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Boston University
A Light Scalar in Low-Scale Technicolor

Antonio Delgado\textsuperscript{1*}, Kenneth Lane\textsuperscript{2,3†} and Adam Martin\textsuperscript{4‡}

\textsuperscript{1}Department of Physics, University of Notre Dame
Notre Dame, Indiana 46556
\textsuperscript{2}Department of Physics, Boston University
590 Commonwealth Avenue, Boston, Massachusetts 02215
\textsuperscript{3}LAPTH§, Université de Savoie, CNRS
B.P. 110, F-74941, Annecy-le-Vieux Cedex, France
\textsuperscript{4}Fermi National Accelerator Laboratory
Batavia, Illinois 60510

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Abstract

In addition to the narrow spin-one resonances $\rho_T$, $\omega_T$ and $a_T$ occurring in low-scale technicolor, there will be relatively narrow scalars in the mass range 200 to 600–700 GeV. We study the lightest isoscalar state, $\sigma_T$. In several important respects it is like a heavy Higgs boson with a small vev. It may be discoverable with high luminosity at the LHC where it is produced via weak boson fusion and likely has substantial $W^+W^-$ and $Z^0Z^0$ decay modes.

\textsuperscript{*}antonio.delgado@nd.edu
\textsuperscript{†}lane@physics.bu.edu
\textsuperscript{‡}aomartin@fnal.gov
\textsuperscript{§}Laboratoire de Physique Théorique d’Annecy-le-Vieux, UMR5108
1. Introduction  Walking technicolor [1] [2] [3] [4] is an asymptotically free gauge theory whose coupling $\alpha_{TC}$ runs very slowly for 100s, perhaps 1000s, of TeV above the electroweak breaking scale of a few 100 GeV. This is necessary so that extended technicolor (ETC) can generate sizable quark and lepton masses while suppressing otherwise fatal flavor-changing neutral current interactions [5]. A suitable walking $\alpha_{TC}$ occurs if the critical coupling for chiral symmetry breaking lies just below an (approximate) infrared fixed point [6] [7]. This requires a large number $N_D \gg 1$ of technifermion doublets in the fundamental representation of the TC gauge group or a few doublets in higher-dimensional representations [8] [9]. In the latter case, the constraints on ETC representations [5] almost always imply other technifermions in the fundamental representation as well. Thus, it is expected that there are technifermions whose technipions ($\pi_T$) have a decay constant $F_1^2 \ll F_\pi^2 = (246 \text{ GeV})^2$. This implies that bound states of the lightest technifermion doublet ($T_U, T_D$) have masses well below a TeV — greater than the experimental lower limit of 225–250 GeV [10] [11] and probably not more than 600–700 GeV. We refer to this as low-scale technicolor (LSTC). The most experimentally accessible bound states are the technivectors $V_T = \rho_T(I^G J^{PC} = 1^+ 1^-)$, $\omega_T(0^- 1^-)$ and $a_T(1^- 1^+)$, all of which may be produced as s-channel resonances of the Drell-Yan process in hadron and lepton colliders. The technipions $\pi_T(1^- 0^-)$ may be accessed through $V_T$ decays[1]. A central assumption of LSTC is that the lightest technihadrons may be treated in isolation, without significant mixing or other interference from higher-mass states. In a model with $N_D$ equivalent technifermion doublets, this requires that ETC leaves just one of them significantly lighter than the others. We also assume that the lightest technifermions are ordinary $SU(3)$-color singlets. For a more extensive discussion of LSTC, see Ref. [12].

Walking technicolor has another important phenomenological consequence: It enhances $M_{\pi_T}$ relative to the lightest $M_{\rho_T}$ so that the all-$\pi_T$ decay channels of $\rho_T$, $\omega_T$ and $a_T$ are closed [8]. The light $V_T$ then are very narrow, $\lesssim 1$ GeV, with principal decays to $\pi_T + W$ or $Z$, a pair of electroweak bosons (including one photon), and fermion-antifermion pairs [13] [14]. Many of these provide striking signatures, visible above backgrounds, within a limited mass range at the Tevatron and probably up to 600–700 GeV at the LHC [15] [16].

A walking $\alpha_{TC}$ also invalidates [17] [18] the QCD-based assumptions made to estimate the $S$-parameter for technicolor models [19] [20] [21] [22]. Further, it has been suggested that walking may cause $\rho_T$ and $a_T$ to be closer in mass than their QCD counterparts and to have approximately equal couplings to the vector and axial-vector parts of the electroweak currents [23] [24] [25] [26] [27] [28] [29]. That would eliminate the low-scale contribution to $S$. Determining the viability of this conjecture is the object of several recent papers employing lattice-gauge techniques [30] [31] [32]. Consequently, LSTC phenomenology now assumes that $M_{\rho_T} \simeq M_{\omega_T} \lesssim M_{a_T} < 2M_{\pi_T}$, with isospin-symmetric masses.

The main point of this paper is that, given this pattern of LSTC masses, there will be

1The isoscalar axial-vector $f_T(0^+ 1^{++})$ can be produced via weak vector boson fusion but not by the much stronger Drell-Yan process. We will not consider it further in this paper.
scalars $\sigma_0(0^+0^{++})$ and $\sigma_1(1^-0^{++})$, analogs of the $f_0$ and $a_0$ of QCD. As with the technivetors, we expect that ETC enhancements to their masses are less important than they are for technipions, so that $M_{\rho_T} \approx M_{\omega_T} \approx M_{\sigma_0,1} \approx M_{\alpha_T} \approx 2M_{\pi_T}$. Thus, these scalars are also relatively narrow (compared to a standard-model Higgs boson of the same mass). Like technipions, their couplings to $\bar{t}t$ are induced by ETC boson exchange. As for the $\pi_T$, they are of order $m_f/F_1$ — except for the top quark because ETC produces no more than 5–10 GeV of the top’s mass [33, 12]. Therefore, so long as the $\sigma_0$’s constituent fermions are color singlets, its main production mechanism at the LHC is weak vector boson fusion, not gluon fusion. The $\sigma_1$ is not produced by either mechanism. Therefore, in this Letter we concentrate on $\sigma_0$, which we refer to as $\sigma_T$ from now on. Its principal production and decay modes are similar to those of a heavy Higgs boson in a type-I two-Higgs-doublet model. It has a small vev, couples to the $W$ and $Z$, but only weakly to quarks and leptons. The LHC with $\sqrt{s} = 14$ TeV can produce $\sim 1$ to 100 $\sigma_T$ per fb$^{-1}$ in the mass range 200–700 GeV and decaying to $WW$ and $ZZ$ in all-leptonic or semileptonic channels. However, we shall see that several 100 fb$^{-1}$ will likely be required to discover $\sigma_T$ in these modes.

A light scalar with vacuum quantum numbers in walking technicolor has been proposed by a number of authors; see, e.g., Refs. [3, 34, 35, 36, 37, 38]. Usually, this is a “techni-dilaton”, a pseudo-Goldstone boson argued to arise as a consequence of spontaneous breaking of the theory’s approximate conformal invariance. We do not believe that $\sigma_T$ is a techni-dilaton. The main phenomenological difference between the two is in their couplings to matter and gauge fields. The techni-dilaton couples as $F_\pi/F_c$, where $F_c \gg F_\pi$ is the scale at which the conformal symmetry is broken while, as we see next, $\sigma_T$ couples as $F_1/F_\pi$. Although these couplings may be numerically similar, they have different origins. Furthermore, the technihadron partners of $\sigma_T$ and its place in their spectrum is specific to the version of LSTC considered here.

2. Effective Lagrangian for $\sigma_T$ in Low-Scale Technicolor In Ref. [12] two of us constructed an effective Lagrangian for LSTC. It describes the interactions at energies $\lesssim M_{\rho_T}$ of the lowest-lying technihadrons. A principal motivation for constructing $\mathcal{L}_{\text{eff}}$ was to provide a consistent treatment of the weak bosons, including the longitudinal $W_L$ and $Z_L$, which are common products of $V_T$ decays. The $V_T$ are included using the hidden local symmetry (HLS) formalism of Bando, et al. [39, 40]. The Lagrangian is based on $\mathcal{G} = SU(2) \otimes U(1) \otimes U(2)_L \otimes U(2)_R$, where $SU(2) \otimes U(1)$ is the electroweak gauge group and $U(2)_L \otimes U(2)_R$ is the HLS gauge group.

To describe the lightest technihadrons and mock up the heavier TC states contributing most to electroweak symmetry breaking (i.e., the isovector technipions of the other $N_D - 1$ technifermion doublets or the higher-scale states of a two-scale TC model), and to break all the gauge symmetries down to electromagnetic $U(1)$, we used nonlinear $\Sigma$-model fields $\Sigma_2$.

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2The $\sigma_0$ mixing with the light $\bar{t}t$ scalar that would be expected in a topcolor-assisted technicolor model is small because it is proportional to the $\sigma_0\bar{t}t$ coupling induced by ETC. Also, if QCD is any guide, $0^+0^{++}$ techni-glueballs are considerable heavier than $\sigma_0$ so that mixing between these scalars is not significant.
\[ \xi_L, \xi_R, \xi_M \text{ and } \Sigma_1 \equiv \xi_L \xi_M \xi_R, \text{ with covariant derivatives} \]
\[ D_\mu \xi_L = \partial_\mu \xi_L - ig t \cdot W_\mu + g' y_1 t_0 B_\mu ) \xi_L + ig T \xi_L t \cdot L_\mu \]
\[ D_\mu \xi_M = \partial_\mu \xi_M - ig T ( t \cdot L_\mu ) \xi_M - \xi_M t \cdot R_\mu \]
\[ D_\mu \xi_R = \partial_\mu \xi_R - ig T t \cdot R_\mu \xi_R + ig' \xi_R ( t_3 + y_1 t_0 ) B_\mu \]
\[ D_\mu \Sigma_{1,2} = \partial_\mu \Sigma_{1,2} - ig T \cdot W_\mu \Sigma_{1,2} + ig' \Sigma_{1,2} t_3 B_\mu , \]

where \( t \cdot L_\mu = \sum_{\alpha=0}^2 t_\alpha L_\mu^\alpha \) and \( t = \frac{1}{2} T \), \( t_0 = \frac{1}{2} I_1 \). The HLS gauge coupling \( g_L = g_R = g_T \) reflects the parity invariance of TC interactions and the expectation that \( I = 0, 1 \) technivectors are nearly degenerate. The hypercharge \( y_1 = Q_U + Q_D \) is the sum of electric charges of \( T_U \) and \( T_D \). The field \( \Sigma_2 \) contains the technipions that get absorbed by the \( W \) and \( Z \) bosons. We represent them as an isoriplet of \( F_2 \)-scale Goldstone bosons, where \( F_2^2 = ( F_\pi \cos \chi )^2 \gg F_1^2 \), and \( \chi \) is a mixing angle with, e.g., \( \sin \chi \approx 1/\sqrt{N_D} \) in an \( N_D \) doublet model.

Although \( \sigma_T \) was not included in the nonlinear fields in Ref. [12], it is easy to incorporate. In the unitary gauge, in which \( \Sigma_2, \xi_L, \xi_R \to 1 \) and \( \Sigma_1 = \xi_M \), we write \( F_1 \Sigma_1 = ( \sigma_T + F_1 ) E \) where \( E = \exp ( 2i t \cdot \pi_T ) / F_1 \). We do not consider other light scalars to arise from \( \xi_{L,R} \) because they are not expected in the low-lying spectrum of \( T \)-hadrons. The complete effective Lagrangian is

\[ \mathcal{L}_{\text{eff}} = \mathcal{L}_\Sigma + \mathcal{L}_{\text{gauge}} + \mathcal{L}_{f} + \mathcal{L}_{\text{WZW}} + \mathcal{L}_{M^2} + \mathcal{L}_{\pi_T f} , \]

where, in unitary gauge,

\[ \mathcal{L}_\Sigma = \frac{1}{4} F_2^2 \text{Tr} | g t \cdot W_\mu - g' t_3 B_\mu |^2 + \frac{1}{4} (a + c) ( \partial_\mu \sigma_T )^2 \]
\[ + \frac{1}{4} ( \sigma_T + F_1 )^2 \left\{ a \text{Tr} ( \partial_\mu E - i g t \cdot W_\mu E - g' t_3 B_\mu ) \right\}^2 \]
\[ + b \left\{ \text{Tr} [ g t \cdot W_\mu + g' y_1 t_0 B_\mu - g T t \cdot L_\mu ]^2 + \text{Tr} [ g' ( t_3 + y_1 t_0 ) B_\mu - g T t \cdot R_\mu ]^2 \right\} \]
\[ + c \text{Tr} ( \partial_\mu E + ig T ( E t \cdot R_\mu - t \cdot L_\mu ) E )^2 \]
\[ + d \text{Tr} [ ( g E^t t \cdot W_\mu + g' t_3 E^t B_\mu + g T ( t \cdot R_\mu E^t - E^t t \cdot L_\mu ) ) \]
\[ \times ( \partial_\mu E + ig T ( E t \cdot R_\mu - t \cdot L_\mu ) E ) ] \}
\[ - \frac{i f}{2 g_T F_1^2 } ( \sigma_T + F_1 )^2 \]
\[ \times \left\{ \text{Tr} [ ( \partial_\mu E + i g T ( E t \cdot R_\mu - t \cdot L_\mu ) ) E^t ( \partial_\nu E + i g T ( E t \cdot R_\nu - t \cdot L_\nu ) ) E^t t \cdot L_\mu \nu \]
\[ + E^t ( \partial_\mu E + i g T ( E t \cdot R_\mu - t \cdot L_\mu ) ) E^t ( \partial_\nu E + i g T ( E t \cdot R_\nu - t \cdot L_\nu ) ) t \cdot R_\mu \nu ] \right\} . \]

A “simplicity principle” was adopted in writing \( \mathcal{L}_\Sigma \): only the lowest-dimension operators needed to describe the experimentally important LSTC processes — mainly \( V_T \) two-body production and decay vertices — were kept. The dimensionless parameters \( a, \ldots, f \) are

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3These \( \sigma_T \) and \( \pi_T \) fields are not yet canonically normalized. Also, we have assumed the isoscalar \( \pi_T^0 \) to be much heavier than \( \rho_T \) and integrated it out. See Ref. [12] for discussions of these points.
nominally of $\mathcal{O}(1)$. The gauge-field Lagrangian $\mathcal{L}_{\text{gauge}}$ has the standard form; $\mathcal{L}_{ff}$ is the usual coupling of quarks and leptons to the $SU(2) \otimes U(1)$ gauge bosons only; $\mathcal{L}_{M^2}$ describes $\pi_T$ and $\sigma_T$ masses and $\mathcal{L}_{\pi_T f f}$ the $\pi_T f f$ couplings. The Wess-Zumino-Witten interaction $\mathcal{L}_{\text{WZW}}$ reflects the anomalous global and HLS gauge symmetries of the underlying theory $\text{[11]}$. It is essential for describing radiative decays of $\rho_T$ and $\omega_T$ as well as $\pi_T^0 \to \gamma \gamma$. Processes computed in tree approximation, including production and scattering of $W_L$ and $Z_L$, behave at high energies, $\sqrt{s} \gg M_{\rho_T}$, as they do in the standard model without a Higgs boson.\footnote{By itself, this $\sigma_T$ does not alter this situation much because of its small $\mathcal{O}(F_1)$ coupling to $WW$.} The masses of the $W$, $Z$ and $V_T$, with $M_W/M_Z \cos \theta_W = 1$ and $M_{\omega_T} = M_{\rho_T} + \mathcal{O}((gy_1 \sin \theta_W / g_T)^2)$, follow from $\mathcal{L}_\Sigma$. See Ref $\text{[12]}$ for details.

The mixing angle $\chi$ characterizing the contribution of the low $F_1$-scale to electroweak symmetry breaking is

$$\sin \chi \simeq F_1/F_\pi,$$  

where $F_\pi = \sqrt{F_2^2 + AF_1^2} = 246$ GeV.

In the $U(2)_V$ limit in which $M_{\omega_T} = M_{\rho_T}$, all quantities of phenomenological interest can be expressed in terms of $M_{\rho_T}$, $M_{\omega_T}$, $F_\pi$, the number of technicolors $N_{TC}$, $\sin \chi$, $y_1 = Q_U + Q_D$, and three mass parameters — $M_{V_1}$, $M_{A_1}$, $M_{A_2}$ — that control the strength of dimension-five operators involved in decays of technifermions to photons and transversely-polarized weak bosons.\footnote{We assume for simplicity that the TC gauge group is $SU(N_{TC})$ and the technifermions transform as the fundamental $N_{TC}$. This affects WZW interactions whose dimension-five operators are $\propto 1/M_{V_1}$.} In particular, the effective coupling $g_{\rho_T \pi_T \pi_T}$ for $\rho_T \to \pi_T \pi_T$, $\pi_T W_L$ and $W_L W_L$ is \footnote{This is motivated by the mass controlling ordinary $\rho, \omega \to \gamma \pi^0$, namely, $M_{V_1(QCD)} \simeq 700$ MeV $\simeq M_{\rho}$.} \footnote{This ability to express unknown couplings in terms of the natural inputs of LSTC is convenient. In the commonly used case, $M_{V_1} = M_{A_1} = M_{\rho_T}$, $g_{\rho_T \pi_T \pi_T}$ is given by the KSFR-like relation $g_{\rho_T \pi_T \pi_T} = M_{\rho_T} / 2 F_\pi \sin \chi$, while $g_T = 8 \sqrt{2} \pi F_\pi / M_{\rho_T}$. For light $\rho_T$, this relation has the consequence that $\rho_T \to W_L \pi_T$, $W_L W_L$ are significantly suppressed (and radiative $\rho_T$ decay branching ratios enhanced!) relative to those calculated using PYTHIA \text{[42]} where the default, scaled from QCD, is $g_T = g_{\rho_T \pi_T \pi_T} = \sqrt{4 \pi(2.16)(3/N_{TC})}$.}

$$g_{\rho_T \pi_T \pi_T} = \frac{M_{\rho_T}^2}{\sqrt{2} g_T (F_\pi \sin \chi)^2} \left[ 1 + (f - 1) \frac{M_{A_2}^2}{M_{A_1}^2} \right],$$

in which

$$g_T = \frac{16 \sqrt{2} \pi^2 M_{A_1} F_\pi \sin \chi}{N_{TC} M_{V_1} (M_{A_1} + M_{A_2})}, \quad f = \frac{(4 \pi M_{A_1} F_\pi \sin \chi)^2}{N_{TC} M_{V_1} M_{A_2}^2 (M_{A_1} + M_{A_2})}.$$

This ability to express unknown couplings in terms of the natural inputs of LSTC is convenient. In the commonly used case, $M_{V_1} = M_{A_1} = M_{\rho_T}$, $g_{\rho_T \pi_T \pi_T}$ is given by the KSFR-like relation $g_{\rho_T \pi_T \pi_T} = M_{\rho_T} / 2 F_\pi \sin \chi$, while $g_T = 8 \sqrt{2} \pi F_\pi / M_{\rho_T}$. For light $\rho_T$, this relation has the consequence that $\rho_T \to W_L \pi_T$, $W_L W_L$ are significantly suppressed (and radiative $\rho_T$ decay branching ratios enhanced!) relative to those calculated using PYTHIA \text{[42]} where the default, scaled from QCD, is $g_T = g_{\rho_T \pi_T \pi_T} = \sqrt{4 \pi(2.16)(3/N_{TC})}$.

At this point, we mention that the contribution of the $F_1$-scale $\rho_T$ and $a_T$ to the $S$-parameter is \footnote{This ability to express unknown couplings in terms of the natural inputs of LSTC is convenient. In the commonly used case, $M_{V_1} = M_{A_1} = M_{\rho_T}$, $g_{\rho_T \pi_T \pi_T}$ is given by the KSFR-like relation $g_{\rho_T \pi_T \pi_T} = M_{\rho_T} / 2 F_\pi \sin \chi$, while $g_T = 8 \sqrt{2} \pi F_\pi / M_{\rho_T}$. For light $\rho_T$, this relation has the consequence that $\rho_T \to W_L \pi_T$, $W_L W_L$ are significantly suppressed (and radiative $\rho_T$ decay branching ratios enhanced!) relative to those calculated using PYTHIA \text{[42]} where the default, scaled from QCD, is $g_T = g_{\rho_T \pi_T \pi_T} = \sqrt{4 \pi(2.16)(3/N_{TC})}$.}

$$S_1(\rho_T, a_T) = \frac{8 \pi}{g_T^2} \left( 1 - \frac{M_{A_2}^2}{M_{A_1}^2} \right).$$

(6)
This can be made as small as desired, and it vanishes in the commonly used case mentioned above. The contribution of \( \sigma_T \) and the \( \pi_T \) is

\[
S_1(\sigma_T, \pi_T) = \frac{1}{2\pi} \int_0^1 dx \, x(1-x) \ln \left[ \frac{M_{\pi_T}^2 x + M_{\sigma_T}^2 (1-x)}{M_{\pi_T}^2} \right].
\]  

(7)

This is positive for \( M_{\sigma_T} > M_{\pi_T} \), but also quite small, \( \mathcal{O}(10^{-2}) \), for reasonable choices of these masses. Because of the built-in isospin symmetry, \( T_1 = 0 \) and other precision parameters are negligibly small \([12]\).

3. \( \sigma_T \) at the LHC Since \( \sigma_T \) couples weakly to \( \bar{t}t \), its major production modes at the LHC are weak vector boson fusion (VBF), \( W^+W^- \) and \( Z^0 Z^0 \to \sigma_T \), and associated production, \( W^\pm, Z^0 \to \sigma_T + W^\pm, Z^0 \). Its principal decay modes are \( \sigma_T \to W^+W^- \), \( Z^0 Z^0 \), \( W^\pm \pi_T^\mp, Z^0 \pi_T^0 \). The Lagrangian describing these couplings comes from \( \mathcal{L}_\Sigma \) and is given by

\[
\mathcal{L}_{\sigma_T} = \mathcal{K}_{\sigma_T} \left[ \frac{1}{4} F_\pi \sin \chi \left( 2g^2 W^+\mu W^-_\mu + (g^2 + g'^2)Z^\mu Z_\mu \right) \right. \\
\left. - \cos \chi \left( g(W^+_\mu \partial^\mu \pi^-_T + W^-_\mu \partial^\mu \pi^+_T) + \sqrt{g^2 + g'^2} Z_\mu \partial^\mu \pi^0_T \right) \right],
\]  

(8)

where

\[
\mathcal{K} = \frac{(4\pi M_{A_1} F_\pi \sin \chi)^4 - (N_{TC} M_{\rho_T} M_{\mu_1} (M_{A_1}^2 - M_{A_2}^2))^2}{(4\pi M_{A_1} F_\pi \sin \chi)^2 \left[ (4\pi M_{A_1} F_\pi \sin \chi)^4 + (N_{TC} M_{\sigma_T} M_{\mu_1} (M_{A_1}^2 - M_{A_2}^2))^2 \right]^2} \to 1 \text{ as } M_{A_1} - M_{A_2} \to 0.
\]  

(9)

In the limit \( \mathcal{K} = 1 \), \( \mathcal{L}_{\sigma_T} \) has the same form as the corresponding interaction of a Higgs boson with vev \( v_1 = F_\pi \sin \chi \). This lowers the \( \sigma_T \) production and decay rates by \( \sin^2 \chi \), nominally an order of magnitude for the value \( \sin \chi \simeq 1/3 \) assumed in most LSTC studies. Including the \( W/Z\pi_T \) decay modes, which a Higgs boson does not have, while ignoring the small \( \bar{f}f \) contributions to its width, we find \( \Gamma_{\sigma_T} \simeq 5 \text{ (65) GeV for } M_{\sigma_T} = 300 \text{ (600) GeV. These widths assume } M_{\pi_T} = 0.55 M_{\sigma_T}. \) They are roughly half this large for \( M_{\pi_T} = 0.65 M_{\sigma_T}. \) The width of a 300 (600) GeV standard-model Higgs is about twice this large, 9 (125) GeV.

The VBF and associated production rates of \( \sigma_T \) in \( pp \) collisions at 14 TeV and decay to \( W^+W^- \) in the semileptonic \( jj\ell\nu \) and leptonic \( \ell\nu\ell'\nu' \) modes are shown in Fig. [1]. Here, \( \ell, \ell' = e, \mu \). The upper limit \( B(\sigma_T \to W^+W^-) = 2/3 \) was assumed. We show \( \sigma_T \) decay branching ratios in Fig. [2] for \( M_{\pi_T}/M_{\sigma_T} = 0.55 \) and 0.65.

The raw VBF rates to \( WW \) are large, but tagging the energetic forward jets is essential. The cleanest \( WW \) mode is the leptonic one. The possibility of discovering the standard-model Higgs boson in this way, including a Higgs in the mass range of interest to us, was studied in Ref. [13]. Rescaling the signal significances found there to a luminosity of 300 fb\(^{-1}\) yields Fig. [3]. This result is not promising for \( \sigma_T \) discovery.

A less explored and more promising possibility is VBF production followed by the semileptonic modes of \( WW \) and \( ZZ \). A complication for the semileptonic mode is that central jets
Figure 1: The $\sigma_T \rightarrow W^+W^-$ production rates in $pp$ collisions at $\sqrt{s} = 14$ TeV for $B(\sigma_T \rightarrow W^+W^-) = 2/3$ and $K = 1$ in Eq. (9); $\sin \chi = 1/3$ (black) or 1/2 (gray). Left: Weak boson fusion rates at $\int L dt = 1$ fb$^{-1}$. Right: Associated production $W^\pm \rightarrow W^\pm + \sigma_T$ at 100 fb$^{-1}$.

Leptons include $e$ and $\mu$. Must be vetoed to suppress $t\bar{t}$ + jets and $W/Z$ + jets backgrounds while retaining enough hadronic activity to reconstruct the hadronically decaying gauge boson. This issue was addressed in CMS Note 2001/050, which studied VBF production of a heavy Higgs ($M_H = 300$ and 600 GeV) at the LHC with $\int L dt = 30$ fb$^{-1}$ at $\mathcal{L} = 10^{33}$ cm$^{-2}$ s$^{-1}$. The note found a central-jet veto effective in suppressing $t\bar{t}$ and a $W$-mass cut on $W \rightarrow jj$ suppressed $W/Z +$ jets. Translating their results to $\sigma_T$-production rates at a luminosity of 300 fb$^{-1}$, we expect $S/\sqrt{S+B} = 0.7$ (1.6) for $M_{\pi} = 300$ GeV, $\sin \chi = 1/3$ and $M_{\pi}/M_{\pi} = 0.55$ (0.65). The corresponding significances are 2.9 (4.7) for $\sin \chi = 1/2$. They are 1.5 to 2 times smaller than these for $M_{\pi} = 600$ GeV.

This CMS study was traditional in that it did not employ recently developed jet substructure techniques; see e.g., Refs. [15, 44, 45, 46, 47, 48]. If a hadronically decaying gauge boson (more generally, any resonance) is sufficiently boosted, its daughter partons and their corresponding final state radiation can be captured within a single “fat” jet with radius $R \sim 1.2 – 1.5$. This fat jet will have high mass, $\sim M_{W,Z}$ and contain interior structure — two subjet hotspots. Jet substructure techniques are easy to incorporate into VBF analyses. After tagging the forward jets, the remaining hadronic activity can be grouped into fat jets and analyzed for substructure. Then, once a hadron gauge boson is identified, a jet veto can be applied to the remaining hadronic energy to further suppress $t\bar{t}$ + jets, etc.

One problem facing substructure techniques is that the large jet area captures a lot of contamination from initial state radiation and the underlying event — energy not associated
with the resonance. However, in VBF there is no color information exchanged between the initial quarks \cite{49, 50, 51}. This makes VBF events less prone to these effects and therefore well-suited to the use of substructure.

In a pioneering paper on jet substructure \cite{52}, the method was applied to the search for strongly-interacting \(WW\) resonances produced via VBF. The analysis focused on discovering scalar and vector resonances with masses greater than 1 TeV. We cannot simply recycle the backgrounds of that analysis because it combined substructure and VBF cuts with hard kinematic cuts on the \(p_T\) of the reconstructed \(W\) that resonances like \(\sigma_T\) in the \(\sim 300-700\) GeV range will not pass. Reference \cite{52} concluded that very heavy resonances could be discovered and their decay angular distributions studied with 100 fb\(^{-1}\) of LHC data. We expect a similar conclusion for a lighter, more weakly interacting resonance. The analysis in Ref. \cite{52} was done at the particle level, meaning that showering and hadronization were included, but no detector effects other than fiducial volume cuts and rudimentary particle ID efficiencies were applied. A detailed study incorporating more realistic detector effects is needed.

The \(\sigma_T \to ZZ\) process suffers from the small \(Z \to \ell^+\ell^-\) branching ratio and a production cross section half as large as for \(WW\), but has the advantages of no missing \(E_T\) and and no \(W \to \ell\nu\) reconstruction ambiguity. In Ref. \cite{53}, jet substructure techniques were applied to the decays of heavy Higgs bosons into \(ZZ \to jj\ell^+\ell^-\). Both gluon fusion and VBF production modes were studied and saw promising results. However, this analysis focused on substructure and did not combine hadronic \(Z\) identification with the usual VBF selections.
and cuts, particularly on forward jets. Again, a more detailed and VBF-specific study is needed.

We have argued that the spectrum of low-scale technicolor has a relatively light and relatively narrow $0^+0^{++}$ scalar, $\sigma_T$. Like a heavy Higgs boson, $\sigma_T$ decays mainly to $WW$ and $ZZ$ and it is produced at the LHC via $WW$ and $ZZ$ fusion, albeit at a rate suppressed by $\sin^2 \chi \simeq F_1^2/F_\pi^2 \sim 0.1$. The most promising final states are the semileptonic ones in which one $W/Z$ decays hadronically. Tagging the forward jets of the fusion process and using jet substructure to identify the hadronic $W$ or $Z$ decay may make $\sigma_T$ discovery possible, although it seems likely that a luminosity of several $100 \text{fb}^{-1}$ will be required. Detailed, detector-specific simulations are necessary to decide this question, one we think is well worth answering. If a standard-model-like Higgs boson is not found at the LHC with luminosities typical of gluon fusion, it will be important to determine whether any scalar exists. If the light technivectors $\rho_T$, $\omega_T$, $\sigma_T$ are found and they have nearly equal masses, the discovery of $\sigma_T$ is then important to understanding the spectroscopy of low-scale technicolor.

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