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Discovering walking technirho mesons at the LHC

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We formulate a scale-invariant hidden local symmetry (HLS) as a low-energy effective theory of walking technicolor (WTC) which includes the technidilaton, technipions, and technirho mesons as the low-lying spectra. As a benchmark for LHC phenomenology, we in particular focus on the one-family model of WTC having eight technifermion flavors, which can be—at energy scales relevant to the reach of the LHC—described by the scale-invariant HLS based on the manifold $[\mathrm{SU}(8)_L \times \mathrm{SU}(8)_R]_{\mathrm{global}} \times \mathrm{SU}(8)_{\mathrm{local}}/\mathrm{SU}(8)_V$, where $\mathrm{SU}(8)_{\mathrm{local}}$ is the HLS and the global $\mathrm{SU}(8)_L \times \mathrm{SU}(8)_R$ symmetry is partially gauged by the $\mathrm{SU}(3) \times \mathrm{SU}(2)_L \times U(1)_Y$ of the standard model. Based on the scale-invariant HLS, we evaluate the coupling properties of the technirho mesons and place limits on the masses from the current LHC data. Then, implications for future LHC phenomenology are discussed by focusing on the technirho mesons produced through the Drell-Yan process. We find that the color-octet technirho decaying to the technidilaton along with the gluon is of interest as the discovery channel at the LHC, which would provide a characteristic signature to probe the one-family WTC.

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I. INTRODUCTION

The Higgs boson with a mass of 125 GeV was discovered at the LHC. However, the dynamical origins of electroweak (EW) symmetry breaking and of the Higgs are still mysterious and may be explained by physics beyond the standard model (SM). Technicolor (TC) [1–3] is a well-motivated model for the dynamical origin of EW symmetry breaking in a way similar to the established mechanism in QCD which breaks the chiral symmetry (and hence the EW symmetry as well) dynamically via the fermion pair condensate. However, the original TC was ruled out a long time ago by the notorious flavor-changing neutral current (FCNC) problem, and more dramatically by the recent discovery of the 125 GeV Higgs which cannot be accounted for by the TC dynamics of a simple QCD scale-up.

Fortunately, both problems are simultaneously solved by walking technicolor (WTC) [4,5], which was proposed based on the scale-symmetric dynamics of the ladder Schwinger-Dyson equation: with the scale symmetry, WTC predicted a large anomalous dimension $\gamma_m = 1$ as a solution to the FCNC problem,[1] and at the same time predicted a light composite Higgs—known as a "technidilaton" (TD) [4,7]—that is a pseudo-Nambu-Goldstone (NG) boson of the scale symmetry that is broken spontaneously (and also explicitly) by the technifermion condensate. It was shown that the TD can account for the 125 GeV Higgs, with couplings that are consistent with the current LHC data of the 125 GeV Higgs (see below) [8–11]. Thus the origin of the Higgs mass is dynamically explained by the scale of the chiral condensate in WTC.

The mass of the TD as a pseudo-NG boson comes from the nonperturbative trace anomaly due to the chiral condensate and can be estimated through the partially conserved dilatation current (PCDC) relation [4,7]. A precise ladder evaluation of $m_\phi F_\phi$ based on this PCDC relation reads [12] $(m_\phi F_\phi)^2 \simeq 0.154 \cdot N_f N_c \cdot m_D^4 \simeq (2.5 \cdot v_{\mathrm{EW}}^2)^2 \cdot [(8/N_f)(4/N_c)]$, where $v_{\mathrm{EW}}^2 = (246\ \mathrm{GeV})^2 = N_D F_\pi^2 \simeq 0.028 \cdot N_f N_c \cdot m_D^2$ (Pagels-Stokar formula), with $N_D (= N_f/2)$ being the number of the electroweak doublets. Note the scaling $m_\phi / v_{\mathrm{EW}} \sim 1/\sqrt{N_f N_c}$, which implies that a light TD $m_\phi / v_{\mathrm{EW}} \ll 1$ is naturally realized for $N_f \gg 1$ (as well as $N_c \gg 1$) as in large-N_f QCD.[2,3]

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[1]A similar solution to the FCNC problem was given without the notion of the scale symmetry/technidilaton and the anomalous dimension [6].

[2]Thus the mass of the LHC Higgs, $m_\phi \simeq 125\ \mathrm{GeV} \simeq v_{\mathrm{EW}}/2$, can be obtained [8] when we take $v_{\mathrm{EW}}/F_\phi = 2F_\pi/F_\phi \simeq 1/5 = 0.2$ ($v_{\mathrm{EW}} = 2F_\pi$ for $N_c = 4, N_f = 8$ in the one-family model; see below). Amazingly, this value of F_ϕ turned out to be consistent with the LHC Higgs data (best fit: $v_{\mathrm{EW}}/F_\phi \simeq 0.22$) [10].

[3]One might think that such a large N_f (and N_c) would result in the so-called S-parameter problem [13]. The large S from the TC sector, however, is not necessarily in conflict with the experimental value of S from the precision EW measurements, since the contributions from the TC sector can easily be canceled by strong mixing with the SM fermion contribution through the ETC interactions, as in the fermion delocalization of the Higgsless model [14]. Moreover, even within the TC sector alone there exists a way to resolve this problem, as demonstrated in the holographic model, where we can reduce $S \sim 4\pi(N_D F_\pi^2)/M_\rho^2 = 4\pi v_{\mathrm{EW}}^2/M_\rho^2$ by tuning the holographic parameter (roughly corresponding to increasing the technirho mass M_ρ) in a way consistent with the TD mass of 125 GeV and all the current LHC data for the 125 GeV Higgs [11].

More recently, in another approach using holographic WTC [15] it was shown [11] that the strong gluon dynamics via the large technigluon condensate can realize a parametrically massless TD limit $m_\phi/v_{\rm EW} \to 0+$ and hence naturally realize $m_\phi \simeq v_{\rm EW}/2 \simeq 125$ GeV, consistent with the LHC Higgs data. Similar arguments for realizing a parametrically light dilaton were given in somewhat different contexts [16].

Amazingly, the recent lattice results [17] in fact indicate that the SU(3) gauge theory with eight fundamental fermions ($N_f = 8$ QCD) possesses a walking nature, with the anomalous dimension $\gamma_m \simeq 1$. Most remarkably, it has been shown in lattice $N_f = 8$ QCD [18] that there in fact exists a flavor-singlet scalar meson that is as light as the pions for a small fermion-mass region, which thus can be a composite Higgs (in the form of the TD) in the chiral limit.

Thus, special interest in the context of lattice studies has recently been paid to the one-family model (or Farhi-Susskind model) as a candidate theory for WTC [3,19]. The model is the most straightforward version of the extended TC (ETC) model [20] which incorporates the mechanism of producing masses for the SM fermions. The one-family model consists of $N_{\rm TC}$ copies of a whole generation of the SM fermions; therefore, the TC sector of the model is a SU($N_{\rm TC}$) gauge theory with eight fundamental Dirac fermions, $N_f = 8$ (i.e., four weak doublets, $N_D = N_f/2 = 4$, with the NG boson decay constant $F_\pi \simeq 246\ {\rm GeV}/\sqrt{N_D} = 123$ GeV).

The global chiral symmetry-breaking pattern of the one-family model is $G/H = {\rm SU}(8)_L \times {\rm SU}(8)_R/{\rm SU}(8)_V$, which is described by the usual nonlinear chiral Lagrangian based on the manifold G/H in terms of the 63 NG bosons. It is further straightforwardly extended to a scale-invariant version so as to incorporate the TD, ϕ, as a composite Higgs [9]. The chiral perturbation theory (ChPT) of the scale-invariant version can also be formulated by assigning the chiral order counting $m_\phi^2 = \mathcal{O}(p^2)$ [21].

Three of the 63 NG bosons are eaten by the SM weak gauge bosons when the SM gauge interactions are switched on, while the other 60 remain as physical states ("technipions"), all acquiring mass to become pseudo-NG bosons by the SM gauging and the ETC gauging, which explicitly breaks the chiral symmetry G down to the symmetry corresponding to the EW symmetry. The gauge couplings of these explicit breakings are small and perturbative, and therefore the masses of the technipions may be estimated by the Dashen formula at the lowest-order perturbation, just like the estimate of the $\pi^+ - \pi^0$ mass difference in QCD. It turns out that the masses of all the technipions are drastically enhanced by the walking dynamics of WTC [4,22,23]. In the case of the one-family model, the masses of the walking technipions were explicitly estimated [24,25] to be of $\mathcal{O}$ (TeV) (see Sec. III below), suggesting a new possibility that the technirhos decay directly to the SM particles, rather than through the technipions.

In this paper, we consider another type of technihadron—vector resonances (here we are confined to the flavor-nonsinglet ones, called "technirho mesons")—which are expected to exist as typical bound states in the generic dynamical EW symmetry-breaking scenarios, not restricted to WTC or the one-family model. In the case of WTC we extend the scale-invariant version of the low-energy effective theory [8–10]—i.e., the case of the one-family model based on $G/H = {\rm SU}(8)_L \times {\rm SU}(8)_R/{\rm SU}(8)_V$ [24]—in a way that includes the technirhos by using the hidden local symmetry (HLS) [26,27].[4] Similar discussions of the technirhos based on the HLS were done for the one-doublet model with $G/H = {\rm SU}(2)_L \times {\rm SU}(2)_R/{\rm SU}(2)_V$ without the scale symmetry/TD [the "breaking electroweak symmetry strongly" (BESS) model] [28]. This is the first study of the HLS for the one-family model as well as its scale-invariant version, which implies a novel salient LHC phenomenology involving the TD and the colored technirho. It is to be noted that the ChPT was formulated for the HLS Lagrangian [27] and can be extended to the scale-invariant version of the HLS in the same way as in the case without the HLS [21], although we do not include the loop effects in this paper. Based on the scale-invariant HLS, we first evaluate the constraint on the technirho meson masses from the current LHC data, and then future LHC phenomenology is discussed. We find that the color-octet technirho produced via the Drell-Yan (DY) process which decays to the TD along with the gluon is especially interesting as a discovery channel at the LHC to probe the one-family WTC.

The paper is organized as follows. In Sec. II, we formulate the scale-invariant HLS based on the manifold $[{\rm SU}(8)_L \times {\rm SU}(8)_R]_{\rm global} \times {\rm SU}(8)_{\rm local}/{\rm SU}(8)_V$, including the TD, technipions, and technirho mesons. In Sec. III the decay widths and branching ratios of the technirho mesons in the one-family WTC are discussed. In Sec. IV we explore the LHC phenomenology of the technirho mesons and place limits on the masses from the currently available LHC data on searches for new spin-1 resonances. We then discuss the discovery channels of the technirho mesons, which include the TD as the daughter particle of the parent technirho mesons. A summary is given in Sec. V. The explicit forms of the technirho couplings relevant to the LHC phenomenology and the partial decay widths are presented in Appendices A and B, respectively.

[4] The flavor-singlet technivector meson can be incorporated into the HLS Lagrangian by taking $G/H = {\rm U}(8)_L \times {\rm U}(8)_R/{\rm U}(8)_V$ instead of $G/H = {\rm SU}(8)_L \times {\rm SU}(8)_R/{\rm SU}(8)_V$. As we discuss later, however, it would be less interesting compared with the technirho in the LHC phenomenology and is not discussed in this paper.

II. THE SCALE-INVARIANT HLS FOR WALKING TECHNIRHO MESONS OF THE ONE-FAMILY MODEL

In this section, based on the HLS formalism [26,27], we give a formulation of the scale-invariant HLS for the one-family WTC model, which includes the TD and technipions, as well as the technirho mesons as the HLS gauge bosons. The HLS formalism makes it straightforward to simultaneously incorporate both the external (SM) gauge and HLS (vector-meson) interactions, in contrast to other approaches for the inclusion of the vector mesons into the chiral Lagrangian. It actually turns out to be crucial for studying the LHC phenomenology, as will be seen below. Furthermore, the ChPT for the systematic loop expansion has been developed only in the HLS formalism, although the vector mesons can also be incorporated into the chiral Lagrangian by other formalisms which are equivalent to the HLS Lagrangian at the on-shell tree level [27].

The Lagrangian for the technipion is expressed as the usual nonlinear sigma model based on the manifold $G/H = \mathrm{SU}(8)_L \times \mathrm{SU}(8)_R/\mathrm{SU}(8)_V$. The TD is incorporated by forcing the chiral effective theory to be scale-invariant through the introduction of the compensating nonlinear field $\chi(x) = e^{\phi(x)/F_\phi}$, which transforms under scale transformation as $\delta\chi(x) = (1 + x^\nu \partial_\nu)\chi(x)$, so that $\phi(x)$ does scale nonlinearly as $\delta\phi(x) = F_\phi + x^\nu \partial_\nu \phi(x)$, where $\phi(x)$ and F_ϕ are the TD field and its decay constant, respectively [8–10]. The resultant one-family scale-invariant action is explicitly given by the Lagrangian [24]

$$\mathcal{L} = \frac{F_\pi^2}{4} \cdot \chi^2(x) \cdot \mathrm{tr}[\mathcal{D}_\mu U^\dagger \mathcal{D}^\mu U] + \frac{F_\phi^2}{2} \partial_\mu \chi(x) \partial^\mu \chi(x), \qquad (1)$$

where $U(x) = e^{i\frac{2\pi(x)}{F_\pi}}$ (with F_π being the decay constant of the NG bosons) transforms as $U \to g_L \cdot U \cdot g_R^\dagger$, with $(g_L, g_R) \in G = \mathrm{SU}(8)_L \times \mathrm{SU}(8)_R$, as does the covariant derivative $\mathcal{D}^\mu U(x) = \partial^\mu U(x) - i\mathcal{L}^\mu(x) U(x) + iU(x)\mathcal{R}^\mu(x)$, where G is formally fully gauged by the external gauge fields $(\mathcal{L}^\mu(x), \mathcal{R}^\mu(x))$. The action for the Lagrangian (1) is invariant under the gauged-G symmetry as well as the scale symmetry. In the realistic application to WTC, the external gauge fields are restricted to the SM gauge fields of $\mathrm{SU}(3) \times \mathrm{SU}(2)_L \times U(1)_Y$. Besides the scale-invariant term, there exists a scale-anomaly term which reproduces the TD mass and as well as terms involving the TD coupling to the SM fields [9].

Now, it is straightforward to introduce the technirho into the Lagrangian (1) in the standard manner of the HLS formalism [26]. The HLS can be made explicit by writing $U(x) = e^{i\frac{2\pi(x)}{F_\pi}} = \xi_L^\dagger(x) \cdot \xi_R(x)$, where the $\xi_{L,R}$ are parametrized as

$$\xi_{L,R}(x) = e^{\frac{i\sigma(x)}{F_\sigma}} e^{\mp\frac{i\pi(x)}{F_\pi}},$$
$$(\pi(x) = \pi^A(x) X^A, \qquad \sigma(x) = \sigma^A(x) X^A), \qquad (2)$$

with the broken generators X^A and the fictitious NG bosons $\sigma^A(x)$ (not to be confused with the scalar meson) along with the decay constant F_σ, which are to be absorbed into the HLS. The $\xi_{L,R}(x)$ transform as $\xi_{L,R} \to h(x) \cdot \xi_{L,R} \cdot g_{L,R}^\dagger$ under $G_{\mathrm{global}} \times H_{\mathrm{local}} = [\mathrm{SU}(8)_L \times \mathrm{SU}(8)_R]_{\mathrm{global}} \times \mathrm{SU}(8)_{\mathrm{local}}$, where $h \in H_{\mathrm{local}} = \mathrm{SU}(8)_{\mathrm{local}}$ is the HLS, and $g_{L,R} \in G_{\mathrm{global}} = [\mathrm{SU}(8)_L \times \mathrm{SU}(8)_R]_{\mathrm{global}}$. When we fix the gauge [unitary gauge $\sigma(x) = 0$] as $\xi_L^\dagger(x) = \xi_R(x) = \xi(x) = e^{i\frac{\pi(x)}{F_\pi}}$, H_{local} and $H_{\mathrm{global}} (\subset G_{\mathrm{global}})$ are both spontaneously broken down to a single H which is a diagonal sum of both of them, and accordingly G_{global} is reduced back to the original chiral symmetry G in the model based on G/H: ξ transforms as $\xi \to h(g,\pi)\xi g_{L,R}^\dagger$, with $h(g,\pi)$ being the $\pi(x)$-dependent (global) H transformation of G/H.

The technirho mesons are introduced as the gauge bosons of the HLS $H_{\mathrm{local}} = \mathrm{SU}(8)_{\mathrm{local}}$ through the covariant derivative $D_\mu \xi_{L,R}(x) = \partial_\mu \xi_{L,R}(x) - iV_\mu(x)\xi_{L,R}(x) + i\xi_{L,R}\mathcal{L}_\mu(x)(\mathcal{R}_\mu(x))$, with the HLS gauge field V_μ and the external gauge fields $\mathcal{L}_\mu$ and $\mathcal{R}_\mu$. G_{global} is again fully gauged for formal discussion to make the invariance transparent. The resulting form of the Lagrangian is as follows:

$$\mathcal{L} = \chi^2(x) \cdot (F_\pi^2 \mathrm{tr}[\hat{\alpha}_{\mu\perp}^2] + F_\sigma^2 \mathrm{tr}[\hat{\alpha}_{\mu\parallel}^2]) - \frac{1}{2g^2} \mathrm{tr}[V_{\mu\nu}^2], \qquad (3)$$

where

$$\hat{\alpha}_{\mu\perp,\parallel} = \frac{D_\mu \xi_R \cdot \xi_R^\dagger \mp D_\mu \xi_L \cdot \xi_L^\dagger}{2i}. \qquad (4)$$

The covariantized Maurer-Cartan 1-forms $\hat{\alpha}_{\mu\perp,\parallel}$ transform as $\hat{\alpha}_{\mu\perp,\parallel} \to h(x) \cdot \hat{\alpha}_{\mu\perp,\parallel} \cdot h^\dagger(x)$. Without the kinetic term of the HLS gauge fields $V_\mu(x)$ (namely, by integrating out the V_μ), the Lagrangian is reduced to the nonlinear sigma model based on G/H in the unitary gauge $\sigma(x) = 0$ $(\xi_L^\dagger(x) = \xi_R(x) = \xi(x) = e^{i\frac{\pi(x)}{F_\pi}})$.

The 63 chiral NG bosons are embedded in the adjoint representation of the SU(8) group [24]:

$$\begin{aligned}\sum_{A=1}^{63} \pi^A(x) X^A &= \sum_{i=1}^{3} \Pi^i(x) X_\Pi^i + \sum_{i=1}^{3} P^i(x) X_P^i + P^0(x) X_P \\ &+ \sum_{i=1}^{3}\sum_{a=1}^{8} \theta_a^i(x) X_{\theta a}^i + \sum_{a=1}^{8} \theta_a^0(x) X_{\theta a} \\ &+ \sum_{c=r,g,b}\sum_{i=1}^{3} [T_c^i(x) X_{Tc}^i + \bar{T}_c^i(x) X_{\bar{T}c}^i] \\ &+ \sum_{c=r,g,b} [T_c^0(x) X_{Tc} + \bar{T}_c^0(x) X_{\bar{T}c}], \qquad (5)\end{aligned}$$

where $(\tau^i = \sigma^i/2)$

$$X^i_\Pi = \frac{1}{2}\left(\begin{array}{c|c} \tau^i \otimes \mathbf{1}_{3\times 3} & \\ \hline & \tau^i \end{array}\right),$$

$$X^i_P = \frac{1}{2\sqrt{3}}\left(\begin{array}{c|c} \tau^i \otimes \mathbf{1}_{3\times 3} & \\ \hline & -3\cdot\tau^i \end{array}\right),$$

$$X_P = \frac{1}{4\sqrt{3}}\left(\begin{array}{c|c} \mathbf{1}_{6\times 6} & \\ \hline & -3\times\mathbf{1}_{2\times 2} \end{array}\right),$$

$$X^i_{\theta a} = \frac{1}{\sqrt{2}}\left(\begin{array}{c|c} \tau^i \otimes \lambda_a & \\ \hline & 0 \end{array}\right),$$

$$X_{\theta_a} = \frac{1}{2\sqrt{2}}\left(\begin{array}{c|c} \mathbf{1}_{2\times 2} \otimes \lambda_a & \\ \hline & 0 \end{array}\right),$$

$$X^i_{Tc} = \frac{1-i}{2}\left(\begin{array}{c|c} & \tau^i \otimes \mathbf{e}_c \\ \hline \tau^i \otimes \mathbf{e}_c^\dagger & \end{array}\right), \quad X^i_{\bar{T}c} = (X^i_{Tc})^\dagger,$$

$$X_{Tc} = \frac{1-i}{4}\left(\begin{array}{c|c} & \mathbf{1}_{2\times 2} \otimes \mathbf{e}_c \\ \hline \mathbf{1}_{2\times 2} \otimes \mathbf{e}_c^\dagger & \end{array}\right), \quad X_{\bar{T}c} = (X_{Tc})^\dagger, \tag{6}$$

with $\mathbf{e}_c$ being a three-dimensional unit vector in color space and the generators are normalized as $\mathrm{Tr}[X^A X^B] = \delta^{AB}/2$. Among the above, Π^i become longitudinal degrees of freedom of the SM $W^\pm$ and Z bosons. It is convenient to express π in a blocked 8×8 matrix form as

$$\pi^A X^A = \left(\begin{array}{c|c} (\pi_{QQ})_{6\times 6} & (\pi_{QL})_{2\times 6} \\ \hline (\pi_{LQ})_{6\times 2} & (\pi_{LL})_{2\times 2} \end{array}\right), \tag{7}$$

where

$$\pi_{QQ} = \left[\sqrt{2}\theta + \frac{1}{\sqrt{2}}\theta^0\right] + \left(\frac{1}{2}\Pi + \frac{1}{2\sqrt{3}}P + \frac{1}{4\sqrt{3}}P^0\right)\otimes\mathbf{1}_{3\times 3},$$

$$\pi_{QL} = T + \frac{1}{2}T^0,$$

$$\pi_{LQ} = \pi^\dagger_{QL} = \bar{T} + \frac{1}{2}\bar{T}^0,$$

$$\pi_{LL} = \left(\frac{1}{2}\Pi - \frac{\sqrt{3}}{2}P - \frac{\sqrt{3}}{4}P^0\right),$$

$$\theta = \theta^i_a \tau^i \frac{\lambda_a}{2}, \quad \theta^0 = \theta^0_a \cdot \mathbf{1}_{2\times 2}\cdot\frac{\lambda_a}{2},$$

$$T = T^i_c \mathbf{e}_c \tau^i, \quad T^0 = T^0_c \mathbf{e}_c,$$

$$P = P^i \tau^i, \quad P^0 = P^0 \cdot \mathbf{1}_{2\times 2},$$

$$\Pi = \Pi^i \tau^i.$$

The technirho meson fields are also parametrized in a way similar to π:

$$\begin{aligned}\sum_{A=1}^{63}\rho^A(x)X^A &= \sum_{i=1}^{3}\rho^i_\Pi(x)X^i_\Pi + \sum_{i=1}^{3}\rho^i_P(x)X^i_P + \rho^0_P(x)X_P \\ &\quad + \sum_{i=1}^{3}\sum_{a=1}^{8}\rho^i_{\theta a}(x)X^i_{\theta a} + \sum_{a=1}^{8}\rho^0_{\theta_a}(x)X_{\theta a} \\ &\quad + \sum_{c=r,g,b}\sum_{i=1}^{3}[\rho^i_{Tc}(x)X^i_{Tc} + \bar{\rho}^i_{Tc}(x)X^i_{\bar{T}c}] \\ &\quad + \sum_{c=r,g,b}[\rho^0_{Tc}(x)X_{Tc} + \bar{\rho}^0_{T_c}(x)X_{\bar{T}c}].\end{aligned} \tag{8}$$

They are embedded in a 8×8 block-diagonal form, $V_\mu = V^A_\mu X^A$, as

$$\rho^\mu = \frac{V^{\mu A}X^A}{g} = \left(\begin{array}{c|c} (\rho^\mu_{QQ})_{6\times 6} & (\rho^\mu_{QL})_{2\times 6} \\ \hline (\rho^\mu_{LQ})_{6\times 2} & (\rho^\mu_{LL})_{2\times 2} \end{array}\right), \tag{9}$$

with

$$\begin{aligned}\rho^\mu_{QQ} &= \left[\sqrt{2}\rho^\mu_\theta + \frac{1}{\sqrt{2}}\rho^{\mu 0}_\theta\right] \\ &\quad + \left(\frac{1}{2}\rho^\mu_\Pi + \frac{1}{2\sqrt{3}}\rho^\mu_P + \frac{1}{4\sqrt{3}}\rho^{0\mu}_P\right)\otimes\mathbf{1}_{3\times 3}, \\ \rho^\mu_{QL} &= \rho^\mu_T + \frac{1}{2}\rho^{0\mu}_T, \\ \rho^\mu_{LQ} &= (\rho^\mu_{QL})^\dagger = \bar{\rho}^\mu_T + \frac{1}{2}\bar{\rho}^{0\mu}_T, \\ \rho^\mu_{LL} &= \left(\frac{1}{2}\rho^\mu_\Pi - \frac{\sqrt{3}}{2}\rho^\mu_P - \frac{\sqrt{3}}{4}\rho^{0\mu}_P\right),\end{aligned}$$

$$\rho^\mu_\theta = \rho^{i\mu}_{\theta a}\tau^i\frac{\lambda_a}{2}, \qquad \rho^{0\mu}_\theta = \rho^{0\mu}_{\theta a}\cdot\mathbf{1}_{2\times 2}\cdot\frac{\lambda_a}{2},$$

$$\rho^\mu_T = \rho^{i\mu}_{Tc}\mathbf{e}_c\tau^i, \qquad \rho^{0\mu}_T = \rho^{0\mu}_{Tc}\mathbf{e}_c,$$

$$\rho^\mu_P = \rho^{i\mu}_P\tau^i, \qquad \rho^{0\mu}_P = \rho^{0\mu}_P\cdot\mathbf{1}_{2\times 2},$$

$$\rho^\mu_\Pi = \rho^{i\mu}_\Pi\tau^i.$$

Here we used the same basis of the $\mathrm{SU}(8)_V$ matrix as that of π since ρ^i_Π are produced only by the Drell-Yan processes through the mixing with W and Z bosons, which absorb Π^i. With this base, the other color-singlet isotriplet technirhos, ρ^i_P, are not produced due to the orthogonality of $\mathrm{tr}[X^i_\Pi X^j_P] = 0$, as will be seen later.

The external gauge fields $\mathcal{L}_\mu$ and $\mathcal{R}_\mu$ involve the $\mathrm{SU}(3)_c$, $\mathrm{SU}(2)_W$ and $U(1)_Y$ gauge fields (G_μ, W_μ, B_μ) in the SM as follows:

$$\mathcal{L}_\mu = 2g_W W^i_\mu X^i_\Pi + \frac{2}{\sqrt{3}}g_Y B_\mu X_P + \sqrt{2}g_s G^a_\mu X_{\theta_a},$$

$$\mathcal{R}_\mu = 2g_Y B_\mu\left(X^3_\Pi + \frac{1}{\sqrt{3}}X_P\right) + \sqrt{2}g_s G^a_\mu X_{\theta_a}. \tag{10}$$

Through the standard diagonalization procedure, the left and right gauge fields are expressed in terms of the mass eigenstates $(W^{\pm}, Z, \gamma, g)$ as

$$\begin{aligned}\mathcal{L}_\mu &= g_s G^a_\mu \Lambda^a + eQ_{\rm em}A_\mu + \frac{e}{sc}(I_3 - s^2 Q_{\rm em})Z_\mu \\ &\quad + \frac{e}{\sqrt{2}s}(W^+_\mu I^+ + {\rm H.c.}),\\ \mathcal{R}_\mu &= g_s G^a_\mu \Lambda^a + eQ_{\rm em}A_\mu - \frac{es}{c}Q_{\rm em}Z_\mu,\end{aligned} \tag{11}$$

where s ($c^2 = 1 - s^2$) denotes the standard weak mixing angle defined by $g_W = e/s$ and $g_Y = e/c$, and

$$\begin{aligned}\Lambda_a &= \sqrt{2}X_{\theta_a}, \qquad I_3 = 2X^3_\Pi, \qquad Q_{\rm em} = I_3 + Y,\\ Y &= \frac{2}{\sqrt{3}}X_P, \qquad I_+ = 2(X^1_\Pi + iX^2_\Pi), \qquad I_- = (I_+)^\dagger.\end{aligned} \tag{12}$$

It is convenient to define the vector and axial-vector gauge fields $\mathcal{V}_\mu$ and $\mathcal{A}_\mu$ as

$$\mathcal{V}_\mu = \frac{\mathcal{R}_\mu + \mathcal{L}_\mu}{2}, \qquad \mathcal{A}_\mu = \frac{\mathcal{R}_\mu - \mathcal{L}_\mu}{2}, \tag{13}$$

so that they are expressed in a blocked-8 × 8 matrix form:

$$\mathcal{V}^\mu = \left(\begin{array}{c|c} (\mathcal{V}^\mu_{QQ})_{6\times 6} & \mathbf{0}_{2\times 6} \\ \mathbf{0}_{6\times 2} & (\mathcal{V}^\mu_{LL})_{2\times 2}\end{array}\right), \tag{14}$$

$$\mathcal{A}^\mu = \left(\begin{array}{c|c} (\mathcal{A}^\mu_{QQ})_{6\times 6} & \mathbf{0}_{2\times 6} \\ \mathbf{0}_{6\times 2} & (\mathcal{A}^\mu_{LL})_{2\times 2}\end{array}\right), \tag{15}$$

where

$$\begin{aligned}\mathcal{V}^\mu_{QQ} &= \mathbf{1}_{2\times 2}\cdot g_s G^\mu_a \frac{\lambda_a}{2}\\ &\quad + \left[eQ^q_{\rm em}A^\mu + \frac{e}{2sc}z^q_V Z^\mu + \frac{e}{2\sqrt{2}s}(\tau^+ W^{\mu +} + {\rm H.c.})\right]\\ &\quad \cdot \mathbf{1}_{3\times 3},\\ \mathcal{V}^\mu_{LL} &= \left[eQ^l_{\rm em}A^\mu + \frac{e}{2sc}z^l_V Z^\mu + \frac{e}{2\sqrt{2}s}(\tau^+ W^{\mu +} + {\rm H.c.})\right],\\ \mathcal{A}^\mu_{QQ} &= -\left(\frac{e}{2sc}\tau^3 Z^\mu + \frac{e}{2\sqrt{2}s}(\tau^+ W^{\mu +} + {\rm H.c.})\right)\cdot \mathbf{1}_{3\times 3},\\ \mathcal{A}^\mu_{LL} &= -\left(\frac{e}{2sc}\tau^3 Z^\mu + \frac{e}{2\sqrt{2}s}(\tau^+ W^{\mu +} + {\rm H.c.})\right),\\ Q^q_{\rm em} &= \begin{pmatrix} 2/3 & 0 \\ 0 & -1/3\end{pmatrix}, \qquad Q^l_{\rm em} = \begin{pmatrix} 0 & 0 \\ 0 & -1\end{pmatrix},\\ z^{q,l}_V &= \tau^3 - 2s^2 Q^{q,l}_{\rm em}, \qquad \tau^+ = \begin{pmatrix} 0 & 1 \\ 0 & 0\end{pmatrix}, \qquad \tau^- = (\tau^+)^\dagger.\end{aligned}$$

Interactions among technihadrons and SM particles can be obtained by expanding the Lagrangian in Eq. (3) in powers of π and ϕ with the HLS fixed as unitary gauge ($\sigma = 0$). Basic HLS relations include [26]

$$M^2_\rho = ag^2F^2_\pi, \tag{16}$$

$$g_{\rho\pi\pi} = \frac{a}{2}g, \tag{17}$$

$$g_{\mathcal{V}\pi\pi} = \left(1 - \frac{a}{2}\right), \tag{18}$$

with

$$a \equiv \frac{F^2_\sigma}{F^2_\pi}, \tag{19}$$

where M_ρ, $g_{\rho\pi\pi}$, and $g_{\gamma\pi\pi}$ have been read off from the following terms:

$$\mathcal{L}_{M^2_\rho} = M^2_\rho {\rm tr}[\rho_\mu \rho^\mu], \tag{20}$$

$$\mathcal{L}_{\mathcal{V}\pi\pi} = 2ig_{\mathcal{V}\pi\pi}{\rm tr}[\mathcal{V}^\mu[\partial_\mu \pi, \pi]], \tag{21}$$

$$\mathcal{L}_{\rho\pi\pi} = 2ig_{\rho\pi\pi}{\rm tr}[\rho^\mu[\partial_\mu \pi, \pi]]. \tag{22}$$

Note that, from Eq. (18), direct couplings of SM gauge bosons to two pions vanish when we take $a = 2$:

$$g_{\mathcal{V}\pi\pi} = 0 \quad \text{for } a = 2. \tag{23}$$

In that case, the couplings of the SM gauge bosons ($\mathcal{V}_\mu$) to two pions arise only from the ρ-$\mathcal{V}$ mixing (vector-meson dominance).

The explicit forms of interactions relevant to the current study are summarized in Appendix A. Among these interaction terms, the most relevant terms in the present study are $\mathcal{V}$-ρ and $\mathcal{V}$-ρ-ϕ vertices. These terms arise from the $\chi^2 F_\sigma^2 \mathrm{tr}[\hat{\alpha}_{\mu\|}^2]$ term in Eq. (3) by expanding χ and $\hat{\alpha}_{\mu\|}$ in terms of the technidilaton (ϕ) and pion (π) fields:

$$\chi^2 F_\sigma^2 \mathrm{tr}[\hat{\alpha}_{\mu\|}^2] = F_\sigma^2 \left(1 + \frac{2\phi}{F_\phi} + \cdots\right) \times \mathrm{tr}\left[(\mathcal{V}_\mu - V_\mu)^2 + \frac{i}{F_\pi^2} V_\mu [\partial^\mu \pi, \pi] + \frac{2i}{F_\pi} V_\mu [\mathcal{A}^\mu, \pi] + \cdots\right]. \tag{24}$$

From the first term in the square bracket, and by using Eqs. (9) and (14), one can readily read off the $\mathcal{V}$-ρ and $\mathcal{V}$-ρ-ϕ vertices as

$$\begin{aligned}\mathcal{L}_{\mathcal{V}\rho} &= -2gF_\sigma^2 \mathrm{tr}[\mathcal{V}_\mu \rho^\mu] \\ &= -2gF_\sigma^2 \left[\frac{g_s}{\sqrt{2}} G_\mu^a \rho_\theta^{0a\mu} + eA_\mu \left\{\rho_\Pi^{3\mu} + \frac{1}{\sqrt{3}} \rho_P^{0\mu}\right\} + \frac{e}{2sc} Z_\mu \left\{(c^2 - s^2)\rho_\Pi^{3\mu} - \frac{2}{\sqrt{3}} s^2 \rho_P^{0\mu}\right\} + \frac{e}{2s}\{W_\mu^- \rho_\Pi^{\mu+} + \mathrm{H.c.}\}\right]\end{aligned} \tag{25}$$

and

$$\begin{aligned}\mathcal{L}_{\mathcal{V}\rho\phi} &= -\frac{4gF_\sigma^2}{F_\phi} \phi \mathrm{tr}[\mathcal{V}_\mu \rho^\mu] \\ &= -\frac{4gF_\sigma^2}{F_\phi} \phi \left[\frac{g_s}{\sqrt{2}} G_\mu^a \rho_\theta^{0a\mu} + eA_\mu \left\{\rho_\Pi^{3\mu} + \frac{1}{\sqrt{3}} \rho_P^{0\mu}\right\} + \frac{e}{2sc} Z_\mu \left\{(c^2 - s^2)\rho_\Pi^{3\mu} - \frac{2}{\sqrt{3}} s^2 \rho_P^{0\mu}\right\} + \frac{e}{2s}\{W_\mu^- \rho_\Pi^{\mu+} + \mathrm{H.c.}\}\right].\end{aligned} \tag{26}$$

Here we have defined the charged rho-meson fields as

$$\rho_\Pi^{\mu\pm} = \frac{\rho_\Pi^{1\mu} \mp i\rho_\Pi^{2\mu}}{\sqrt{2}}. \tag{27}$$

Note the absence of $A - \rho_P^3, Z - \rho_P^3$, $W^\pm - \rho_P^\mp$, $A - \rho_P^3 - \phi, Z - \rho_P^3 - \phi$, and $W^\pm - \rho_P^\mp - \phi$ terms due to the orthogonality of the $\mathrm{SU}(8)_V$ generators. As will be discussed more explicitly, the terms in Eq. (25) are crucial for technirho-meson productions through the Drell-Yan process and decays to the SM fermions via the vector-meson dominance, as in Eq. (23). The terms in Eq. (26) are relevant for decays which involve the Higgs boson (TD).

III. DECAY WIDTHS AND BRANCHING RATIOS OF TECHNIRHO MESONS

Having formulated the low-energy effective Lagrangian of the one-family model and derived relevant interactions among technihadrons and SM fields, in this section we study the decay widths and branching ratios of the technirho mesons and related collider phenomenology that are based on them. The Lagrangian in Eq. (3) has four parameters: F_π, F_σ, F_ϕ, and g. We fix $F_\pi = 123$ GeV, which is set by the EW scale $v_{\mathrm{EW}} \simeq 246$ GeV through the relation

$$F_\pi = \frac{v_{\mathrm{EW}}}{\sqrt{N_D}} \simeq 123 \text{ GeV} \quad (\text{for } N_D = 4). \tag{28}$$

As for the TD decay constant F_ϕ, we use the best-fit value by which the TD can be fitted well to the current LHC Higgs data (see Ref. [10]):

$$F_\phi = F_\phi|_{\mathrm{best}} \simeq v_{\mathrm{EW}}/0.22. \tag{29}$$

Also, we impose vector-meson dominance, which is achieved by taking $a = F_\sigma^2/F_\pi^2$ to be [see Eq. (23)]

$$a = 2. \tag{30}$$

The remaining HLS parameter, g, will be fixed by the input value of the technirho mass, M_ρ, through the relation in Eq. (16).

In the next section, we will study technirho production through the mixing with the SM gauge bosons which are produced by the Drell-Yan process.[5] The types of technirho mesons produced by such a process are ρ_θ^0, ρ_P^0, and $\rho_\Pi^{\pm,3}$, as illustrated in Fig. 1. Note that $\rho_P^{\pm,3}$ are not produced via the Drell-Yan process because they do not mix with the SM gauge bosons due to the orthogonality of the $\mathrm{SU}(8)_V$ symmetry [see Eq. (A2) in Appendix A]. Thus, in this

[5]It should be noted that we do not include the gluon-gluon fusion process for ρ_θ^0 production since there is an accidental cancellation of the $g - g - \rho_\theta^0$ on-shell amplitude due to the presence of the contribution from the non-Abelian $\rho_\theta^0 - \rho_\theta^0 - \rho_\theta^0$ vertex [29,30]. [This is true for the leading and the next-to-leading order in the derivative expansion of the Lagrangian in Eq. (3).] Also, we have no $g - \rho_T - \rho_T$ vertex at the leading order due to the $\mathrm{SU}(8)_V$ invariance, so the current limit [31] on vector leptoquarks through the Drell-Yan process is not applicable to ρ_T.

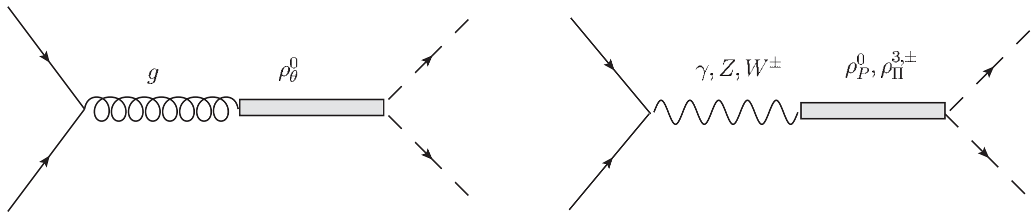

FIG. 1. An illustration of the relevant LHC production and two-body decay processes for $\rho^0_\theta, \rho^0_P, \rho^{\pm,3}_\Pi$.

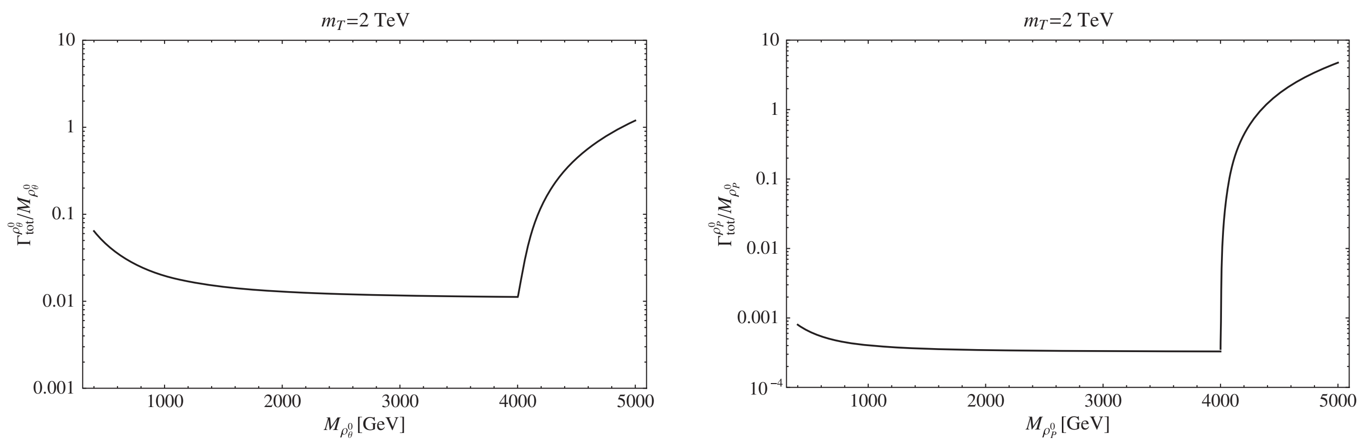

FIG. 2. The total widths (normalized by each mass) of the isosinglet technirho mesons, ρ^0_θ (left panel) and ρ^0_P (right panel), as functions of respective technirho masses. Here we take the masses of the relevant technipions, $M_{T^{3,0,\pm}}$, to be 2 TeV.

section we show several coupling properties of ρ^0_θ, ρ^0_P, and $\rho^{\pm,3}_\Pi$.

Partial decay rates of technirho mesons are calculated by using relevant vertex terms, which are summarized in Appendix A, and we show the explicit expressions of them for ρ^0_θ, ρ^0_P, and $\rho^{\pm,3}_\Pi$ in Appendix B. The total decay widths are plotted in Figs. 2 and 3 as functions of respective technirho masses. The branching fractions for ρ^0_θ, ρ^0_P, and $\rho^{\pm,3}_\Pi$ are shown in Figs. 4 and 5, respectively. It should also be noted [25] that the technipion masses are severely constrained by the LHC data to be of the order of TeV.[6]

Here we take reference values of the masses of technipions relevant to the calculations as $(M_{T^{3,0,\pm}}, M_{P^{\pm,3}}) = (2\ \text{TeV}, 1\ \text{TeV})$. This choice is motivated by the results of Refs. [24,25], in which it was shown that color-triplet technipions are heavier than the color-singlet technipions, and can be as heavy as $O(1)$ TeV. Changing the reference values of $M_{T^{3,0,\pm}}$ and $M_{P^{\pm,3}}$ does not affect the LHC analysis in this paper, as will be shown later.

From these figures, we see the following general tendency: technirhos dominantly decay to $\pi\pi$ or $\pi W_L(\pi Z_L)$ above the thresholds (besides $W_L W_L/W_L Z_L$ channels for $\rho^{\pm,3}_\Pi$), and decay widths become large; on the other hand, below these thresholds, the decay widths are small enough that usual resonance search strategies can be applied. In particular, the ρ^0_θ (ρ^0_P) search below the technipion-pair-decay threshold is interesting in the sense that it dominantly decays to the TD and g (γ/Z). The associate production of the Higgs and the weak gauge boson (W, Z) from the resonant vector meson were discussed in the literature in the context of certain kinds of dynamical EW symmetry-breaking scenarios [33–35]. On the other hand, the associated production of the Higgs (TD) and the gluon from the color-octet vector resonance is a characteristic process of the one-family model, which plays an important and complementary role together with existing studies of color-octet signals (see, e.g., Ref. [36]).

[6]The perturbative treatment of such a large explicit breaking effect might sound questionable. Actually, this is a typical phenomenon of the "amplification of a symmetry violation" by the large anomalous dimension in the dynamics near the criticality, resulting in huge violation effects for a small violation parameter, as was most dramatically shown in the top-quark condensate model [32] (see, e.g., the first reference in Ref. [5]). One also might suspect that since all the massive pseudo-NG bosons are decoupled, leaving only three exact NG bosons absorbed into the W/Z bosons, the theory would be equivalent to the model based on the $G/H = \text{SU}(2)_L \times \text{SU}(2)_R/\text{SU}(2)_V$, i.e., the one-doublet model. However, the fictitious NG bosons (absorbed into W/Z) as well as the TD are composites of all eight flavors of technifermions (not a particular subset of them), which contribute to the dynamics on the same footing. Thus the walking dynamics responsible for the lightness of the TD as well as the FCNC solution is still operative in contrast to the one-doublet model.

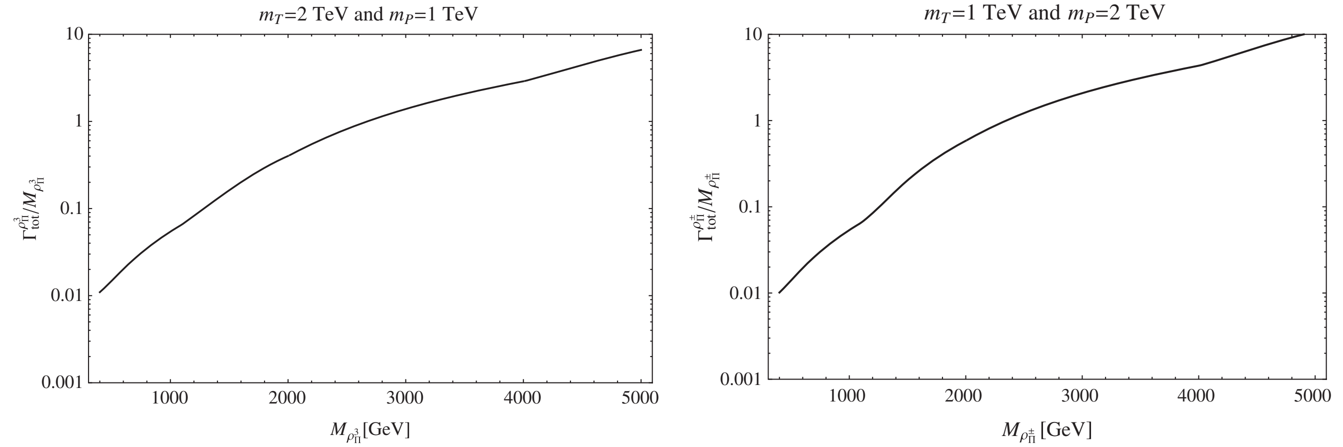

FIG. 3. The total widths (normalized by each mass) of the isotriplet technirho mesons, ρ_Π^3 (left panel) and $\rho_\Pi^\pm$ (right panel), as functions of respective technirho masses. Here we take the masses of the relevant technipions as $(M_{T^{3,0,\pm}}, M_{P^{\pm,3}}) = (2\ \text{TeV}, 1\ \text{TeV})$.

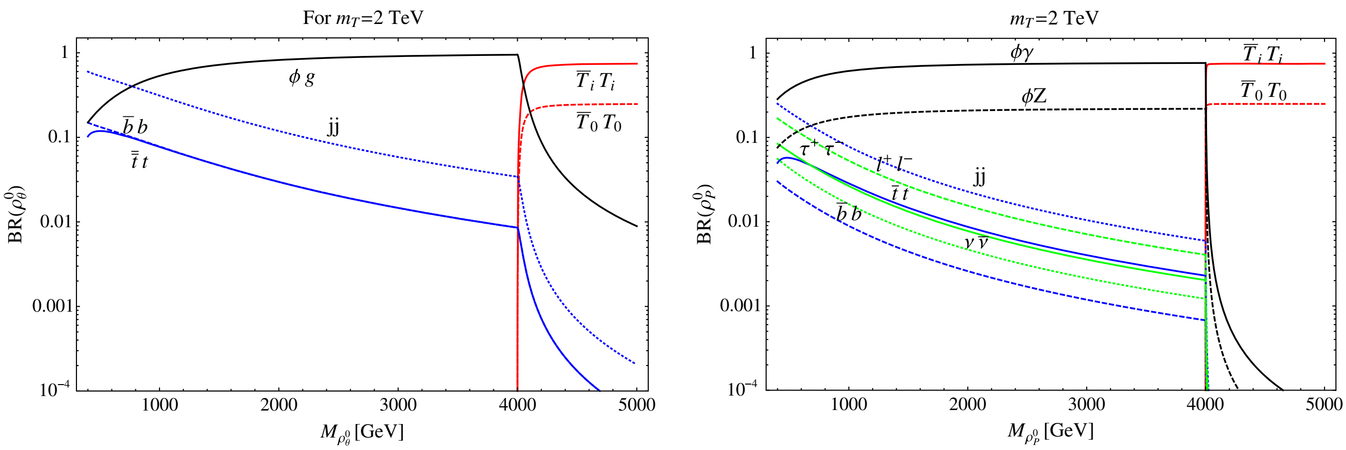

FIG. 4 (color online). The branching ratios of the isosinglet technirho mesons, ρ_θ^0 (left panel) and ρ_P^0 (right panel), as functions of respective technirho masses. Here we take the masses of the relevant technipions, $M_{T^{3,0,\pm}}$, to be 2 TeV. Note that j and l in the figure represent the sum of light quarks $(j = u, d, s, c)$ and that of leptons $l = (e, \mu)$, respectively.

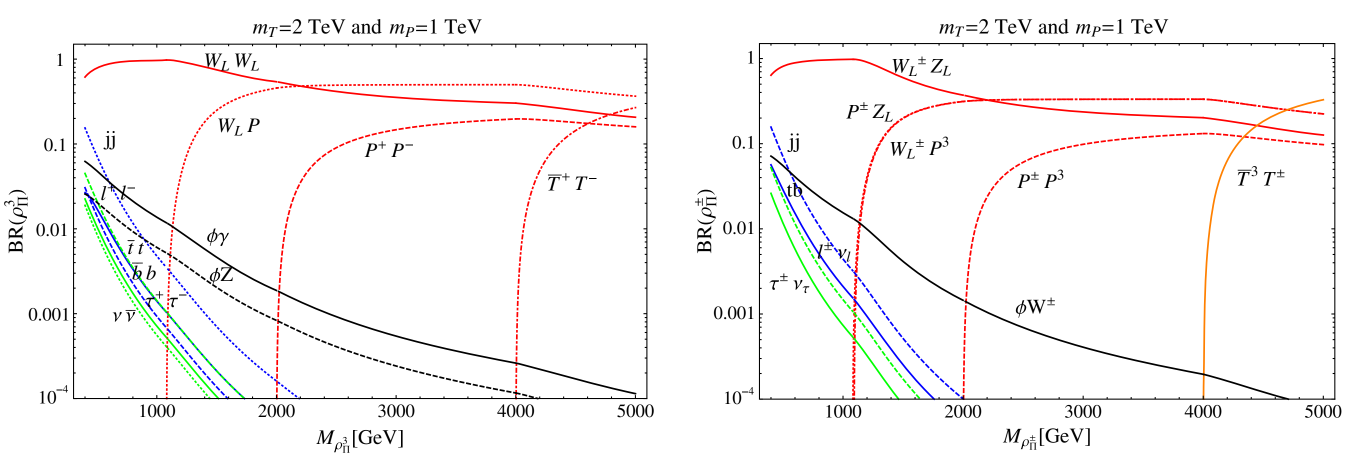

FIG. 5 (color online). The branching ratios of the isotriplet technirho mesons, ρ_Π^3 (left panel) and $\rho_\Pi^\pm$ (right panel), as functions of respective technirho masses. Here we take the masses of the relevant technipions as $(M_{T^{3,0,\pm}}, M_{P^{\pm,3}}) = (2\ \text{TeV}, 1\ \text{TeV})$. Note that j and l in the figure represent the sum of light quarks $(j = u, d, s, c)$ and that of leptons $l = (e, \mu)$, respectively.

IV. TECHNIRHO PHENOMENOLOGY AT THE LHC

As mentioned in the previous section, we consider the production of the technirho mesons through the mixing with the SM gauge bosons, which are produced by the Drell-Yan process (Fig. 1). The LHC cross section of technirho mesons with mass M_ρ is thus calculated to be (for a review, see, e.g., Ref. [37])

$$\frac{d\sigma(pp\to\rho X\to ABX)}{d\eta dM^2} = \frac{32\pi}{s}\sum_{a,b}C_{ab}f_{a/p}\left(\frac{M}{\sqrt{s}}e^{\eta},M^2\right)f_{b/p}\left(\frac{M}{\sqrt{s}}e^{-\eta},M^2\right) \times\frac{C_\rho\cdot(2S_\rho+1)}{(2S_a+1)(2S_b+1)}\frac{M^2}{M_\rho^2}\frac{\Gamma(\rho\to ab)\Gamma(\rho\to AB)}{(M^2-M_\rho^2)^2+M_\rho^2(\Gamma_{\rm tot}^\rho)^2}, \tag{31}$$

where the function $f_{a/p}$ denotes the parton distribution function for parton a in the proton, which is available from Ref. [38] (for CTEQ6M); $\Gamma_{\rm tot}^\rho$ is the total width of ρ; $\sqrt{s}$ is the center-of-mass energy at the LHC ($\sqrt{s}=8$ or 14 TeV); η is the rapidity of the a-b system in the p-p center-of-mass frame; C_{ab} is a multiplication factor regarding the SU(3) gauge group [e.g., $C_{gg}=(1/8)^2, C_{qq}=(1/3)^2$]; M^2 is an invariant mass squared associated with particles A and B coming out of the ρ; $C_\rho=1(8)$ for the color-singlet (-octet) technirho meson; and $(2S_a+1)$ is a multiplication factor for spin degeneracy [e.g., $(2S_a+1)=2$ for $a=q$, and $(2S_\rho+1)=3$ for vector mesons].

Figures 2 and 3 show that $\Gamma_{\rm tot}^\rho/M_\rho \ll 1$ when $M_\rho<4$ TeV for ρ_θ^0,ρ_P^0 and $M_\rho<2$ TeV for $\rho_\Pi^{3,\pm}$, so we may apply the narrow-width approximation when evaluating Eq. (31) by replacing the ρ-resonance function $1/[(M^2-M_\rho^2)^2+M_\rho^2(\Gamma_{\rm tot}^\rho)^2]$ with $\pi/(M_\rho\Gamma_{\rm tot}^\rho)\delta(M^2-M_\rho^2)$. For the Drell-Yan production ($\rho=\rho_\theta^0,\rho_\Pi^{\pm,3},\rho_P^0$), we thus have

$$\sigma_{\rm DY}(pp\to\rho\to AB) = \frac{16\pi^2}{3s}C_\rho\frac{{\rm BR}(\rho\to AB)}{M_\rho}\times\sum_{q={\rm quarks}}\Gamma(\rho\to q\bar{q}) \times\int_{-Y_B}^{Y_B}d\eta f_{q/p}\left(\frac{M_\rho}{\sqrt{s}}e^{\eta},M_\rho^2\right)f_{\bar{q}/p}\left(\frac{M_\rho}{\sqrt{s}}e^{-\eta},M_\rho^2\right), \tag{32}$$

where $Y_B=-\frac{1}{2}\ln(M_\rho^2/s)$. The predicted production cross sections of ρ_θ^0, ρ_P^0, and $\rho_\Pi^{3,\pm}$ are plotted in Figs. 6 and 7, respectively.

A. Current LHC limits

The 8 TeV LHC data have already placed stringent constraints on masses of hypothetical heavy resonances, such as Z', W', colorons, Kaluza-Klein (KK) gluons, etc. Here, by using these results, we give a rough estimate of the lower bound on the masses of technirho mesons (ρ_θ^0,ρ_P^0, and $\rho_\Pi^{\pm,3}$) in the one-family model under the assumption that the kinematics of the technirho production and decay process is more or less the same as that of, e.g., Z' and W'.

The resonance search in the dijet mass distribution by CMS [39] places the strongest constraint on the ρ_θ^0 mass, while the dilepton resonance search by the ATLAS [40] and CMS [41] experiments are the most relevant for ρ_P^0 and ρ_Π^3. As for $\rho_\Pi^\pm$, studies of resonant $WZ\to 3\ell+\nu$ production by the ATLAS [42] and CMS [43] experiments place the strongest constraint on its mass. The $t\bar{t}$ resonance search by ATLAS [42] and CMS [43] also places a strong constraint on the ρ_θ^0 mass, though the bound is slightly weaker than

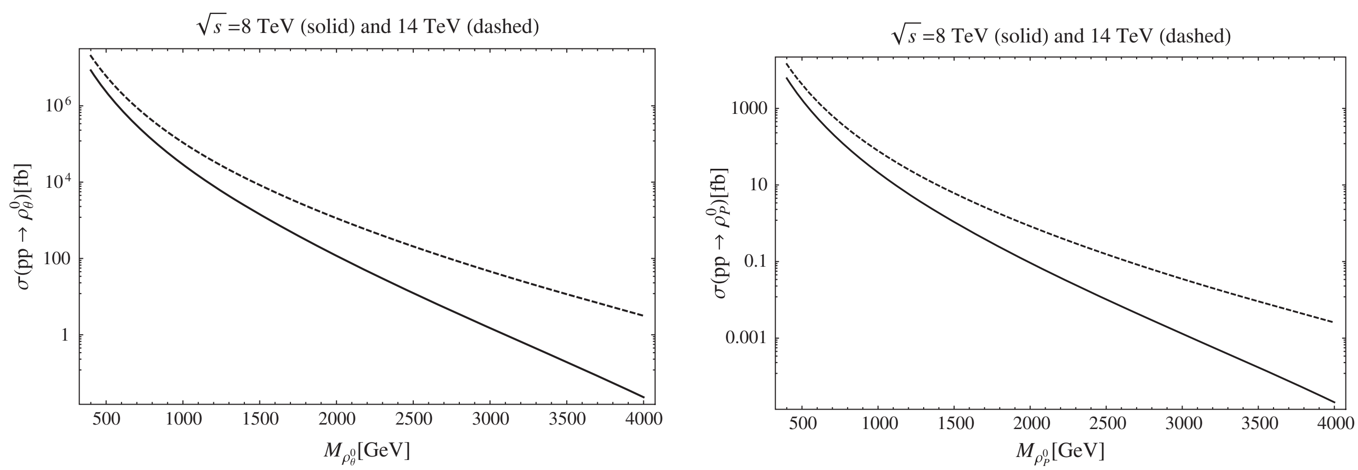

FIG. 6. The LHC production cross sections of the isosinglet technirho mesons, ρ_θ^0 (left panel) and ρ_P^0 (right panel), as functions of respective technirho masses. Here we take the masses of the relevant technipions as $(M_{T^{3,0,\pm}},M_{P^{\pm,3}})=(2$ TeV, 1 TeV$)$. The solid and dashed curves correspond to $\sqrt{s}=8$ and 14 TeV, respectively.

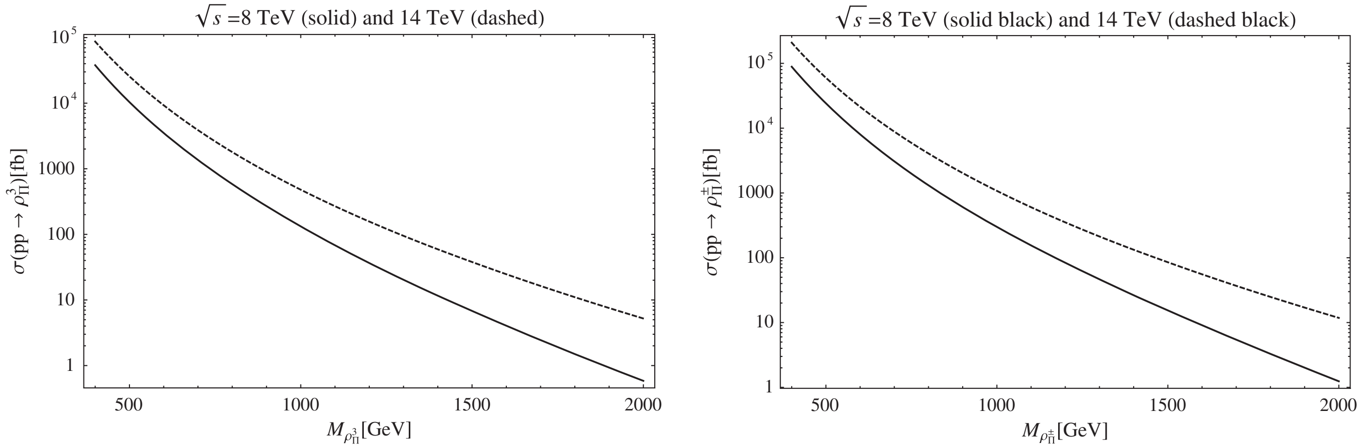

FIG. 7. The LHC production cross sections of the isotriplet technirho mesons, ρ_Π^3 (left panel) and $\rho_\Pi^\pm$ (right panel), as functions of respective technirho masses. Here we take the mass of the relevant technipion, $M_{T^{3,0,\pm}}$, to be 2 TeV. The solid and dashed curves correspond to $\sqrt{s} = 8$ and 14 TeV, respectively.

that obtained from the dijet search. ρ_P^0 and ρ_Π^3 also decay to $\bar{t}t$, though the cross sections are so small that the current LHC data do not give any constraint.

In Figs. 8 and 9, we show the cross sections for isosinglet rho mesons (ρ_θ^0, ρ_P^0) and isotriplet rho mesons ($\rho_\Pi^3, \rho_\Pi^\pm$), respectively, for the most constrained decay processes mentioned above together with the experimental upper bounds. Here, we take the technipion masses as $M_{T^{\pm,3}} =$ 2 TeV and $M_{P^{\pm,3}} = 1$ TeV, so that all the relevant energy regions are below the threshold of the decay channels that involve technipion(s). From these figures, we find that the current LHC experiments constrain the masses of the technirho mesons to be

$$M_{\rho_\theta^0} \gtrsim 1.7 \text{ TeV}, \qquad M_{\rho_P^0} \gtrsim 1.3 \text{ TeV},$$
$$M_{\rho_\Pi^3} \gtrsim 1.0 \text{ TeV}, \qquad M_{\rho_\Pi^\pm} \gtrsim 1.4 \text{ TeV}. \qquad (33)$$

As we mentioned in the previous section, these results are insensitive to the precise values of $M_{T^{\pm,3}}$ and $M_{P^{\pm,3}}$ as long as the relevant mass range of the technirho is below the technipion thresholds.

Note that the limits on the technirho masses are milder than those on other hypothetical spin-1 resonances, such as W', Z', KK gluons, and colorons. This is due to the fact that the technirho mesons have no direct couplings to the SM quarks, and hence the Drell-Yan productions necessarily

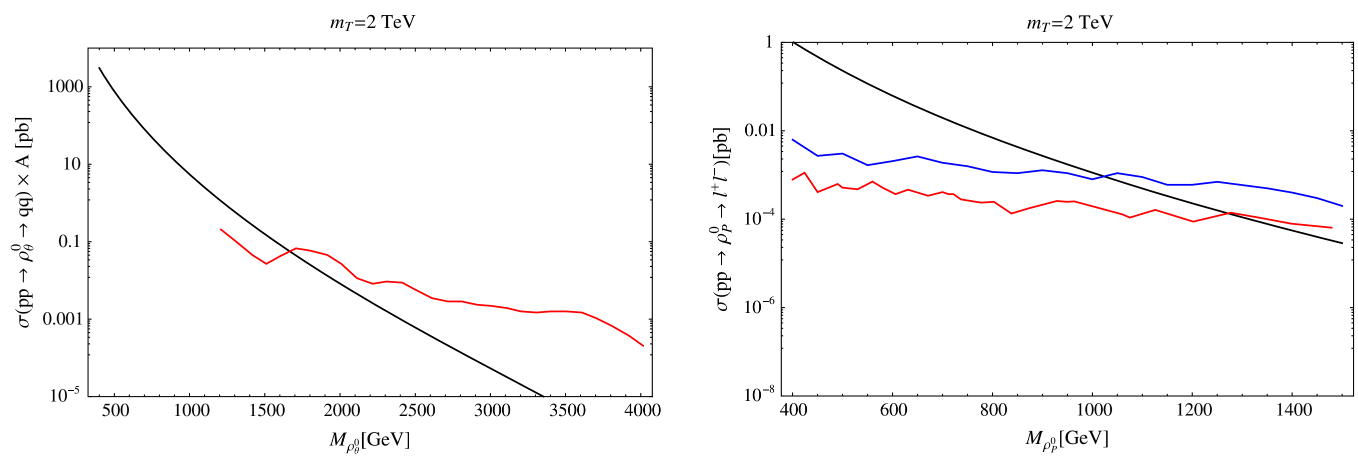

FIG. 8 (color online). Left panel: $\sigma_{\rm DY}(pp \to \rho_\theta^0 \to qq)$, multiplied by the acceptance $A \simeq 0.6$ [39,44], for $q = u, d, s, c$, in units of pb as a function of $M_{\rho_\theta^0}$. The black curve (smooth decreasing function) corresponds to the prediction of the one-family model for $\sqrt{s} = 8$ TeV. The 95% C.L. upper limit on a generic cross section times the acceptance set by searches for a new resonance in the dijet mass distribution by the CMS experiment with $\sqrt{s} = 8$ TeV [39] is shown by the red curve (nonsmooth line). Right panel: $\sigma_{\rm DY}(pp \to \rho_P^0 \to l^+l^-)$ with $l = e, \mu$ as a function of $M_{\rho_P^0}$. The black curve (smooth decreasing function) corresponds to the prediction of the one-family model for $\sqrt{s} = 8$ TeV. The 95% C.L. upper limits on the $Z' \to l^+l^-$ cross section reported by the ATLAS [40] and CMS [41] experiments are shown by the blue and red curves (upper and lower of two nonsmooth lines), respectively.

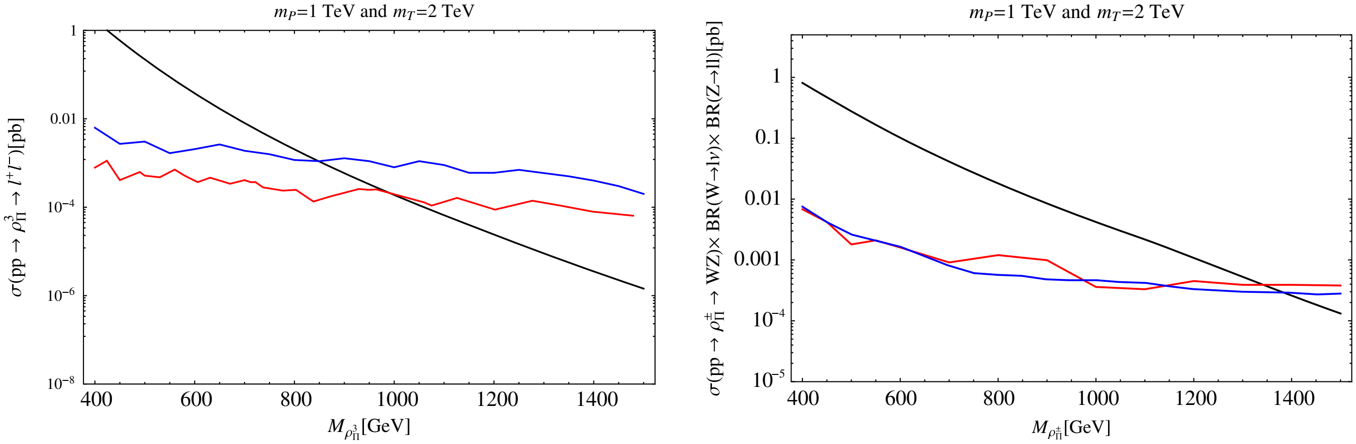

FIG. 9 (color online). Left panel: $\sigma_{\rm DY}(pp \to \rho_\Pi^3 \to l^+l^-)$ with $l = e, \mu$ as a function of $M_{\rho_\Pi^3}$. The black curve (smooth decreasing function) corresponds to the prediction of the one-family model for $\sqrt{s} = 8$ TeV. The 95% C.L. upper limits on the $Z' \to l^+l^-$ cross section reported by the ATLAS [40] and CMS [41] experiments are shown by the blue and red curves (upper and lower of two nonsmooth lines), respectively. Right panel: $\sigma_{\rm DY}(pp \to \rho_\Pi^\pm \to W^\pm Z) \times {\rm BR}(W^\pm \to l^\pm \nu) \times {\rm BR}(Z \to l^+l^-)$ with $l = e, \mu$. The black curve (smooth decreasing function) corresponds to the prediction of the one-family model for $\sqrt{s} = 8$ TeV. The 95% C.L. upper limits on the $W'/\rho \to WZ$ cross section reported by the ATLAS [42] and CMS [43] experiments are shown by the blue and red curves (middle and upper curves at the right edge of the plot), respectively.

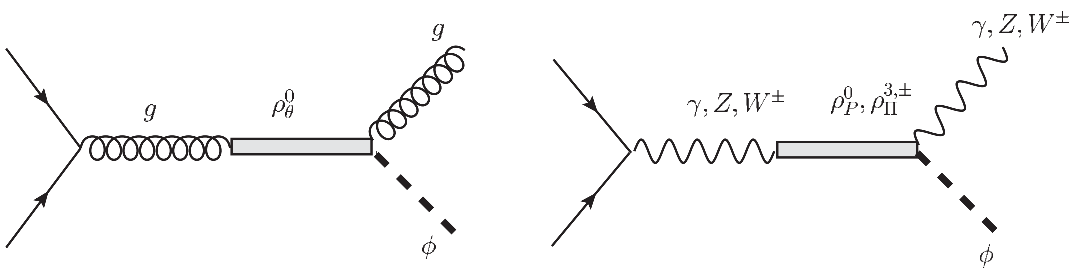

FIG. 10. DY production of the technirho meson which decays into the Higgs (TD) and the SM gauge boson.

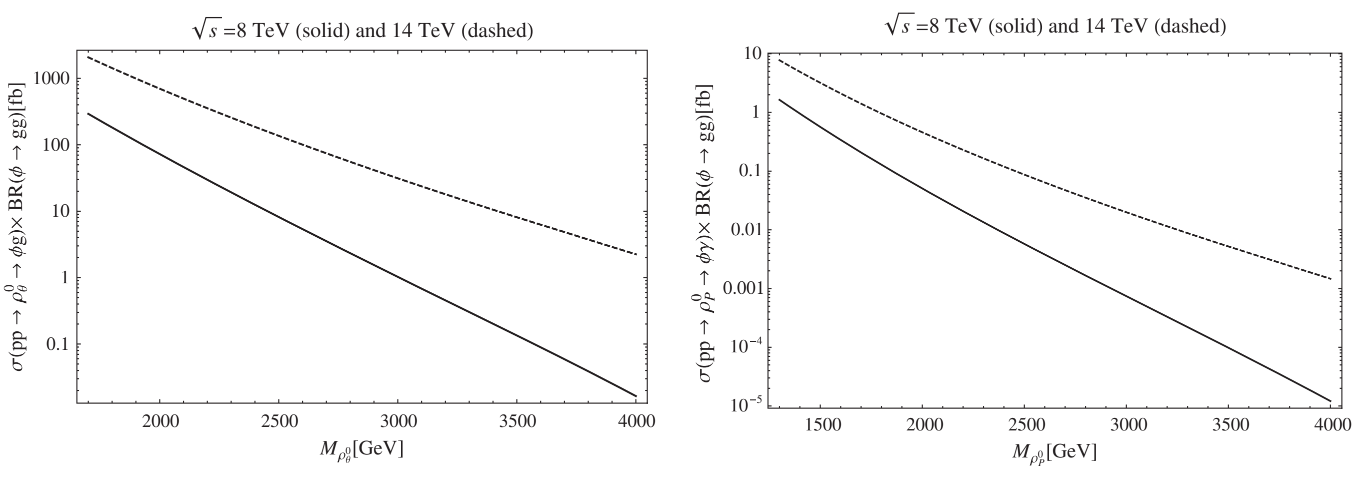

FIG. 11. Left panel: $\sigma_{\rm DY}(pp \to \rho_\theta^0 \to \phi g) \times {\rm BR}(\phi \to gg)$ in units of fb as a function of $M_{\rho_\theta^0}$. Right panel: $\sigma_{\rm DY}(pp \to \rho_P^0 \to \phi\gamma) \times {\rm BR}(\phi \to gg)$ in units of fb as a function of $M_{\rho_P^0}$. In both plots, the solid and dashed curves correspond to cross sections for $\sqrt{s} = 8$ and 14 TeV, respectively. Also, the branching ratio of the TD decaying into two gluons is taken to be ${\rm BR}(\phi \to gg) = 75\%$ [10].

arise through the $\rho - \mathcal{V}$ mixing, as in Eq. (A2), leading to the significant suppression by α_s or $\alpha_{\rm em}$ in the amplitudes compared to the case for other hypothetical spin-1 resonances.

B. Associated production of the technidilaton and the SM gauge boson through a resonant technirho

As we mentioned at the end of Sec. III, the most interesting search channel of the one-family model is the DY production of the technirho meson which decays into the Higgs (TD) and the SM gauge boson (see Fig. 10). The processes consist of vertices in Eqs. (25) and (26) as a consequence of the scale-invariant extension of the HLS formalism.

Since the TD dominantly decays to two gluons [10], we consider the process where the produced TD subsequently decays into two gluons. In Fig. 11, we plot the cross sections of $pp \to \rho^0_\theta \to \phi g \to ggg$ (left panel) and $pp \to \rho^0_P \to \phi\gamma \to gg\gamma$ (right panel) as functions of each technirho mass. Here, we take ${\rm BR}(\phi \to gg) = 75\%$, which can be read off from Ref. [10]. The mass ranges (horizontal axes) of the plots are chosen in such a way that they are above the current LHC limit derived in the previous subsection and below the threshold of decay channels which involve the technipion(s). We can see that the cross section for the color-singlet channel (right panel in Fig. 11) is rather small, and it may be challenging even at the 14 TeV LHC. We also estimated similar cross sections for an isotriplet technirho production followed by its decay into the TD and the electroweak gauge boson ($W/Z/\gamma$), and found that these cross sections are even smaller than that of the ρ^0_P case. Meanwhile, the color-octet channel (left panel in Fig. 11) has large cross sections, and can be a promising search channel at the 14 TeV LHC. A detailed collider study of this channel will be published elsewhere [45].

V. SUMMARY

In this paper, we formulated a scale-invariant hidden local symmetry as a low-energy effective theory of walking technicolor, which includes the technidilaton, technipions, and technirho mesons as the low-lying spectra. As a benchmark for LHC phenomenology, our discussions have in particular focused on the one-family model of walking technicolor with eight technifermion flavors, which can be—at energy scales relevant to the reach of the LHC—described by the scale-invariant hidden local symmetry based on the manifold $[{\rm SU}(8)_L \times {\rm SU}(8)_R]_{\rm global} \times {\rm SU}(8)_{\rm local}/{\rm SU}(8)_V$, where ${\rm SU}(8)_{\rm local}$ is the hidden local symmetry and the global ${\rm SU}(8)_L \times {\rm SU}(8)_R$ symmetry is partially gauged by the ${\rm SU}(3) \times {\rm SU}(2)_L \times U(1)_Y$ of the SM. Based on the scale-invariant hidden local symmetry, we evaluated the coupling properties of the technirho mesons and placed limits on the masses from the current LHC data. Then, implications for future LHC phenomenology were discussed by focusing on the technirho mesons produced through the Drell-Yan process. We found that the color-octet technirho decaying to the technidilaton along with the gluon is of interest as a discovery channel at the LHC, which would provide a characteristic signature to probe the one-family model of walking technicolor. More detailed collider studies are in progress.

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APPENDIX A: INTERACTIONS

In this appendix, we summarize interactions that involve one technirho meson, which can be obtained by expanding the $\chi^2 F_\sigma^2 {\rm tr}[\hat{\alpha}^2_{\mu\|}]$ term in Eq. (3):

$$\chi^2 F_\sigma^2 {\rm tr}[\hat{\alpha}^2_{\mu\|}] = F_\sigma^2\left(1 + \frac{2\phi}{F_\phi} + \cdots\right) \times {\rm tr}\left[(\mathcal{V}_\mu - V_\mu)^2 + \frac{i}{F_\pi^2} V_\mu[\partial^\mu \pi, \pi] + \frac{2i}{F_\pi} V_\mu[\mathcal{A}^\mu, \pi] + \cdots\right]. \tag{A1}$$

Here, we focus on a set of spectra, $(\rho^0_\theta, \rho^{\pm,3}_\Pi, \rho^0_P)$, which are expected to be produced through the Drell-Yan processes at the LHC. Using Eqs. (7), (9), and (15), we thus derive the technirho couplings relevant for the LHC phenomenology.

1. ρ-$\mathcal{V}$ mixing terms

$$\begin{aligned}\mathcal{L}_{\mathcal{V}\rho} &= -2gF_\sigma^2 {\rm tr}[\mathcal{V}_\mu \rho^\mu] \\ &= -2gF_\sigma^2\Bigg[\frac{g_s}{\sqrt{2}} G^a_\mu \rho^{0a\mu}_\theta + eA_\mu\left\{\rho^{3\mu}_\Pi + \frac{1}{\sqrt{3}}\rho^{0\mu}_P\right\} \\ &\quad + \frac{e}{2sc} Z_\mu\left\{(c^2 - s^2)\rho^{3\mu}_\Pi - \frac{2}{\sqrt{3}} s^2 \rho^{0\mu}_P\right\} \\ &\quad + \frac{e}{2s}\{W^-_\mu \rho^{\mu+}_\Pi + {\rm H.c.}\}\Bigg],\end{aligned} \tag{A2}$$

where the charged rho-meson fields have been defined as

$$\rho^{\mu\pm}_\Pi = \frac{\rho^{1\mu}_\Pi \mp i\rho^{2\mu}_\Pi}{\sqrt{2}}. \tag{A3}$$

Note the absence of $A - \rho^3_P, Z - \rho^3_P$, and $W^\pm - \rho^\mp_P$ terms due to a coincident cancellation of contributions between the techniquark and lepton sectors, which follows from the orthogonality of the ${\rm SU}(8)_V$ generators. These terms are crucial for technirho meson productions through the Drell-Yan process at the LHC and allow the decays to the SM fermions by assuming vector-meson dominance.

2. ρ-f-f couplings

The ρ-$\mathcal{V}$ terms in Eq. (A2) allow the decays to the SM fermions by assuming vector-meson dominance via the couplings induced by the SM gauge-boson exchanges, evaluated at the ρ on shell:

$$\mathcal{L}_{\rho-f-f} = -2gF_\sigma^2\left[\frac{g_s^2}{\sqrt{2}}\frac{1}{M^2_{\rho^0_\theta}}\bar{q}\gamma_\mu\rho^{\mu 0}_{\theta a}\left(\frac{\lambda_a}{2}\right)q + e\frac{1}{M^2_{\rho^3_\Pi}}J^{\rm em}_\mu\rho^{3\mu}_\Pi + \frac{e(c^2-s^2)}{2sc}\frac{1}{M^2_{\rho^3_\Pi}-m_Z^2}J^Z_\mu\rho^{3\mu}_\Pi \right.$$
$$\left. + \frac{e}{2s}\frac{1}{M^2_{\rho^\pm_\Pi}-m_W^2}(J^{W^+}_\mu\rho^{+\mu}_\Pi + \text{H.c.}) + \frac{e}{\sqrt{3}}\frac{1}{M^2_{\rho^0_P}}J^{\rm em}_\mu\rho^{0\mu}_P - \frac{es}{\sqrt{3}c}\frac{1}{M^2_{\rho^0_P}-m_Z^2}J^Z_\mu\rho^{0\mu}_P\right], \tag{A4}$$

where

$$J^{\rm em}_\mu = e\sum_f \bar{f}\gamma_\mu Q^f_{\rm em}f, \tag{A5}$$

$$J^Z_\mu = \frac{e}{sc}\sum_f[\bar{f}_L\gamma_\mu(\tau^f_3 - s^2Q^f_{\rm em})f_L + \bar{f}_R\gamma_\mu(-s^2Q^f_{\rm em})f_R], \tag{A6}$$

$$J^{W^\pm}_\mu = \frac{e}{\sqrt{2}s}\sum_f \bar{f}_L\gamma_\mu\tau^f_\pm f_L. \tag{A7}$$

3. ρ-$\mathcal{V}$-ϕ terms

$$\mathcal{L}_{\mathcal{V}\rho\phi} = -\frac{4gF_\sigma^2}{F_\phi}\phi\,\text{tr}[\mathcal{V}_\mu\rho^\mu]$$
$$= -\frac{4gF_\sigma^2}{F_\phi}\phi\left[\frac{g_s}{\sqrt{2}}G^a_\mu\rho^{0a\mu}_\theta + eA_\mu\left\{\rho^{3\mu}_\Pi + \frac{1}{\sqrt{3}}\rho^{0\mu}_P\right\}\right.$$
$$\left. + \frac{e}{2sc}Z_\mu\left\{(c^2-s^2)\rho^{3\mu}_\Pi - \frac{2}{\sqrt{3}}s^2\rho^{0\mu}_P\right\} + \frac{e}{2s}\{W^-_\mu\rho^{\mu+}_\Pi + \text{H.c.}\}\right]. \tag{A8}$$

Note again the absence of $A-\rho^3_P-\phi, Z-\rho^3_P-\phi$, and $W^\pm-\rho^\mp_P-\phi$ terms due to the orthogonality, namely, the cancellation of contributions between the techniquark and lepton sectors.

4. ρ-π-π terms

The ρ-π-π terms are decomposed into four parts:

$$\mathcal{L}_{\rho-\pi-\pi} = 2ig_{\rho\pi\pi}\text{tr}[\rho_\mu[\partial^\mu\pi,\pi]] = \frac{igF_\sigma^2}{F_\pi^2}\text{tr}[\rho_\mu[\partial^\mu\pi,\pi]]$$
$$= \frac{igF_\sigma^2}{F_\pi^2}\text{tr}[\rho^\mu_{QQ}(\overleftrightarrow{\partial}_\mu\pi_{QQ}\pi_{QQ} + \overleftrightarrow{\partial}_\mu\pi_{QL}\pi_{LQ}) + \rho^\mu_{QL}(\overleftrightarrow{\partial}_\mu\pi_{LQ}\pi_{QQ} + \overleftrightarrow{\partial}_\mu\pi_{LL}\pi_{LQ})$$
$$+ \rho^\mu_{LQ}(\overleftrightarrow{\partial}_\mu\pi_{QQ}\pi_{QL} + \overleftrightarrow{\partial}_\mu\pi_{QL}\pi_{LL}) + \rho^\mu_{LL}(\overleftrightarrow{\partial}_\mu\pi_{LQ}\pi_{QL} + \overleftrightarrow{\partial}_\mu\pi_{LL}\pi_{LL})]$$
$$= \mathcal{L}_{\rho-\pi-\pi} + \mathcal{L}_{\rho_P-\pi-\pi} + \mathcal{L}_{\rho_T-\pi-\pi} + \mathcal{L}_{\rho_\theta-\pi-\pi}, \tag{A9}$$

where we have defined, for arbitrary fields $\bar{A}$ and B,

$$\overleftrightarrow{\partial}_\mu AB \equiv \partial_\mu AB - A\partial_\mu B. \tag{A10}$$

The ρ_Π-π-π terms (arising from the ρ_{QQ} and ρ_{LL} terms) are

$$\mathcal{L}_{\rho_\Pi-\pi-\pi} = \frac{igF_\sigma^2}{F_\pi^2}\left[\frac{1}{4}i\epsilon^{ijk}\partial_\mu\Pi^i\Pi^j\rho^{k\mu}_\Pi + \frac{1}{4}i\epsilon^{ijk}\overleftrightarrow{\partial}_\mu\Pi^iP^j\rho^{k\mu}_\Pi + \frac{1}{4}i\epsilon^{ijk}\partial_\mu P^iP^j\rho^{k\mu}_\Pi + \frac{1}{4}i\epsilon^{ijk}\overleftrightarrow{\partial}_\mu\bar{T}^iT^j\rho^{k\mu}_\Pi\right]. \tag{A11}$$

Note the absence of $\rho_\Pi^i - T^i - T^0$ terms due to the accidental cancellation between the techniquark and technilepton sectors. The ρ_P-π-π terms (arising from the ρ_{QQ} and ρ_{LL} terms) are

$$\mathcal{L}_{\rho_P-\pi-\pi} = \frac{igF_\sigma^2}{F_\pi^2}\bigg[-\frac{1}{2\sqrt{3}}i\epsilon^{ijk}\partial_\mu P^i P^j \rho_P^{k\mu} - \frac{1}{4\sqrt{3}}i\epsilon^{ijk}\overset{\leftrightarrow}{\partial}_\mu \bar{T}^i T^j \rho_P^{k\mu} - \frac{1}{2\sqrt{3}}(\overset{\leftrightarrow}{\partial}_\mu \bar{T}^i T^0 + \overset{\leftrightarrow}{\partial}_\mu \bar{T}^0 T^i)\rho_P^{i\mu} - \frac{1}{2\sqrt{3}}(\overset{\leftrightarrow}{\partial}_\mu \bar{T}^i T^i + \overset{\leftrightarrow}{\partial}_\mu \bar{T}^0 T^0)\rho_P^{0\mu}\bigg]. \tag{A12}$$

Note that there is no ρ_P-Π-Π term due to an accidental cancellation between the techniquark and technilepton sectors.

The ρ_T-π-π terms (arising from the ρ_{QL} and ρ_{LQ} terms) are expressed as

$$\begin{aligned}\mathcal{L}_{\rho_T-\pi-\pi} = \frac{igF_\sigma^2}{F_\pi^2}\bigg[&\frac{i}{2\sqrt{2}}\epsilon^{ijk}\overset{\leftrightarrow}{\partial}_\mu \bar{T}^i \theta_a^j\left(\frac{\lambda_a}{2}\right)\rho_T^{k\mu} + \frac{1}{2\sqrt{2}}\overset{\leftrightarrow}{\partial}_\mu \bar{T}^0 \theta_a^i\left(\frac{\lambda_a}{2}\right)\rho_T^{i\mu} + \frac{1}{2\sqrt{2}}\overset{\leftrightarrow}{\partial}_\mu \bar{T}^i \theta_a^0\left(\frac{\lambda_a}{2}\right)\rho_T^{i\mu}\\ &+ \frac{i}{2\sqrt{3}}\epsilon^{ijk}\overset{\leftrightarrow}{\partial}_\mu \bar{T}^i P^j \rho_T^{k\mu} + \frac{1}{2\sqrt{3}}\overset{\leftrightarrow}{\partial}_\mu \bar{T}^0 P^i \rho_T^{i\mu} + \frac{1}{2\sqrt{3}}\overset{\leftrightarrow}{\partial}_\mu \bar{T}^i P^0 \rho_T^{i\mu} + \frac{1}{2\sqrt{2}}\overset{\leftrightarrow}{\partial}_\mu \bar{T}^i \theta_a^i\left(\frac{\lambda_a}{2}\right)\rho_T^{0\mu}\\ &+ \frac{1}{2\sqrt{2}}\overset{\leftrightarrow}{\partial}_\mu \bar{T}^0 \theta_a^0\left(\frac{\lambda_a}{2}\right)\rho_T^{0\mu} + \frac{1}{2\sqrt{3}}\overset{\leftrightarrow}{\partial}_\mu \bar{T}^i P^i \rho_T^{0\mu} + \frac{1}{2\sqrt{3}}\overset{\leftrightarrow}{\partial}_\mu \bar{T}^0 P^0 \rho_T^{0\mu}\bigg] + \text{H.c.}.\end{aligned} \tag{A13}$$

The ρ_θ-π-π terms (arising from the ρ_{QQ} term) are expressed as

$$\begin{aligned}\mathcal{L}_{\rho_\theta-\pi-\pi} = \frac{igF_\sigma^2}{F_\pi^2}\bigg[&-\frac{1}{\sqrt{2}}f^{abc}\epsilon^{ijk}\partial_\mu \theta_a^i \theta_b^j \rho_{\theta c}^{k\mu} - \frac{1}{2\sqrt{2}}\epsilon^{ijk}\overset{\leftrightarrow}{\partial}_\mu \bar{T}^j\left(\frac{\lambda_a}{2}\right)T^i\rho_{\theta a}^{k\mu} - \frac{1}{2\sqrt{2}}\overset{\leftrightarrow}{\partial}_\mu \bar{T}^i\left(\frac{\lambda_a}{2}\right)T^0\rho_{\theta a}^{i\mu} - \frac{1}{2\sqrt{2}}\overset{\leftrightarrow}{\partial}_\mu \bar{T}^0\left(\frac{\lambda_a}{2}\right)T^i\rho_{\theta a}^{i\mu}\\ &- \frac{1}{2\sqrt{2}}\overset{\leftrightarrow}{\partial}_\mu \bar{T}^i\left(\frac{\lambda_a}{2}\right)T^i\rho_{\theta a}^{0\mu} - \frac{1}{2\sqrt{2}}\overset{\leftrightarrow}{\partial}_\mu \bar{T}^0\left(\frac{\lambda_a}{2}\right)T^0\rho_{\theta a}^{0\mu}\bigg].\end{aligned} \tag{A14}$$

5. ρ-$\mathcal{A}$-π terms

The ρ-$\mathcal{A}$-π terms are constructed from three parts:

$$\begin{aligned}\mathcal{L}_{\rho-\mathcal{A}-\pi} &= \frac{2igF_\sigma^2}{F_\pi}\,\mathrm{tr}[\rho_\mu[\mathcal{A}^\mu,\pi]]\\ &= \mathcal{L}_{\rho-\mathcal{A}-\pi} + \mathcal{L}_{\rho_P-\mathcal{A}-\pi} + \mathcal{L}_{\rho_T-\mathcal{A}-\pi}.\end{aligned} \tag{A15}$$

The ρ_Π-$\mathcal{A}$-π terms are

$$\mathcal{L}_{\rho_\Pi-\mathcal{A}-\pi} = \frac{2igF_\sigma^2}{F_\pi}\left[\frac{e}{4sc}Z_\mu(\rho_\Pi^{+\mu}\Pi^- - \rho_\Pi^{-\mu}\Pi^+) + \frac{e}{4s}[W_\mu^+(\rho_\Pi^{-\mu}\Pi^3 - \rho_\Pi^{3\mu}\Pi^-) + \text{H.c.}].\right] \tag{A16}$$

The ρ_P-$\mathcal{A}$-π terms are

$$\mathcal{L}_{\rho_P-\mathcal{A}-\pi} = \frac{2igF_\sigma^2}{F_\pi}\left[\frac{e}{4sc}Z_\mu(\rho_P^{+\mu}P^- - \rho_P^{-\mu}P^+) + \frac{e}{4s}\{W_\mu^+(\rho_P^{-\mu}P^3 - \rho_P^{+\mu}P^3) + W_\mu^-(\rho_P^{3\mu}P^+ - \rho_P^{+\mu}P^-)\}\right], \tag{A17}$$

where

$$\rho_P^{\pm\mu} \equiv \frac{\rho_P^{1\mu} \mp i\rho_P^{2\mu}}{\sqrt{2}}. \tag{A18}$$

The ρ_T-$\mathcal{A}$-π terms are

$$\mathcal{L}_{\rho_T-\mathcal{A}-\pi} = \frac{2igF_\sigma^2}{F_\pi}\left[\frac{e}{4sc}Z_\mu(\bar{T}^-\rho_T^{+\mu} - \bar{T}^+\rho_T^{-\mu}) + \frac{e}{4s}\{W_\mu^+(\bar{T}^3\rho_T^{-\mu} - \bar{T}^-\rho_T^{3\mu}) + W_\mu^-(\bar{T}^+\rho_T^{3\mu} - \bar{T}^3\rho_T^{+\mu})\}\right] + \text{H.c.}, \tag{A19}$$

where

$$T^{\pm} \equiv \frac{T^1 \mp iT^2}{\sqrt{2}}, \qquad \bar{T}^{\pm} = (T^{\pm})^{\dagger} \tag{A20}$$

$$\rho_T^{\pm\mu} \equiv \frac{\rho_T^{1\mu} \mp i\rho_T^{2\mu}}{\sqrt{2}}. \tag{A21}$$

6. ρ-$\mathcal{A}$-$\mathcal{V}$ terms

It turns out that all the terms involving the $\rho_\theta^0, \rho_\Pi^{\pm,3}$, and ρ_P^0 fields vanish due to the $SU(8)_V$ symmetry, so that

$$\mathcal{L}_{V\mathcal{A}\mathcal{V}}|_{\rho_\theta^0, \rho_\Pi^{\pm,3}, \rho_P^0} = 0. \tag{A22}$$

APPENDIX B: PARTIAL DECAY WIDTHS OF THE TECHNIRHO MESONS

In this appendix, we summarize the partial decay rates of the technirho mesons studied in this paper (ρ_θ^0, ρ_P^0, and $\rho_\Pi^{\pm,3}$) that are relevant for collider phenomenology.

1. The ρ_θ^0 partial decay rates

$$\Gamma(\rho_\theta^0 \to \bar{T}^0 T^0) = \frac{1}{96\pi}\left(\frac{gF_\sigma^2}{2\sqrt{2}F_\pi^2}\right)^2 \frac{[M_{\rho_\theta^0}^2 - 4M_{T^0}^2]^{3/2}}{M_{\rho_\theta^0}^2}, \tag{B1}$$

$$\Gamma(\rho_\theta^0 \to \bar{q}q) = \frac{1}{24\pi}\left(\frac{\sqrt{2}g_s^2 gF_\sigma^2}{M_{\rho_\theta^0}^2}\right)^2 \left(\frac{M_{\rho_\theta^0}^2 + 2m_q^2}{M_{\rho_\theta^0}^2}\right)[M_{\rho_\theta^0}^2 - 4m_q^2]^{1/2}, \tag{B2}$$

$$\Gamma(\rho_\theta^0 \to g\phi) = \frac{1}{8\pi}\left(\frac{2\sqrt{2}g_s gF_\sigma^2}{F_\phi}\right)^2 \left(\frac{M_{\rho_\theta^0}^2 - M_\phi^2}{M_{\rho_\theta^0}^3}\right). \tag{B3}$$

2. The ρ_P^0 partial decay rates

$$\Gamma(\rho_P^0 \to \bar{T}^i T^i) = \frac{3}{16\pi}\left(\frac{gF_\sigma^2}{2\sqrt{3}F_\pi^2}\right)^2 \frac{[M_{\rho_P^0}^2 - 4M_{T^i}^2]^{3/2}}{M_{\rho_P^0}^2}, \tag{B4}$$

$$\Gamma(\rho_P^0 \to \bar{T}^0 T^0) = \frac{1}{16\pi}\left(\frac{gF_\sigma^2}{2\sqrt{3}F_\pi^2}\right)^2 \frac{[M_{\rho_P^0}^2 - 4M_{T^0}^2]^{3/2}}{M_{\rho_P^0}^2}, \tag{B5}$$

$$\Gamma(\rho_P^0 \to f\bar{f}) = \frac{N_c^{(f)}}{12\pi}\left(\frac{2e^2 gF_\sigma^2}{\sqrt{3}}\right)^2 \left[\frac{[G_V^{\rho_P^0}]^2(M_{\rho_P^0}^2 + 2m_f^2) + [G_A^{\rho_P^0}]^2(M_{\rho_P^0}^2 - 4m_f^2)}{M_{\rho_P^0}^2}\right]\sqrt{M_{\rho_P^0}^2 - 4m_f^2}, \tag{B6}$$

$$\Gamma(\rho_P^0 \to \phi\gamma) = \frac{1}{16\pi}\left(\frac{4egF_\sigma^2}{\sqrt{3}F_\phi}\right)^2 \left(\frac{M_{\rho_P^0}^2 - M_\phi^2}{M_{\rho_P^0}^3}\right), \tag{B7}$$

$$\Gamma(\rho_P^0 \to \phi Z) = \frac{1}{16\pi}\left(\frac{4esgF_\sigma^2}{\sqrt{3}cF_\phi}\right)^2 \frac{\sqrt{(M_{\rho_P^0}^2 - (M_\phi + m_Z)^2)(M_{\rho_P^0}^2 - (M_\phi - m_Z)^2)}}{M_{\rho_P^0}^3}, \tag{B8}$$

where $N_c^{(f)} = 1(3)$ for leptons (quarks) and

$$G_V^{\rho_P^0} = \frac{Q_{\rm em}^f}{M_{\rho_P^0}^2} - \frac{\tau_3^f - 2s^2 Q_{\rm em}^f}{2c^2(M_{\rho_P^0}^2 - m_Z^2)}, \tag{B9}$$

$$G_A^{\rho_P^0} = \frac{\tau_f^3}{2c^2(M_{\rho_P^0}^2 - m_Z^2)}. \tag{B10}$$

3. The ρ_Π^3 partial decay rates

$$\Gamma(\rho_\Pi^3 \to W_L W_L) = \frac{1}{48\pi}\left(\frac{gF_\sigma^2}{4F_\pi^2}\right)^2 \frac{\left[M_{\rho_\Pi^3}^2 - 4m_W^2\right]^{3/2}}{M_{\rho_\Pi^3}^2}, \tag{B11}$$

$$\Gamma(\rho_\Pi^3 \to P^+P^-) = \frac{1}{48\pi}\left(\frac{gF_\sigma^2}{4F_\pi^2}\right)^2 \frac{\left[M_{\rho_\Pi^3}^2 - 4M_{P_\Pi^\pm}^2\right]^{3/2}}{M_{\rho_\Pi^3}^2}, \tag{B12}$$

$$\Gamma(\rho_\Pi^3 \to \bar{T}^\pm T^\mp) = \frac{1}{8\pi}\left(\frac{gF_\sigma^2}{4F_\pi^2}\right)^2 \frac{\left[M_{\rho_\Pi^3}^2 - 4M_{T^\pm}^2\right]^{3/2}}{M_{\rho_\Pi^3}^2}, \tag{B13}$$

$$\Gamma(\rho_\Pi^3 \to W_L^\pm P^\mp) = \frac{1}{24\pi}\left(\frac{gF_\sigma^2}{4F_\pi^2}\right)^2 \frac{\left[(M_{\rho_\Pi^3}^2 - (m_W + M_{P^\pm})^2)(M_{\rho_\Pi^3}^2 - (m_W - M_{P^\pm})^2)\right]^{3/2}}{M_{\rho_\Pi^3}^5}, \tag{B14}$$

$$\Gamma(\rho_\Pi^3 \to f\bar{f}) = \frac{N_c^{(f)}}{12\pi}(2e^2 gF_\sigma^2)^2 \left[\frac{[G_V^{\rho_\Pi^3}]^2(M_{\rho_\Pi^3}^2 + 2m_f^2) + [G_A^{\rho_\Pi^3}]^2(M_{\rho_\Pi^3}^2 - 4m_f^2)}{M_{\rho_\Pi^3}^2}\right]\sqrt{M_{\rho_\Pi^3}^2 - 4m_f^2},$$

$$\Gamma(\rho_\Pi^3 \to \phi\gamma) = \frac{1}{16\pi}\left(\frac{4egF_\sigma^2}{F_\phi}\right)^2 \left(\frac{M_{\rho_\Pi^3}^2 - M_\phi^2}{M_{\rho_\Pi^3}^3}\right), \tag{B15}$$

$$\Gamma(\rho_\Pi^3 \to \phi Z) = \frac{1}{16\pi}\left(\frac{2e(c^2-s^2)gF_\sigma^2}{scF_\phi}\right)^2 \frac{\sqrt{(M_{\rho_\Pi^3}^2 - (m_Z + M_\phi)^2)(M_{\rho_\Pi^3}^2 - (m_Z - M_\phi)^2)}}{M_{\rho_\Pi^3}^3}, \tag{B16}$$

where $W_L^\pm \equiv \Pi^\pm$ and

$$G_V^{\rho_\Pi^3} = \frac{Q_{\rm em}^f}{M_{\rho_\Pi^3}^2} + \frac{c^2 - s^2}{s^2c^2}\frac{\tau_3^f - 2s^2 Q_{\rm em}^f}{4(M_{\rho_\Pi^3}^2 - m_Z^2)}, \tag{B17}$$

$$G_A^{\rho_\Pi^3} = -\frac{c^2 - s^2}{s^2c^2}\frac{\tau_f^3}{4(M_{\rho_\Pi^3}^2 - m_Z^2)}. \tag{B18}$$

4. The $\rho_\Pi^\pm$ partial decay rates

$$\Gamma(\rho_\Pi^\pm \to W_L^\pm Z_L) = \frac{1}{48\pi}\left(\frac{gF_\sigma^2}{4F_\pi^2}\right)^2 \frac{\left[(M_{\rho_\Pi^\pm}^2 - (m_W + m_Z)^2)(M_{\rho_\Pi^\pm}^2 - (m_W - m_Z)^2)\right]^{3/2}}{M_{\rho_\Pi^\pm}^5}, \tag{B19}$$

$$\Gamma(\rho_\Pi^\pm \to P^\pm P^3) = \frac{1}{48\pi}\left(\frac{gF_\sigma^2}{4F_\pi^2}\right)^2 \frac{\left[(M_{\rho_\Pi^\pm}^2 - (M_{P^\pm} + M_{P^3})^2)(M_{\rho_\Pi^\pm}^2 - (M_{P^\pm} - M_{P^3})^2)\right]^{3/2}}{M_{\rho_\Pi^\pm}^5}, \tag{B20}$$

$$\Gamma(\rho_\Pi^\pm \to \bar{T}^\pm T^3) = \Gamma(\rho_\Pi^\pm \to \bar{T}^3 T^\pm)$$
$$= \frac{1}{8\pi}\left(\frac{gF_\sigma^2}{4F_\pi^2}\right)^2 \frac{\left[(M_{\rho_\Pi^\pm}^2 - (M_{T^\pm} + M_{T^3})^2)(M_{\rho_\Pi^\pm}^2 - (M_{T^\pm} - M_{T^3})^2)\right]^{3/2}}{M_{\rho_\Pi^\pm}^5}, \tag{B21}$$

$$\Gamma(\rho_\Pi^\pm \to W_L^\pm P^3) = \frac{1}{24\pi}\left(\frac{gF_\sigma^2}{4F_\pi^2}\right)^2 \frac{\left[(M_{\rho_\Pi^\pm}^2 - (m_W + M_{P^3})^2)(M_{\rho_\Pi^\pm}^2 - (m_W - M_{P^3})^2)\right]^{3/2}}{M_{\rho_\Pi^\pm}^5}, \tag{B22}$$

$$\Gamma(\rho_\Pi^\pm \to Z_L P^\pm) = \frac{1}{24\pi}\left(\frac{gF_\sigma^2}{4F_\pi^2}\right)^2 \frac{\left[(M_{\rho_\Pi^\pm}^2 - (m_Z + M_{P^\pm})^2)(M_{\rho_\Pi^\pm}^2 - (m_Z - M_{P^\pm})^2)\right]^{3/2}}{M_{\rho_\Pi^\pm}^5}, \tag{B23}$$

$$\Gamma(\rho_\Pi^\pm \to f_1 \bar{f}_2) = \frac{N_c^{(f)}}{48\pi} \left(\frac{e^2 g F_\sigma^2}{\sqrt{2} s^2 (M_{\rho_\Pi^\pm}^2 - m_W^2)} \right)^2 \frac{(2M_{\rho_\Pi^\pm}^4 - (m_{f_1}^2 + m_{f_2}^2) M_{\rho_\Pi^\pm}^2 - (m_{f_1}^2 - m_{f_2}^2)^2)}{M_{\rho_\Pi^\pm}^4} \times \frac{\sqrt{(M_{\rho_\Pi^\pm}^2 - (m_{f_1} + m_{f_2})^2)(M_{\rho_\Pi^\pm}^2 - (m_{f_1} - m_{f_2})^2)}}{M_{\rho_\Pi^\pm}}, \tag{B24}$$

$$\Gamma(\rho_\Pi^\pm \to \phi W^\pm) = \frac{1}{16\pi} \left(\frac{2egF_\sigma^2}{sF_\phi} \right)^2 \frac{\sqrt{(M_{\rho_\Pi^\pm}^2 - (m_W + M_\phi)^2)(M_{\rho_\Pi^\pm}^2 - (m_W - M_\phi)^2)}}{M_{\rho_\Pi^\pm}^3}, \tag{B25}$$

where $Z_L \equiv \Pi^3$.

[1] S. Weinberg, Phys. Rev. D **13**, 974 (1976).
[2] L. Susskind, Phys. Rev. D **20**, 2619 (1979).
[3] For a review see E. Farhi and L. Susskind, Phys. Rep. **74**, 277 (1981).
[4] K. Yamawaki, M. Bando, and K. Matumoto, Phys. Rev. Lett. **56**, 1335 (1986).
[5] For reviews, see, e.g., K. Yamawaki, in Proceedings of the 14th Symposium on Theoretical Physics: Dynamical Symmetry Breaking and Effective Field Theory, Cheju, Korea, 21 July, 1995 (to be published), arXiv:hep-ph/9603293; C. T. Hill and E. H. Simmons, Phys. Rep. **381**, 235 (2003); **390**, 553(E) (2004); F. Sannino, Acta Phys. Pol. B **40**, 3533 (2009); K. Yamawaki, in *Proceedings of the XLVIIth Rencontres de Moriond: Electroweak Interactions and Unified Theories, La Thuile, Italy, 2–9 March, 2013*, edited by E. Augé, J. Dumarchez, and J. T. T. Vaân (ARISF, Paris, France, 2013), p. 497.
[6] T. Akiba and T. Yanagida, Phys. Lett. **169B**, 432 (1986); T. W. Appelquist, D. Karabali, and L. C. R. Wijewardhana, Phys. Rev. Lett. **57**, 957 (1986); See also B. Holdom, Phys. Lett. **150B**, 301 (1985).
[7] M. Bando, K. Matumoto, and K. Yamawaki, Phys. Lett. B **178**, 308 (1986).
[8] S. Matsuzaki and K. Yamawaki, Phys. Rev. D **85**, 095020 (2012).
[9] S. Matsuzaki and K. Yamawaki, Phys. Rev. D **86**, 035025 (2012).
[10] S. Matsuzaki and K. Yamawaki, Phys. Lett. B **719**, 378 (2013); S. Matsuzaki, arXiv:1304.4882.
[11] S. Matsuzaki and K. Yamawaki, Phys. Rev. D **86**, 115004 (2012).
[12] M. Hashimoto and K. Yamawaki, Phys. Rev. D **83**, 015008 (2011).
[13] M. E. Peskin and T. Takeuchi, Phys. Rev. Lett. **65**, 964 (1990); Phys. Rev. D **46**, 381 (1992); B. Holdom and J. Terning, Phys. Lett. B **247**, 88 (1990); M. Golden and L. Randall, Nucl. Phys. **B361**, 3 (1991).
[14] G. Cacciapaglia, C. Csaki, C. Grojean, and J. Terning, Phys. Rev. D **71**, 035015 (2005); R. Foadi, S. Gopalakrishna, and C. Schmidt, Phys. Lett. B **606**, 157 (2005); R. S. Chivukula, E. H. Simmons, H. J. He, M. Kurachi, and M. Tanabashi, Phys. Rev. D **72**, 016008 (2005).
[15] K. Haba, S. Matsuzaki, and K. Yamawaki, Phys. Rev. D **82**, 055007 (2010).
[16] D. Elander, C. Nunez, and M. Piai, Phys. Lett. B **686**, 64 (2010); Z. Chacko and R. K. Mishra, Phys. Rev. D **87**, 115006 (2013); B. Bellazzini, C. Csaki, J. Hubisz, J. Serra, and J. Terning, Eur. Phys. J. C **74**, 2790 (2014); F. Coradeschi, P. Lodone, D. Pappadopulo, R. Rattazzi, and L. Vitale, J. High Energy Phys. 11 (2013) 057.
[17] Y. Aoki, T. Aoyama, M. Kurachi, T. Maskawa, K.-i. Nagai, H. Ohki, A. Shibata, K. Yamawaki, and T. Yamazaki (LatKMI Collaboration), Phys. Rev. D **87**, 094511 (2013); A. Hasenfratz, A. Cheng, G. Petropoulos, and D. Schaich, arXiv:1303.7129; T. Appelquist *et al.* (LSD Collaboration), arXiv:1405.4752 [Phys. Rev. D (to be published)].
[18] Y. Aoki, T. Aoyama, M. Kurachi, T. Maskawa, K. Miura, K.-i. Nagai, H. Ohki, E. Rinaldi, A. Shibata, K. Yamawaki, and T. Yamazaki (LatKMI Collaboration), Phys. Rev. D **89**, 111502(R) (2014); *Proc. Sci.*, LATTICE2013 (2013) 070.
[19] E. Farhi and L. Susskind, Phys. Rev. D **20**, 3404 (1979); S. Dimopoulos, Nucl. Phys. **B168**, 69 (1980).
[20] S. Dimopoulos and L. Susskind, Nucl. Phys. **B155**, 237 (1979); E. Eichten and K. D. Lane, Phys. Lett. **90B**, 125 (1980).
[21] S. Matsuzaki and K. Yamawaki, Phys. Rev. Lett. **113**, 082002 (2014).
[22] B. Holdom, Phys. Lett. B **198**, 535 (1987).
[23] M. Harada, M. Kurachi, and K. Yamawaki, Prog. Theor. Phys. **115**, 765 (2006).
[24] J. Jia, S. Matsuzaki, and K. Yamawaki, Phys. Rev. D **87**, 016006 (2013).
[25] M. Kurachi, S. Matsuzaki, and K. Yamawaki, arXiv:1403.0467 [Phys. Rev. D (to be published)].
[26] M. Bando, T. Kugo, S. Uehara, K. Yamawaki, and T. Yanagida, Phys. Rev. Lett. **54**, 1215 (1985); M. Bando, T. Kugo, and K. Yamawaki, Nucl. Phys. **B259**, 493 (1985); M. Bando, T. Fujiwara, and K. Yamawaki, Prog. Theor.

Phys. **79**, 1140 (1988); M. Bando, T. Kugo, and K. Yamawaki, Phys. Rep. **164**, 217 (1988).
[27] M. Harada and K. Yamawaki, Phys. Rep. **381**, 1 (2003).
[28] R. Casalbuoni, S. De Curtis, D. Dominici, and R. Gatto, Phys. Lett. **155B**, 95 (1985).
[29] A. R. Zerwekh and R. Rosenfeld, Phys. Lett. B **503**, 325 (2001).
[30] R. S. Chivukula, A. Grant, and E. H. Simmons, Phys. Lett. B **521**, 239 (2001).
[31] V. M. Abazov *et al.* (D0 Collaboration), Phys. Lett. B **681**, 224 (2009).
[32] V. A. Miransky, M. Tanabashi, and K. Yamawaki, Phys. Lett. B **221**, 177 (1989); Mod. Phys. Lett. A **04**, 1043 (1989).
[33] A. R. Zerwekh, Eur. Phys. J. C **46**, 791 (2006).
[34] A. Belyaev, R. Foadi, M. T. Frandsen, M. Jarvinen, F. Sannino, and A. Pukhov, Phys. Rev. D **79**, 035006 (2009).
[35] O. Castillo-Felisola, C. Corral, M. Gonzalez, G. Moreno, N. A. Neill, F. Rojas, J. Zamora, and A. R. Zerwekh, Eur. Phys. J. C **73**, 2669 (2013).
[36] A. Atre, R. S. Chivukula, P. Ittisamai, and E. H. Simmons, Phys. Rev. D **88**, 055021 (2013).
[37] R. S. Chivukula, R. Rosenfeld, E. H. Simmons, and J. Terning, in *Electroweak Symmetry Breaking and New Physics at the TeV Scale*, edited by T. L. Barklow, S. Dawson, H. E. Haber, and J. L. Siegrist (World Scientific, Singapore, 1997), p. 352.
[38] J. Pumplin, D. R. Stump, J. Huston, H. L. Lai, P. M. Nadolsky, and W. K. Tung, J. High Energy Phys. 07 (2002) 012.
[39] CMS Collaboration, Report No. CMS-PAS-EXO-12-059.
[40] ATLAS Collaboration, Report No. ATLAS-CONF-2013-017.
[41] CMS Collaboration, Report No. CMS-PAS-EXO-12-061.
[42] ATLAS Collaboration, Report No. ATLAS-CONF-2014-015.
[43] CMS Collaboration, Report No. CMS-PAS-EXO-12-025.
[44] V. Khachatryan *et al.* (CMS Collaboration), Phys. Rev. Lett. **105**, 211801 (2010).
[45] M. Kurachi, S. Matsuzaki, K. Terashi, and K. Yamawaki (to be published).