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(Received 15 July 2016; published 28 October 2016)

We discuss a test of the Standard Model fermion mass origin in models of dynamical electroweak symmetry breaking. The couplings of composite pseudoscalar resonances to top quarks allow us to distinguish high-scale extended-technicolor-type fermion mass generation from fermion partial compositeness and low-scale mass generation via an induced vacuum expectation value of a doublet coupled to the composite sector. These different possible origins of fermion masses are thus accessible via weak-scale physics searched for at the LHC.

I. INTRODUCTION

In this paper, we consider models of dynamical electroweak symmetry breaking (EWSB). We show that the observation of isoscalar CP-odd composite resonances η_X not only sheds light on the underlying composite dynamics but also on the mechanism of mass generation for the Standard Model (SM) fermions. In particular, the would-be Goldstone boson (GB) related to the $U(1)_A$ anomaly that we refer to as the η_1 state is interesting in this respect—in QCD the analogous η_1 and η_8 states mix into the physical η, η' states.

If the top-quark mass arises from a direct coupling of the top quarks to the condensate in the strong sector, mediated by new high-scale states as in extended-technicolor (ETC) models [1–3], this implies Higgs-sized couplings of the new η_1 state to top quarks. This is in contrast to the case where the SM-fermion masses arise via the so-called fermion partial compositeness (PC) mechanism [4], or from a dynamically induced vacuum expectation value (vev) of a weak doublet, i.e., bosonic technicolor (bTC). We note that such a doublet may be either elementary [5,6] or composite [7].

In PC the η_1 couplings to top quarks are suppressed if the compositeness scale is higher than the electroweak scale and by the smallness of the underlying coupling between η_1 and the composite top partner, as discussed in Sec. III C. In bTC the coupling is small because only the components of the doublet acquiring a vev are coupled directly to the SM fermions at interaction-eigenstate level, and the η_1 couples to the SM fermions only via mass mixing with the components of the doublet.

The partial widths into $\bar{t}t$ versus $\gamma\gamma$ and gg of the η_1 resonances thus provide an experimentally accessible diagnostic of the origin of SM-fermion masses. A complementary diagnostic of the fermion-mass origin is the momentum dependence of the Yukawa couplings above the mass-generation scale [8].

II. EFFECTIVE LAGRANGIAN DESCRIPTION OF THE DYNAMICAL EWSB SECTOR

We begin the discussion with a new sector able to break the electroweak symmetry dynamically and generate the observed W and Z masses [9,10]. Minimally, such a sector contains a single weak doublet of left-handed technifermions $Q_L = (U_L, D_L)$ and corresponding right-handed weak-singlet fermions U_R, D_R transforming in some representation under a new confining gauge group G. If the number of weak-doublet fermions is even, then the absence of gauge anomalies fixes the charges in the minimal case to $Q(U) = 1/2$, $Q(D) = -1/2$. If the number of new weak doublet fermions is odd, then gauge anomalies can be satisfied by adding more matter and choosing the hypercharges appropriately, e.g., [11].

Well below the condensation scale, we can describe the new strong sector using an effective Lagrangian invariant under a global symmetry group containing $SU(2)_L \times SU(2)_R$. Restricting to this subgroup, we write the Lagrangian in terms of a nonlinear representation of the three GBs $\pi^{\pm,0}$ and η_1 via

$$\Sigma = \exp(i\pi/F_\pi), \qquad \pi = \begin{pmatrix} \eta_1 + \pi^0 & \sqrt{2}\pi^+ \\ \sqrt{2}\pi^- & \eta_1 - \pi^0 \end{pmatrix}. \quad (1)$$

The transformation of Σ under the global symmetry group is $\Sigma \to g_L \Sigma g_R^\dagger$ where $g_{L,R} \in SU(2)_{L,R}$. The leading-order term of the effective Lagrangian is

$$\mathcal{L}_K = \frac{F_\pi^2}{4} \text{Tr}[D_\mu^\dagger \Sigma D^\mu \Sigma], \quad (2)$$

where D_μ is the electroweak covariant derivative, while the anomaly-induced mass of the would-be GB η_1 can be encoded at the effective-Lagrangian level via the operator

$$\mathcal{L}_m = \frac{m_{\eta_1}^2}{32} F_\pi^2 \text{Tr}[\ln \Sigma - \ln \Sigma^\dagger]^2. \quad (3)$$

The η_1 state is particularly interesting, because its presence in the spectrum is model independent and its mass is sensitive to the compositeness scale and fermion content of the composite sector via $m_{\eta_1}^2 \simeq \frac{N_f}{d(R)} \Lambda^2$, where $d(R)$ is the dimension of the constituent fermion representation under the strongly interacting gauge group, N_f is the number of flavors, and Λ is a scale related to the GB decay constant and the dynamics-dependent anomaly term [12,13].[1] At the same time its couplings are sensitive to the fermion mass mechanism.

Nevertheless, in models with nonminimal global symmetries, other isosinglet pseudo-GBs may be lighter, and the couplings of these states still probe the origin of fermion masses. We collectively denote these η_X.

III. ORIGIN OF SM-FERMION MASSES

To generate SM-fermion masses, the condensate must be communicated to the SM fermions. One way is to extend the model with four-fermion interactions coupling two SM fermions, q, and two technifermions, Q, via dimension-six operators of the schematic form $\mathcal{O}_4 \sim \frac{1}{\Lambda^2} q_L \bar{q}_R Q_L \bar{Q}_R$. These four-fermion operators can, e.g., arise from exchanges of heavy spin-one fields [1,2] (after Fierz rearrangements) or heavy scalar doublets [3]. In the following, we will refer to this as the ETC scenario, independent of the origin of four-fermion operators.

Another possibility is fermion partial compositeness, where dimension-six operators of the form $\mathcal{O}_4 \sim \frac{1}{\Lambda^2} qQQQ$ induce mixing of the SM fermions and composite baryons $B \sim QQQ$.

In these scenarios, such as technicolor or composite Higgs, a CP-even composite resonance could play the role of the observed Higgs boson.

Another possibility is to extend the model with a light scalar doublet Φ, analogous to the SM-Higgs doublet, coupled to the SM fermions and the new strongly interacting fermions via Yukawa interactions. The doublet Φ may then obtain an induced vev from condensation and generate the SM-fermion masses [5,6]. The neutral CP-even component of Φ may be interpreted as the 125-GeV Higgs boson. The scalar doublet Φ may be fundamental and, e.g., embedded in a supersymmetric theory at a higher scale [5,6], or composite and related to some additional composite dynamics above the weak scale [7]. We will refer to this scenario as bTC, as in [5,6].

A. Four-fermion operators (ETC)

Upon integrating out the heavy ETC fields, the $SU(2)_L \times U(1)_Y$-invariant effective operators responsible

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[1]In the Veneziano limit for the $SU(N)$ gauge group, $m_{\eta_1}^2 \simeq 480 \frac{N_f}{N^2} F_\pi^2$ [14].

for generating the SM-fermion masses can be written as [14–16]

$$\mathcal{L}_{\mathrm{ETC}} = -Y_1 F_\pi f_1(h/F_\pi)(\bar{q}_L \Sigma q_R + \mathrm{H.c.})$$
$$- Y_2 F_\pi f_2(h/F_\pi)(\bar{q}_L \Sigma \tau^3 q_R + \mathrm{H.c.}) + \cdots, \quad (4)$$

where we restrict to the third-generation SM-quark doublet $q_{L,R} = (t,b)_{L,R}$ and have absorbed the ETC scale in the Yukawa couplings, $Y_{1,2}$, and the dots represent additional interactions generated at low energies by the ETC interactions. The functions $f_{1,2}$ describe the excitations of the condensate,

$$f_{1,2} = 1 + \kappa_{1,2} \frac{h}{F_\pi} + \cdots. \quad (5)$$

The lightest excitation, h, of the condensate may be interpreted as the 125-GeV Higgs provided the unknown coefficients $\kappa_{1,2}$ reproduce the observed Higgs couplings within uncertainties. The $Y_{1,2}$ couplings are fixed by requiring the correct t and b masses: $m_t = (Y_1 + Y_2)F_\pi$, and $m_b = (Y_1 - Y_2)F_\pi$. If the top mass is generated this way from ETC, then this also fixes the coupling between the quarks and η_X; see, e.g., [14–16].

An example of a composite sector featuring only a single doublet of technifermions is the $SU(3)_S$ next-to-minimal walking technicolor (NMWT) model [11,17,18] with $\eta_{\mathrm{NMWT}} \sim \frac{1}{\sqrt{2}}(\bar{U}U + \bar{D}D)$. Then, as emphasized in [16], we have

$$\mathcal{L}_{\eta\bar{q}q} = -i\frac{m_t}{F_\pi}\eta_{\mathrm{NMWT}}\bar{t}\gamma^5 t - i\frac{m_b}{F_\pi}\eta_{\mathrm{NMWT}}\bar{b}\gamma^5 b, \quad (6)$$

i.e., the top coupling is identical to that of the SM Higgs.

An example of a model with a larger global symmetry group and including colored constituents, is the so-called one-family model [19] for which ETC [2] and fermion partial compositeness [4] extensions have both been explicitly constructed. In the one-family model, the diagonal η_1-type state (in analogy with QCD) is a mixture of colored and uncolored states and comes with an isosinglet partner, η_{63}, the analogue of the η_8 in the pseudoscalar nonet of QCD.

The anomaly coefficients $\mathcal{A}_{V_1 V_2}$, arising from the strongly interacting constituents, that determine the dijet and diphoton decays of η_X are defined as [20]

$$\mathcal{A}_{V_1 V_2} = \mathrm{Tr}[T^a(T_1 T_2 + T_2 T_1)_L + L \leftrightarrow R], \quad (7)$$

where T^a is the generator of the axial-vector current associated with the state in question and L,R denote the left- and right-handed underlying fermion components of the state. The anomaly coefficients and the reduced Yukawa couplings $c_X^{\mathrm{ETC}} \equiv y_X^{\mathrm{ETC}} v_w/m_t$, induced by ETC, for the one-family model are given by

State	A_{gg}	$A_{\gamma\gamma}$	c_X^{ETC}
$\eta_1 \sim \frac{1}{4}(\bar{Q}i\gamma_5 Q + \bar{L}i\gamma_5 L)$	1	$\frac{8}{3}$	1
$\eta_{63} \sim \frac{1}{4\sqrt{3}}(\bar{Q}i\gamma_5 Q - 3\bar{L}i\gamma_5 L)$	$\frac{1}{\sqrt{3}}$	$\frac{-4}{3\sqrt{3}}$	$\frac{1}{\sqrt{3}}$

In a generic model, due to the η_X wave functions, the top-Yukawa couplings involve $\mathcal{O}(1)$ normalization factors as discussed below.

Returning to the low-energy description of ETC-induced mass terms in Eq. (4), the Yukawa couplings $Y_{1,2}$ can, e.g., be generated via integrating out a heavy scalar doublet [3] with interactions

$$\mathcal{L}_{\mathrm{Yuk}} = -y_t \bar{q}_{\mathrm{L}} \tilde{\Phi} t_{\mathrm{R}} - y_b \bar{q}_{\mathrm{L}} \Phi b_{\mathrm{R}} \\ - y_U \bar{Q}_{\mathrm{L}} \tilde{\Phi} U_{\mathrm{R}} - y_D \bar{Q}_{\mathrm{L}} \Phi D_{\mathrm{R}} + \mathrm{H.c.} \quad (8)$$

If we define the average coupling, $Y \equiv \frac{1}{2}(y_U + y_D)$, and the relative difference, $\delta \equiv \frac{y_U - y_D}{y_U + y_D}$, we can write $Y_1 \sim (y_t - y_b)Y F_\pi^2/m_\Phi^2$ and $Y_2 \sim (y_t + y_b)Y F_\pi^2/m_\Phi^2$. The T parameter, measuring the amount of isospin breaking in the technicolor (TC) sector, is proportional to $\delta^2 Y^4$ [21]. Therefore, in the limit $y_U = y_D$, fermion masses can be generated without contributions to the T parameter beyond those of the SM.

Alternatively the interactions in Eq. (4) arise after Fierz-transforming vector-current four-fermion operators [1,2]. The discussion of the η_X couplings is not affected by the difference in origin.

B. Induced vev (bTC)

Another possibility is that the scalar doublet in Eq. (8) cannot be integrated out and obtains a vev, v. To obtain the correct electroweak gauge boson masses, we require

$$v_{\mathrm{w}}^2 = N_D F_\pi^2 + v^2, \quad (9)$$

with $v_{\mathrm{w}} = 246$ GeV, and N_D is the number of families of condensing fermions. The current measurements of the Higgs couplings constrain $\sqrt{N_D} F_\pi \lesssim 0.4 v_{\mathrm{w}}$ if the scalar excitation of Φ is interpreted as the Higgs. On the other hand, it is reasonable to assume $F_\pi \gtrsim 0.25 v_{\mathrm{w}}$ not to have a very light resonance spectrum, in disagreement with direct search constraints [22,23].[2]

The Yukawa terms in Eq. (8) induce the following operators relevant to fermion masses:

$$\mathcal{L}_{\mathrm{mix}} = 4\pi c_1 F_\pi^3(y_U \mathrm{Tr}[U_\Phi \Sigma_U] + y_D \mathrm{Tr}[U_\Phi \Sigma_D]) + \mathrm{H.c} \quad (10)$$

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[2]The composite and elementary doublets can, in fact, mix kinetically, and this can modify Eq. (9), see e.g. [24]. This effect can help lift the compositeness scale in bTC scenarios and therefore help evade the direct search constraints. Alternatively, the compositeness scale can be raised by vacuum misalignment mechanism as in composite Higgs scenarios.

where $c_1 \sim \mathcal{O}(1)$ by naive dimensional analysis [25,26], $U_\Phi = (\Phi \tilde{\Phi})$ is the $\mathrm{SU}(2)_{\mathrm{L}} \times \mathrm{SU}(2)_{\mathrm{R}}$ bidoublet matrix, and we have defined the projections $\Sigma_{U,D} = \Sigma P_{U,D}$ with $P_{U,D} \equiv \frac{1}{2}(1 \pm \tau^3)$.

The above terms generate a mass mixing between the pseudoscalar triplet π_f contained in Φ, the composite triplet of "pions" π_c in Eq. (1) and η_X. The (mostly) η_X mass eigenstate couples to the top quark only via this mass mixing with π_f^0 which is absent in the $y_U = y_D$ limit of unbroken isospin in the composite sector. To lowest order in the isospin-breaking parameter δ, we find

$$c_X \approx 8\sqrt{2}\pi c_1 Y\delta \frac{F_\pi^2}{m_{\eta_X}^2} \frac{v_{\mathrm{w}}}{v} + \mathcal{O}(\delta^2), \quad (11)$$

where $c_X \equiv y_X v_{\mathrm{w}}/m_t$ is the η_X Yukawa coupling in units of the SM Higgs Yukawa coupling. Similarly to the case with ETC, in extended models the η_X wave functions involve $\mathcal{O}(1)$ normalization factors that can further suppress the top coupling. The amount of isospin breaking in the TC sector is highly constrained by the T parameter, which, in the limit where the elementary scalar is much lighter than the compositeness scale, $m_\Phi \ll \Lambda$, is proportional to $\delta^2 Y^2$ [21].

Fixing $m_\Phi = 125$ GeV and $c_1 = 1$, we find that requiring the T parameter within 1σ confidence level from the experimental value constrains $\delta \cdot Y \lesssim 0.2$ for $F_\pi > 0.25 v_{\mathrm{w}}$.

We also note that the physical triplet of CP-odd states couples to top quarks only through its small elementary component. Therefore this top coupling is also very suppressed, and these additional pseudoscalar states do not spoil our ability to distinguish this scenario from ETC-induced fermion masses. However, their mass and properties are further constrained by data and vacuum stability [27].

C. Top partners (PC)

A third possibility to generate masses for the SM fermions is via mixing with composite baryons [4]. Then the low-energy effective theory with the composite baryon $B \sim QQQ$ contains the following operators:

$$\mathcal{L}_{\mathrm{eff}} = \epsilon(qB + \bar{q}\,\bar{B}) - m_B B\bar{B} + \mathrm{H.c.} \quad (12)$$

and diagonalization gives rise to the quark mass q.

The composite interaction $C_X \frac{m_B}{F_\pi}\eta_X B\bar{B}$ generates the coupling between the η_X state and the SM top quark after diagonalization. The result is given by

$$c_X \approx \frac{C_X v_{\mathrm{w}}}{F_\pi}, \quad (13)$$

where F_π is the compositeness scale and C_X includes the form factor between the η_X and the technibaryon and the normalization factor of the state. In composite Higgs

models, where fermion partial compositeness has been extensively employed, the electroweak scale arises due to vacuum misalignment such that $v_w = F_\pi \sin \theta$ and $\theta \ll 1$ to avoid electroweak precision test constraints. The above top coupling is therefore significantly smaller than in the ETC case.

Moreover, in QCD the coupling between proton and the η_1 related to the axial anomaly is small [28,29], implying that $C_1 \ll 1$. Therefore, for top-quark partial compositeness in composite Higgs models, the coupling to axial η_1 is expected to be suppressed compared to the ETC case by both the higher compositeness scale and the small $\eta_1 B \bar{B}$ coupling.

IV. EXPERIMENTAL TEST OF THE FERMION-MASS MECHANISM

We have seen that the top coupling of the isosinglet CP-odd resonance, η_X, in the new gauge sector is essentially binary with respect to the fermion-mass mechanisms considered: In the case of high-scale ETC-type mass generation, the top coupling is SM-Higgs sized. Instead, when a doublet acquiring a vev or mixing with top partners is responsible for fermion masses, the top coupling of η_X is expected to be much smaller—ultimately limited by the contribution to the T parameter or the compositeness scale. Thus, this state presents a sensitive probe of the fermion-mass mechanism in models of dynamical EWSB.

Following the notation in [30], the partial widths for subsequent decays into top quarks, photons, and gluons are then given by

$$\Gamma_{\eta_X \to \bar{t}t} = \frac{3}{8\pi} \frac{m_t^2}{v_w^2} c_X^2 m_{\eta_X} (1 - m_t^2/m_{\eta_X}^2)^{1/2}$$

$$\Gamma_{\eta_X \to \gamma\gamma} \simeq \frac{\alpha^2 m_{\eta_X}^3}{256\pi^3 v_w^2} \left| \frac{4}{3} c_X A_{1/2}^A(\tau_t) + \frac{N_{TC} v_w}{F_\pi} \mathcal{A}_{\gamma\gamma}^\eta \right|^2$$

$$\Gamma_{\eta_X \to gg} \simeq \frac{\alpha_s^2 m_{\eta_X}^3}{128\pi^3 v_w^2} \left| c_X A_{1/2}^A(\tau_t) + \frac{2N_{TC} v_w}{F_\pi} \mathcal{A}_{gg}^\eta \right|^2 \quad (14)$$

where we neglect the small contribution of light fermions and we fix the strong-interaction scale by Eq. (9). The loop function $A_{1/2}^A(\tau_f)$ is given in [30] [Eq. (2.26)].

Relative to a SM Higgs at a given mass $m_{h_{SM}}$, the production cross section of η_X is [20]

$$\frac{\sigma_{gg \to \eta_X}}{\sigma_{gg \to h_{SM}}} = \frac{\Gamma(\eta_X \to gg)}{\Gamma(h_{SM} \to gg)} \equiv R_{gg}. \quad (15)$$

If some of the η_X constituents carry color, and they dominate the gluon-fusion production, we obtain

$$R_{gg} \simeq 4N_{TC}^2 (A_{gg}^\eta)^2 \frac{v_w^2}{F_\pi^2} |A_{1/2}^\mathcal{H}(\tau_t)|^{-2}, \quad (16)$$

where $A_{1/2}^\mathcal{H}(\tau_f)$ is given in [30] [Eq. (2.25)].

The η_X-like state may also be produced in association with other states, e.g., via the decays of spin-one isosinglet or isotriplet resonances like the analogues of the QCD ω and ρ mesons. This requires these states to be heavy enough that the decays $\omega \to \eta_X \gamma$ or $\rho \to \eta_X \gamma$ proceed on shell; see, e.g., [31]. In this study we focus on the gluon-fusion production.

The diagnostic of ETC interactions, given discovery of an η_X resonance, is therefore a relatively broad resonance dominated by the top decay mode.

The $\bar{t}t$ cross sections of the η_X states from the one-family model and the NMWT model are displayed in the top panel of Fig. 1 assuming the top mass arises from ETC-type interactions. Since in this case $\Gamma_{\eta_X \to tt} \simeq \Gamma_{\eta_X}$, the curves also represent the production cross sections of the η_X states to a good approximation. We used the N³LO prediction for the scalar production through gluon fusion [32] and rescaled by R_{gg}. The figure also shows the 95% C.L. limit on the $\bar{t}t$ production cross section of a heavy resonance provided by the ATLAS analysis of highly boosted tops decaying semileptonically [33] with luminosity $L_0 = 3.2$ fb^{-1} (black solid curve). The limit indicates that the LHC is already able to rule out the ETC scenario in the $t\bar{t}$ channel in some composite scenarios, but our analysis does not include interference effects which are known to be important even for widths $\Gamma_{\eta_X}/m_{\eta_X} \gtrsim 2\%$ [34–38] as here.

The middle and lower panels of Fig. 1 show production cross section times branching ratio into dijets and diphotons, respectively, for the ETC (red), the bTC (blue), and the PC (magenta) cases.

For the bTC case we take $v^2 = (3/4)v_w^2$ to satisfy the observed Higgs decay rates. For a four-doublet model such as the one-family model, this implies a very low compositeness scale $F_\pi = v_w/4$ due to the constraint in Eq. (9). For PC, the cross sections are suppressed by the compositeness scale, and we use $F_\pi = 1$ TeV as a benchmark value to avoid constraints from electroweak precision tests in the composite-Higgs case. We note that the η_1 in this case might be too heavy for LHC reach, and a better probe for our diagnostic may be provided by a pseudo-GB, or a more energetic future collider may be necessary.

In the middle panel, we also show the ATLAS 95% C.L. exclusion limit on dijet resonances for low [39] and high [40] resonance mass searches. Since the widths of our η_X resonances satisfy $\Gamma_{\eta_X}/m_{\eta_X} \lesssim 5 \times 10^{-2}$, we used the exclusion limit relevant for a width equal to the detector mass resolution (referred to as the "Res". analysis in [40]). We also assumed (conservatively) an acceptance of 30%. Similarly, in the lower panel, we show the diphoton cross sections limits from ATLAS in [41] which assumes a narrow resonance with $\Gamma/m = 2 \times 10^{-2}$. The integrated luminosities used in the dijet (high mass), dijet (low mass), and diphoton searches are, respectively, $L_0 = 3.4$, 3.6, and 3.2 fb^{-1}. Our projections for the cross-section limits from a very high luminosity (VHL) search with $L = 3$ ab^{-1} are
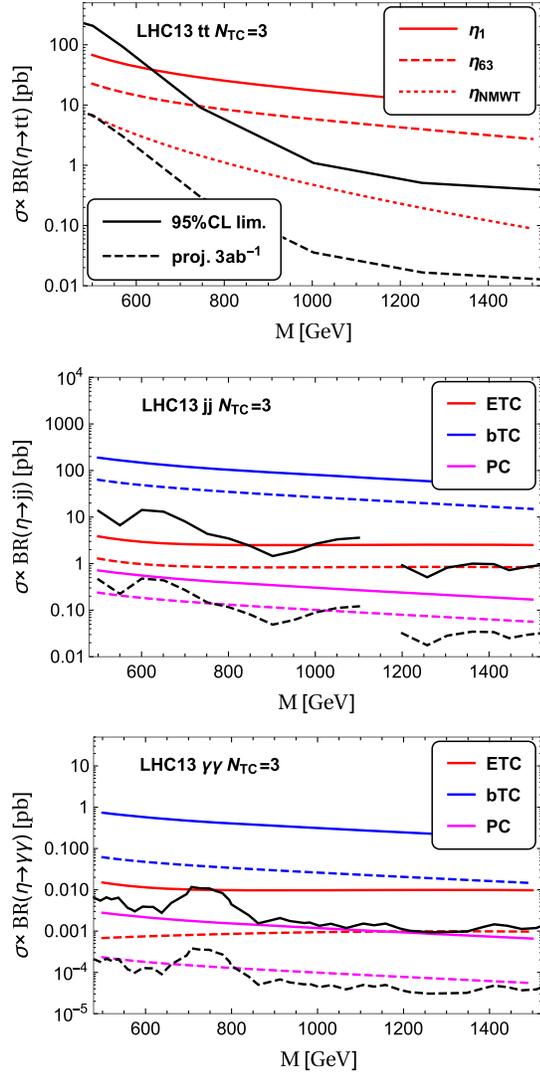
FIG. 1.   Production cross section times branching ratios of the η_X states when ETC (red lines, $c_X^{\rm ETC}$), bTC (blue lines, $c_X = 0$) and PC (magenta lines, $c_X = 0$) is responsible for the top quark mass. Also shown are the LHC limits (solid black lines) and reach projections (black dashed lines): top, $t\bar{t}$ channel; middle, dijet channel; bottom, diphoton channel. We use $N_{\rm TC} = 3$ as a benchmark value and fix $v^2 = \frac{3}{4} v_{\rm w}^2$ for the one-family model ($F_\pi = v_{\rm w}/4$), while $v = 0$ for the NMWT model (dotted line) ($F_\pi = v_{\rm w}$). For PC, we use $F_\pi = 1$ TeV as a benchmark value. We have omitted the NMWT case in the dijet and diphoton channels for clarity, since the cross section is tiny.

obtained by a simple rescaling of the current limit by $\sqrt{L_0/L}$ and depicted as the dashed black lines. It can be seen that in the one-family model setup both ETC and bTC

mechanisms are excluded and only PC may generate fermion masses, which will be probed in the near future for scales F_π up to 1 TeV. Other UV realizations should be studied.

Finally, for a minimal model without colored constituents, such as the $SU(3)_S$ NMWT example above, the gluon-fusion cross section is induced only via the top quark. In the ETC case, the production cross section shown with a dotted line in the top panel of Fig. 1 is then of the order of the SM Higgs with similar mass. This is below current sensitivity in the $\bar{t}t$ channel but the VHL LHC could be able to probe this state in the future. The associated production channels from heavy resonance decays mentioned above will also be relevant to probe this scenario.

V. CONCLUSIONS

There are multiple possible realizations of the SM fermion masses within the framework of dynamical EWSB. In particular, the scale of fermion-mass generation may be much higher than the weak scale as in ETC and fermion PC models. It is therefore very important to establish experimental probes able to test the fermion-mass generation in models with dynamical EWSB.

Here we have shown that a ubiquitous feature of weak-scale dynamical EWSB within reach of the LHC experiment, i.e., isoscalar CP-odd resonances, also provides experimental access to the origin of SM-fermion masses, notably the top-quark mass. We have shown that the value of the SM-top couplings for these resonances is very sensitive to the underlying fermion-mass-generation mechanism: If high-scale interactions generate four-fermion operators that upon EWSB provide the top-quark mass, these resonances acquire an $\mathcal{O}(1)$ Yukawa coupling to the top quark like the SM Higgs [15,16]. In models where, instead, the top mass arises from an induced vev of a scalar doublet, the top coupling is either $\ll 1$ or correlated with a sizable contribution to the T parameter. In composite Higgs models with fermion partial compositeness it is suppressed by the high compositeness scale. The partial width of these resonances into $\bar{t}t$ relative to partial widths into $\gamma\gamma$ and gg therefore offers a diagnostic of the underlying fermion-mass mechanism.

ACKNOWLEDGMENTS

[1] E. Eichten and K. D. Lane, Phys. Lett. **90B**, 125 (1980).
[2] S. Dimopoulos and L. Susskind, Nucl. Phys. **B155**, 237 (1979).
[3] G. 't Hooft, in *Recent Developments in Gauge Theories*, edited by G. 't Hooft *et al.*, NATO ASI, Ser. B, Vol. 59 (Springer, New York, 1980), p. 135.
[4] D. B. Kaplan, Nucl. Phys. **B365**, 259 (1991).
[5] S. Samuel, Nucl. Phys. **B347**, 625 (1990).
[6] M. Dine, A. Kagan, and S. Samuel, Phys. Lett. B **243**, 250 (1990).
[7] R. S. Chivukula, A. G. Cohen, and K. D. Lane, Nucl. Phys. **B343**, 554 (1990).
[8] N. D. Christensen and R. Shrock, Phys. Rev. Lett. **94**, 241801 (2005).
[9] S. Weinberg, Phys. Rev. D **13**, 974 (1976).
[10] L. Susskind, Phys. Rev. D **20**, 2619 (1979).
[11] D. D. Dietrich, F. Sannino, and K. Tuominen, Phys. Rev. D **72**, 055001 (2005).
[12] E. Witten, Nucl. Phys. **B156**, 269 (1979).
[13] G. Veneziano, Nucl. Phys. **B159**, 213 (1979).
[14] P. Di Vecchia and G. Veneziano, Phys. Lett. **95B**, 247 (1980).
[15] J. Jia, S. Matsuzaki, and K. Yamawaki, Phys. Rev. D **87**, 016006 (2013).
[16] E. Molinaro, F. Sannino, and N. Vignaroli, Nucl. Phys. **B911**, 106 (2016).
[17] F. Sannino and K. Tuominen, Phys. Rev. D **71**, 051901 (2005).
[18] A. Belyaev, R. Foadi, M. T. Frandsen, M. Järvinen, F. Sannino, and A. Pukhov, Phys. Rev. D **79**, 035006 (2009).
[19] E. Farhi and L. Susskind, Phys. Rep. **74**, 277 (1981).
[20] R. S. Chivukula, P. Ittisamai, E. H. Simmons, and J. Ren, Phys. Rev. D **84**, 115025 (2011); **85**, 119903(E) (2012).
[21] C. D. Carone and E. H. Simmons, Nucl. Phys. **B397**, 591 (1993).
[22] Z. Fodor, K. Holland, J. Kuti, S. Mondal, D. Nogradi, and C. H. Wong, *Proc. Sci.*, LATTICE2015 (2016) 219 [arXiv:1605.08750].
[23] T. Appelquist *et al.*, Phys. Rev. D **93**, 114514 (2016).
[24] T. Alanne, S. Di Chiara, and K. Tuominen, J. High Energy Phys. 01 (2014) 041.
[25] A. Manohar and H. Georgi, Nucl. Phys. **B234**, 189 (1984).
[26] H. Georgi and L. Randall, Nucl. Phys. **B276**, 241 (1986).
[27] C. D. Carone, Phys. Rev. D **86**, 055011 (2012).
[28] G. Veneziano, Mod. Phys. Lett. A **04**, 1605 (1989).
[29] P. Moskal *et al.*, Phys. Rev. Lett. **80**, 3202 (1998).
[30] A. Djouadi, Phys. Rep. **459**, 1 (2008).
[31] C. D. Carone and M. Golden, Phys. Rev. D **49**, 6211 (1994).
[32] C. Anastasiou, C. Duhr, F. Dulat, E. Furlan, T. Gehrmann, F. Herzog, A. Lazopoulos, and B. Mistlberger, J. High Energy Phys. 09 (2016) 037.
[33] ATLAS Collaboration, Report No. ATLAS-CONF-2016-014, 2016.
[34] B. Hespel, F. Maltoni, and E. Vryonidou, J. High Energy Phys. 10 (2016) 016.
[35] W. Bernreuther, P. Galler, C. Mellein, Z. G. Si, and P. Uwer, Phys. Rev. D **93**, 034032 (2016).
[36] D. Dicus, A. Stange, and S. Willenbrock, Phys. Lett. B **333**, 126 (1994).
[37] R. Frederix and F. Maltoni, J. High Energy Phys. 01 (2009) 047.
[38] S. Jung, J. Song, and Y. W. Yoon, Phys. Rev. D **92**, 055009 (2015).
[39] ATLAS Collaboration, Report No. ATLAS-CONF-2016-030, 2016.
[40] G. Aad *et al.* (ATLAS Collaboration), Phys. Lett. B **754**, 302 (2016).
[41] M. Aaboud *et al.* (ATLAS Collaboration), J. High Energy Phys. 09 (2016) 001.