{2}}\frac{c_\beta^2\upsilon}{s_\beta} + \frac{6}{\sqrt{2}}\mu_2 s_\beta\upsilon. \end{aligned} \quad (5)$$

It has been shown in [18] that the scalar potential (3) could acquire some minima that violate CP symmetry and/or electric charge, and could be deeper than the electroweak vacuum $\{\upsilon_\phi, \sqrt{2}\upsilon_\xi, \upsilon_\xi\}$. Therefore, this part of the parameter space is excluded. Here, we impose the constraints from (1) vacuum stability, (2) unitarity, (3) electroweak precision tests,

(4) the diphoton and undetermined Higgs branching ratios and the total decay width; in addition to (5) constraints from negative searches for light scalar resonances at LEP [19]. In this setup, the SM-like Higgs h (the CP-even scalar with $m_h = 125.18$ GeV) decays mainly into pairs of fermions $(cc, \mu\mu, \tau\tau, b\bar{b})$ and gauge bosons WW and ZZ, in addition to a pair of light scalars $\eta\eta$ when kinematically allowed. Since the Higgs couplings to SM fields are scaled by the coefficients

$$\kappa_{\rm F} = \frac{g_{hff}^{\rm GM}}{g_{hff}^{\rm SM}} = \frac{c_\alpha}{c_\beta},\qquad \kappa_V = \frac{g_{hVV}^{\rm GM}}{g_{hVV}^{\rm SM}} = c_\alpha c_\beta - \sqrt{\frac{8}{3}} s_\alpha s_\beta, \tag{6}$$

that are directly constrained by the Higgs partial strength modifiers [18], by using the recent measurements [20]. Besides the above-mentioned constraints, the negative searches for doubly charged Higgs bosons in the VBF channel $H_5^{++} \to W^+W^+$ and from Drell-Yan production of a neutral Higgs boson $pp \to H_5^0(\gamma\gamma)H_5^+$ give strong bounds on the parameter space [16]. It has been shown in [16] that the search for the doubly charged Higgs boson in the VBF channel leads to a constraint from CMS on $s_\beta^2 \times \mathcal{B}(H_5^{++} \to W^+W^+)$ [21]. Meanwhile, the relevant quantity for the constraints on $H_5^0 \to \gamma\gamma$ is the fiducial cross section multiplied by the branching ratio $\sigma_{\rm fid} = \left(\sigma_{H_5^0 H_5^+} \times \epsilon_+ + \sigma_{H_5^0 H_5^-} \times \epsilon_-\right)\mathcal{B}(H_5^0 \to \gamma\gamma)$, which is constrained by ATLAS at 8 TeV [22] and at 13 TeV [23]. Here, we used the decay rate formulas, as well as the cross section and efficiency values from [16], to include these constraints in our numerical analysis.

The extra CP-even eigenstate η couples to the SM fermions and gauge bosons via the SM couplings scaled by the modifiers

$$\zeta_V = \frac{g_{\eta VV}^{\rm GM}}{g_{hVV}^{\rm SM}} = s_\alpha c_\beta + \sqrt{\frac{8}{3}} c_\alpha s_\beta,\qquad \zeta_F = \frac{g_{\eta FF}^{\rm GM}}{g_{hFF}^{\rm SM}} = \frac{s_\alpha}{c_\beta}, \tag{7}$$

which can be constrained by all negative searches for light scalars. Here, we consider two types of searches for the light scalar η [17]: (1) direct production via $e^-e^+, pp \to \eta + X$, where the scalar could be identified via one of its SM-like decays $\eta \to \gamma\gamma, \mu\mu, \tau\tau, cc, bb$; and (2) indirect production via the Higgs decay $pp \to h \to \eta\eta \to X\overline{X}Y\overline{Y}$, where the light scalar is identified via its SM decays $X, Y = \gamma, \mu, \tau, c, b$. We consider the constraints from negative searches for $pp \to \eta \to \gamma\gamma$ at CMS at $8+13$ TeV Ref. [6], and at ATLAS at 13 TeV with integrated luminosities 80 fb^{-1} Ref. [7] and at 138 fb^{-1} Ref. [8].

Another search for a light SM-like scalar in the diphoton channel with masses in the range $70-110$ GeV has been performed by CMS at 8 TeV and 13 TeV Ref. [6], where upper bounds are established on the production cross section $\sigma(pp \to \eta) \times \mathcal{B}(\eta \to \gamma\gamma)$ scaled by its SM value, i.e., the factor $\zeta_F^2 \zeta_\gamma^2$, with ζ_X defined in Ref. [17]. Regarding the CMS bounds Ref. [6] on the production cross section of the CP-odd scalar $\sigma(pp \to H_3^0) \times \mathcal{B}(H_3^0 \to \gamma\gamma)$, the bounds are automatically fulfilled since $|\vartheta_F \vartheta_\gamma| < |\zeta_F \zeta_\gamma|$ for all the viable parameter space. Since the charged triplet $H_3^\pm$ is partially composed of the SM doublet as shown in (4), it couples to up- and down-type quarks similarly to the W gauge boson. These interactions lead to flavor-violating processes such as $b \to s$ transitions, which depend only on the charged triplet mass m_3 and the mixing angle β. The current experimental value of the $b \to s\gamma$ branching ratio, for a photon energy $E_\gamma > 1.6$ GeV, is $\mathcal{B}(\overline{B} \to X_s\gamma)_{\rm exp} = (3.55 \pm 0.24 \pm 0.09) \times 10^{-4}$, while the two SM predictions are $\mathcal{B}(\overline{B} \to X_s\gamma)_{\rm SM} = (3.15 \pm 0.23) \times 10^{-4}$ Ref. [24] and $\mathcal{B}(\overline{B} \to X_s\gamma)_{\rm SM} = (2.98 \pm 0.26) \times 10^{-4}$ Ref. [25]. In our numerical scan, we consider the most stringent bounds on the $m_3 - \upsilon_\chi$ plane shown in Fig. 1 of Ref. [14].

3. The Excess in the $\gamma\gamma$, $\tau\tau$ and $b\bar{b}$ Channels

Since the Higgs boson discovery with a mass around 125 GeV, see Ref. [1], the question about how electroweak symmetry breaking (EWSB) proceeded, remains open. It is not clear yet whether the EWSB proceed via a single Higgs as in the standard model (SM) or via many scalars as in many SM extensions. Therefore, the LHC physics program devotes many searches and analyses to the search for a di-Higgs signal and the search for additional scalar resonances, whether they are heavier or lighter than the 125 GeV Higgs, for example see Refs. [3, 4].

Although many searches for light scalar (η) have been performed at LEP and the LHC (8 and 13 TeV), where an excess around 95 GeV has been reported in these channels, see Refs. [9, 26–28]

$$\mu_{\gamma\gamma}^{\rm exp} = \frac{\sigma^{\rm exp}(gg \to \eta \to \gamma\gamma)}{\sigma^{\rm SM}(gg \to h \to \gamma\gamma)} = 0.27^{+0.10}_{-0.09},$$
$$\mu_{\tau\tau}^{\rm exp} = \frac{\sigma^{\rm exp}(gg \to \eta \to \tau\tau)}{\sigma^{\rm SM}(gg \to h \to \tau\tau)} = 1.2 \pm 0.5, \qquad (8)$$
$$\mu_{b\bar{b}}^{\rm exp} = \frac{\sigma^{\rm exp}\left(e^+e^- \to Z\eta \to Zb\bar{b}\right)}{\sigma^{\rm SM}\left(e^+e^- \to Zh \to Zb\bar{b}\right)} = 0.117 \pm 0.057,$$

with local significance values of $3.2\,\sigma$ [5], $2.2\ \sigma$ [9], and $2.3\ \sigma$ [10], respectively.

Here, we estimate the excess observed by both LEP and LHC around the 95.4 GeV mass value in the channels $\gamma\gamma$, $\tau\tau$, $b\bar{b}$, where the signal resonance is assumed to be a CP-even scalar for $93 < m_\eta < 97$ GeV, or a superposition of two resonances if $93 < m_\eta, m_{H_3^0} <$ < 97 GeV. Then, the 95 GeV signal strength modifiers can be written in the narrow width approximation (NWA) as

$$\mu_{\gamma\gamma}^{(95)} = \mu_{\gamma\gamma}^{(\eta)} + \mu_{\gamma\gamma}^{(H_3^0)} = \zeta_F^2\zeta_\gamma^2\left(\Gamma_\eta/\Gamma_\eta^{\rm SM}\right)^{-1} +$$
$$+\,\vartheta_F^2\vartheta_\gamma^2\left(\Gamma_{H_3^0}/\Gamma_{H_3^0}^{\rm SM}\right)^{-1},$$
$$\mu_{\tau\tau}^{(95)} = \mu_{\tau\tau}^{(\eta)} + \mu_{\tau\tau}^{(H_3^0)} = \zeta_F^4\left(\Gamma_\eta/\Gamma_\eta^{\rm SM}\right)^{-1} + \qquad (9)$$
$$+\,\vartheta_F^4\left(\Gamma_{H_3^0}/\Gamma_{H_3^0}^{\rm SM}\right)^{-1},$$
$$\mu_{b\bar{b}}^{(95)} = \mu_{b\bar{b}}^{(\eta)} = \zeta_V^2\zeta_F^2\left(\Gamma_\eta/\Gamma_\eta^{\rm SM}\right)^{-1},$$

with $\vartheta_i^2 = g_{H_3^0 ii}^{\rm GM}/g_{hii}^{\rm SM}$ for $i = V, F, \gamma$. Note that $\mu_{b\bar{b}}^{(95)}$ does not include a contribution from H_3^0 since the CP-odd scalar does not couple to the Z gauge boson. Clearly, the CP-even scalar H_5^0 cannot play a similar role as H_3^0 since it does not couple to quarks and therefore cannot be produced via gluon fusion at the LHC, nor does it decay into SM fermions if produced at LEP.

In order to estimate the relative contributions $\rho_{\gamma\gamma,\tau\tau} = \mu_{\gamma\gamma,\tau\tau}^{(H_3^0)}/\mu_{\gamma\gamma,\tau\tau}^{(95)}$ in (9), it is important to note that the H_3^0 total decay width is much smaller than its corresponding SM value, so the factor $\left(\Gamma_{H_3^0}/\Gamma_{H_3^0}^{\rm SM}\right)^{-1}$ may lead to a significant enhancement of $\mu_{\gamma\gamma,\tau\tau}^{(95)}$. In addition, the effective coupling modifier ϑ_γ is highly suppressed due to the absence of gauge and scalar loop contributions. This makes the ratio $\rho_{\tau\tau} =$ $= \mu_{\tau\tau}^{(H_3^0)}/\mu_{\tau\tau}^{(95)}$ comparable to unity, whereas $\rho_{\gamma\gamma} =$ $= \mu_{\gamma\gamma}^{(H_3^0)}/\mu_{\gamma\gamma}^{(95)}$ is strongly suppressed, as will be shown below.

If the 95 GeV excess is confirmed in the di-τ channel when new ATLAS results are reported with more data and/or similar results are released by CMS, this channel could be very useful for confirming whether the 95 GeV excess is a mixture of CP-even and CP-odd resonances. At the detector level, the τ lepton cannot be measured directly but is reconstructed from its decay products, especially the hadronic final states with $\mathcal{B}(\tau \to \text{hadrons}) = 64.79\%$ [20]. The two most important decay channels are $\tau^\pm \to pi^\pm\nu_\tau$ and $\tau^\pm \to \rho^\pm\nu_\tau \to \pi^\pm\pi^0\nu_\tau$, with branching ratios of 10.82% and 25.49%, respectively. The decay $\eta(H_3^0) \to$ $\to \tau^+\tau^- \to \pi^+\pi^-\nu_\tau\bar{\nu}_\tau$ is particularly useful for identifying the scalar CP properties via the so-called acoplanarity angle, defined as $\phi^* = \arccos(\mathbf{n}_+ \cdot \mathbf{n}_-)$, where $\mathbf{n}_\pm$ are unit vectors normal to the decay planes of the charged pions. The ϕ^* distribution is generally used to probe the Higgs CP phase $\Delta_{\rm CP}$ of the tau Yukawa interaction. The normalized acoplanarity angle distributions for pure CP-even and CP-odd cases are given by

$$R_{\rm even,odd}(\phi^*) = \frac{1}{N}\frac{dN}{d\phi^*} = \frac{1}{2\pi}[1 \mp Q\cos(\phi^*)], \qquad (10)$$

with $Q = \frac{\pi^2}{16}$ for $\tau \to \pi\nu_\tau$ and $Q = \frac{\pi^2}{16}\left(\frac{m_\tau^2 - 2m_\rho^2}{m_\tau^2 + 2m_\rho^2}\right)^2$ for $\tau \to \rho\nu_\tau$, respectively [29]. For a degenerate CP-even/CP-odd resonance, or twin peak resonance (TPR) scenario, the distribution is

$$R_{\rm TPR}(\phi^*) = (1 - \rho_{\tau\tau})R_{\rm even}(\phi^*) + \rho_{\tau\tau}R_{\rm odd}(\phi^*), \quad (11)$$

where $\rho_{\tau\tau} = \mu_{\tau\tau}^{(H_3^0)}/\mu_{\tau\tau}^{(95)}$ is defined above.

4. Numerical Results

By considering all the constraints discussed in the previous section, we perform a numerical scan where the masses lie in the ranges $93 < m_\eta, m_3 < 97$ GeV and $78 < m_5 < 2$ TeV, where $m_{3,5}$ are the triplet and quintuplet masses, respectively. Concerning the 95 GeV signal excess (8), we define the function $\chi^2_{(3)}$ as

$$\chi^2_{(3)} = \chi^2_{\gamma\gamma} + \chi^2_{b\bar{b}} + \chi^2_{\tau\tau}, \quad \chi_i^2 = \left(\frac{\mu_i - \mu_i^{\rm exp}}{\Delta\mu_i^{\rm exp}}\right)^2, \qquad (12)$$

which is useful to check whether the excess can be addressed simultaneously in the channels $\gamma\gamma, \tau\tau, b\bar{b}$. In

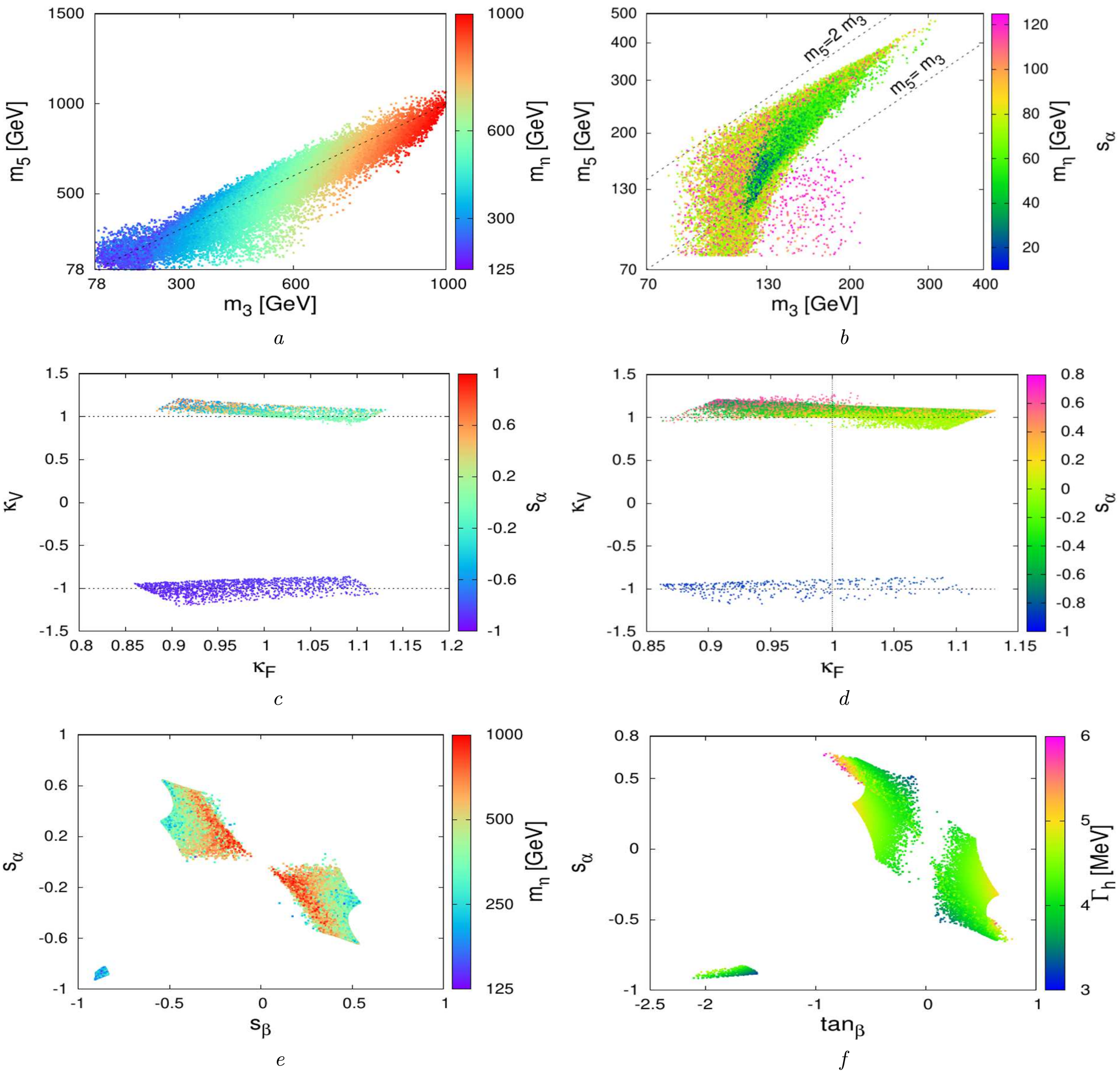
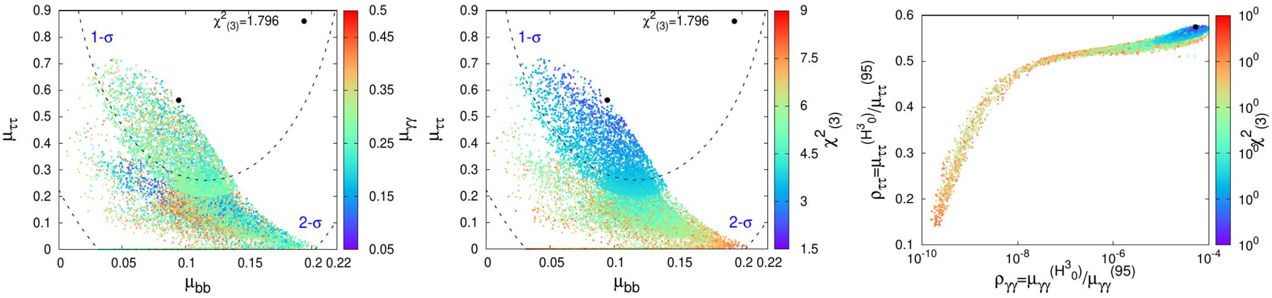


Fig. 1. The masses of the triplet, quintuplet, and singlet η for the case of a light (a) and a heavy (b) SM-like Higgs, considering all theoretical and experimental constraints. The values of the coupling modifiers $\kappa_{F,V}$ for the light (c) and heavy (d) SM-like Higgs cases. The mixing angles for the light (e) and heavy (f) SM-like Higgs cases

our analysis, we consider only the benchmark points (BPs) that address the three channels simultaneously within 2σ, i.e., $\chi^2_{(3)} < 8.02$.

In the following figures, we show the viable parameter space of the model for both the light and heavy SM-like Higgs scenarios, along with several constraints.

The parameter space is well constrained and split into three isolated regions, as shown in Fig. 1, *a*. For instance, the three islands correspond to $\{1.5 < t_\beta < 2.16,\ 0.82 < s_\alpha < 0.93\}$, $\{0.55 < t_\beta < 0.655,\ -0.38 < s_\alpha < -0.22\}$ and $\{0.55 < t_\beta < 0.64,\ -0.65 < s_\alpha < -0.53\}$, respectively. According to Fig. 1, *c*, the κ values for the three islands



a *b* *c*

Fig. 2. The signal strength values in Eq. (9) for 2σ viable BPs with $\chi^2_{(3)} < 8.02$ in the TPR case, considering all constraints (*a*, *b*). The black point represents the best-fit BP corresponding to $\chi^2_{(3)} = 1.796$. The palette shows the ratio $\mu_{\gamma\gamma}$ in the panel *a* and the $\chi^2_{(3)}$ function given in Eq. (12) in the panel *b*. The ratios $\rho_{\tau\tau} = \mu_{\tau\tau}^{(H_3^0)}/\mu_{\tau\tau}^{(95)}$ and $\mu_{\gamma\gamma}^{(H_3^0)}/\mu_{\gamma\gamma}^{(95)}$, where the palette represents the $\chi^2_{(3)}$ function, are shown in panel (*c*). The ratio $\rho_{\tau\tau}$ indicates the relative contribution of the CP-odd resonance to the excess

are $\{0.95 < \kappa_V < 1.09,\ 0.86 < \kappa_F < 1.13\}$, $\{1.018 < \kappa_V < 1.114,\ 1.06 < \kappa_F < 1.08\}$ and $\{1.162 < \kappa_V < 1.221,\ 0.903 < \kappa_F < 0.963\}$, respectively.

For the heavy SM-like Higgs scenario, the parameter space in the $\{m_3, m_5\}$ plane is different from the case with a light scalar η, while the $\{t_\beta, s_\alpha\}$ and $\{\kappa_F, \kappa_V\}$ planes are similar. The Higgs coupling modifier κ_V is tightly constrained and can have either sign, while κ_F can deviate from the SM value by up to 13% for both the heavy and light SM-like Higgs cases. These deviations of $\kappa_{F,V}$ from the SM are possible due to the strength of bounds from experimental constraints such as the Higgs diphoton decay, the total Higgs decay width, and the Higgs signal strength modifiers.

In Fig. 2, we show the signal strength modifier values (9) and the $\chi^2_{(3)}$ function for the considered 5k BPs in the TPR scenario.

According to Fig. 2, the H_3^0 contribution to the signal strengths is crucial for addressing the di-τ excess in the TPR case. Here, the $\chi^2_{(3)}$ values become significantly smaller than in the SPR case, where the 95 GeV excess is addressed only via the CP-even resonance. In the TPR case, about 27.5% of the BPs lie within 1σ, meaning that the excess in all three channels (8) is addressed simultaneously. In contrast, in the SPR case, the $\chi^2_{(3)}$ function values are larger than 1σ for all BPs considered in Fig. 2, with a minimal value of $\chi^{2,\min}_{(3)} = 3.487$.

5. Conclusion

In this work, we have revisited the scalar sector of the Georgi–Machacek model in light of the reported excesses around 95 GeV in the diphoton, di-tau, and $b\bar{b}$ channels. After imposing theoretical constraints (vacuum stability, unitarity, and custodial symmetry) and experimental limits (electroweak precision tests, Higgs signal strengths, direct searches for new scalars and doubly charged Higgs bosons), we performed a numerical scan over the model parameters, considering both the light and heavy SM-like Higgs scenarios.

We find that the GM model can successfully accommodate the 95 GeV excesses, especially in the "twin peak resonance" (TPR) scenario where the CP-even scalar η and the CP-odd scalar H_3^0 are nearly degenerate in mass. In this case, the contributions from both states allow a simultaneous fit to the $\gamma\gamma$, $\tau\tau$, and $b\bar{b}$ signal strengths within 2σ, with a significant fraction of benchmark points even lying within 1σ. By contrast, the "single peak" scenario (with only η) yields a poorer fit, particularly for the di-tau excess.

A key signature of the TPR scenario is the modified acoplanarity distribution in the $\tau^+\tau^-$ decay, which reflects the superposition of CP-even and CP-odd contributions. Future high-statistics measurements of this distribution at the LHC or at a future e^+e^- collider could therefore distinguish between a pure CP-even resonance and a mixed CP-even/CP-odd system.

Our analysis also shows that the coupling modifiers κ_V and κ_F of the 125 GeV Higgs boson can deviate from their SM values by up to $\sim$10% in viable regions of the parameter space, which may be probed with improved precision in future Higgs coupling measurements. Moreover, the allowed mass ranges for the triplet (H_3) and quintuplet (H_5) states are tightly constrained, with the quintuplet mass typically lying above a few hundred GeV.

In summary, the GM model offers a theoretically consistent and phenomenologically viable framework that can explain the 95 GeV excesses while predicting distinctive collider signatures. Further data from the LHC and future colliders will be crucial to test this scenario and to unravel the possible extended nature of the Higgs sector.

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СКАЛЯРНИЙ СЕКТОР
У МОДЕЛІ ДЖОРДЖІ–МАЧАЧЕКА

Модель Джорджі–Мачачека (GM) розширює сектор Хіггса Стандартної моделі (СМ), додаючи один комплексний та один дійсний скалярний триплет, зберігаючи при цьому симетрію $SU(2)_V$. Цей підхід передбачає багатий скалярний спектр, який включає квінтет, триплет та два CP-парні синглети. У цій роботі ми розглядаємо структуру моделі, її спектр мас, а також теоретичні та експериментальні обмеження, пов'язані з пертурбативністю, вакуумною стабільністю, електрослабкими прецизійними тестами, вимірюваннями бозона Хіггса та прямим пошуком на колайдерах в обох випадках, коли СМ-подібний бозон Хіггса має бути легким або важким CP-парним власним станом. Ми також обговорюємо здатність моделі пояснити надлишок 95 ГеВ, який присутній в $\gamma\gamma$, $\tau\tau$ та $b\bar{b}$ каналах у припущенні двох вироджених (CP-парного та CP-непарного) резонансів. Ми показуємо, що природу кандидата на скалярний резонанс з енергією 95 ГеВ можна дослідити за допомогою аналізу властивостей його розпаду на пару τ-лептонів.

Ключові слова: Модель GM, бозон Хіггса, надлишок 95 ГеВ, модифікатори констант зв'язку.