Viability of Strongly Coupled Scenarios with a Light Higgs-like Boson

Antonio Pich,[1] Ignasi Rosell,[1,2] and Juan José Sanz-Cillero[3,4]

[1]*Departament de Física Teòrica, IFIC, Universitat de València-CSIC, Apartat de Correus 22085, 46071 València, Spain*
[2]*Departamento de Ciencias Físicas, Matemáticas y de la Computación, Universidad CEU Cardenal Herrera, 46115 Alfara del Patriarca, València, Spain*
[3]*INFN, Sezione di Bari, Via Orabona 4, 70126 Bari, Italy*
[4]*Department of Physics, Peking University, Beijing 100871, People's Republic of China*

(Received 8 January 2013; published 2 May 2013)

We present a one-loop calculation of the oblique S and T parameters within strongly coupled models of electroweak symmetry breaking with a light Higgs-like boson. We use a general effective Lagrangian, implementing the chiral symmetry breaking $SU(2)_L \otimes SU(2)_R \to SU(2)_{L+R}$ with Goldstone bosons, gauge bosons, the Higgs-like scalar, and one multiplet of vector and axial-vector massive resonance states. Using a dispersive representation and imposing a proper ultraviolet behavior, we obtain S and T at the next-to-leading order in terms of a few resonance parameters. The experimentally allowed range forces the vector and axial-vector states to be heavy, with masses above the TeV scale, and suggests that the Higgs-like scalar should have a WW coupling close to the standard model one. Our conclusions are generic and apply to more specific scenarios such as the minimal $SO(5)/SO(4)$ composite Higgs model.

DOI: 10.1103/PhysRevLett.110.181801 PACS numbers: 12.60.Fr, 12.39.Fe, 12.60.Nz, 12.60.Rc

Introduction.—A new Higgs-like boson around 126 GeV has just been discovered at the LHC [1]. Although its properties are not well measured yet, it complies with the expected behavior and therefore it is a very compelling candidate to be the standard model (SM) Higgs boson. An obvious question to address is to what extent alternative scenarios of electroweak symmetry breaking (EWSB) can already be discarded or strongly constrained. In particular, what are the implications for strongly coupled models where the electroweak symmetry is broken dynamically?

The existing phenomenological tests have confirmed the $SU(2)_L \otimes SU(2)_R \to SU(2)_{L+R}$ pattern of symmetry breaking, giving rise to three Goldstone bosons, which in the unitary gauge become the longitudinal polarizations of the gauge bosons. When the $U(1)_Y$ coupling g' is neglected, the electroweak Goldstone boson dynamics is described at low energies by the same Lagrangian as that of the QCD pions, replacing the pion decay constant by the EWSB scale $v = (\sqrt{2}G_F)^{-1/2} = 246$ GeV [2]. Contrary to the SM, in strongly coupled scenarios the symmetry is nonlinearly realized and one expects the appearance of massive resonances generated by the nonperturbative interaction.

The dynamics of Goldstone bosons and massive resonance states can be analyzed in a generic way by using an effective Lagrangian, based on symmetry considerations. The theoretical framework is completely analogous to the resonance chiral theory description of QCD at GeV energies [3]. Using these techniques, we investigated in Ref. [4] the oblique S parameter [5], characterizing the electroweak boson self-energies, within Higgsless strongly coupled models. By the adoption of a dispersive approach and imposition of a proper UV behavior, it was shown there that it is possible to calculate S at the next-to-leading order, i.e., at one loop. We found that in most strongly coupled scenarios of EWSB, a high resonance mass scale is required, $M_V > 1.8$ TeV, to satisfy the stringent experimental limits.

The recent discovery of a Higgs-like boson makes updating the analysis mandatory, including the light-scalar contributions. In addition, we will also present a corresponding one-loop calculation of the oblique T parameter, which allows us to perform a correlated analysis of both quantities. S measures the difference between the off-diagonal W^3B correlator and its SM value, whereas T parametrizes the breaking of custodial symmetry [5]. More precisely, T measures the difference between the W^3W^3 and W^+W^- correlators, subtracting the SM contribution. The explicit definitions of S and T are given in Refs. [4,5]. Previous one-loop analyses can be found in Refs. [6–8].

Theoretical framework.—We have considered a low-energy effective theory containing the SM gauge bosons coupled to the electroweak Goldstone bosons, one light-scalar state S_1 with mass $m_{S_1} = 126$ GeV, and the lightest vector and axial-vector resonance multiplets $V_{\mu\nu}$ and $A_{\mu\nu}$. We only assume the SM pattern of EWSB; i.e., the theory is symmetric under $SU(2)_L \otimes SU(2)_R$ and becomes spontaneously broken to the diagonal subgroup $SU(2)_{L+R}$. S_1 is taken to be singlet under $SU(2)_{L+R}$, while $V_{\mu\nu}$ and $A_{\mu\nu}$ are triplets (singlet vector and axial-vector contributions are absent at the order we are working). To build the Lagrangian, we only consider operators with the lowest number of derivatives, as higher-derivative terms are either proportional to the equations of motion or tend to violate the expected short-distance behavior of the Green's functions [3]. We will need the interactions

$$\mathcal{L} = \frac{v^2}{4}\langle u_\mu u^\mu\rangle\left(1+\frac{2\omega}{v}S_1\right) + \frac{F_A}{2\sqrt{2}}\langle A_{\mu\nu}f_-^{\mu\nu}\rangle + \frac{F_V}{2\sqrt{2}}\langle V_{\mu\nu}f_+^{\mu\nu}\rangle + \frac{iG_V}{2\sqrt{2}}\langle V_{\mu\nu}[u^\mu, u^\nu]\rangle + \sqrt{2}\lambda_1^{SA}\partial_\mu S_1\langle A^{\mu\nu}u_\nu\rangle, \quad (1)$$

plus the standard gauge boson and resonance kinetic terms. The three Goldstone fields $\vec{\pi}(x)$ are parametrized through the matrix $U = u^2 = \exp\{i\vec{\sigma}\,\vec{\pi}/v\}$, $u^\mu = -iu^\dagger D^\mu U u^\dagger$ with D^μ the appropriate gauge-covariant derivative, and $\langle A\rangle$ stands for the trace of the 2×2 matrix A. We follow the notation from Ref. [4]. The first term in Eq. (1) gives the Goldstone Lagrangian, present in the SM, plus the scalar-Goldstone interactions. For $\omega = 1$, one recovers the $S_1 \to \pi\pi$ vertex of the SM. The calculation will be performed in the Landau gauge, which eliminates the mixing between Goldstone and gauge bosons.

The oblique parameter S receives tree-level contributions from vector and axial-vector exchanges [5], while T is identically zero at lowest order (LO),

$$S_{\rm LO} = 4\pi\left(\frac{F_V^2}{M_V^2} - \frac{F_A^2}{M_A^2}\right), \qquad T_{\rm LO} = 0. \quad (2)$$

To compute the one-loop contributions, we use the dispersive representation of S introduced by Peskin and Takeuchi [5], whose convergence requires a vanishing spectral function at short distances,

$$S = \frac{16\pi}{g^2\tan\theta_W}\int_0^\infty \frac{dt}{t}[\rho_S(t) - \rho_S(t)^{\rm SM}], \quad (3)$$

with $\rho_S(t)$ the spectral function of the W^3B correlator [4,5]. We work at lowest order in g and g', and only the lightest two-particle cuts have been considered, i.e., two Goldstone bosons or one Goldstone boson plus one scalar resonance. $V\pi$ and $A\pi$ contributions were shown to be very suppressed in Ref. [4].

The calculation of T is simplified by noticing that, up to corrections of $\mathcal{O}(m_W^2/M_R^2)$,

$$\alpha T = \frac{Z^{(+)}}{Z^{(0)}} - 1, \quad (4)$$

where $Z^{(+)}$ and $Z^{(0)}$ are the wave-function renormalization constants of the charged and neutral Goldstone bosons computed in the Landau gauge [9]. A further simplification occurs by setting to zero g, which does not break the custodial symmetry, so only the B-boson exchange produces an effect in T. This approximation captures the lowest order contribution to T in its expansion in powers of g and g'. Again, only the lowest two-particle cuts have been considered, i.e., the B boson plus one Goldstone or one scalar resonance.

Figure 1 shows the computed one-loop contributions to S and T. Requiring the W^3B correlator to vanish at high energies also implies a good convergence of the Goldstone boson self-energies, at least for the two-particle cuts we have considered. Therefore, their difference obeys an unsubtracted dispersion relation, which enables us to also compute T through the dispersive integral

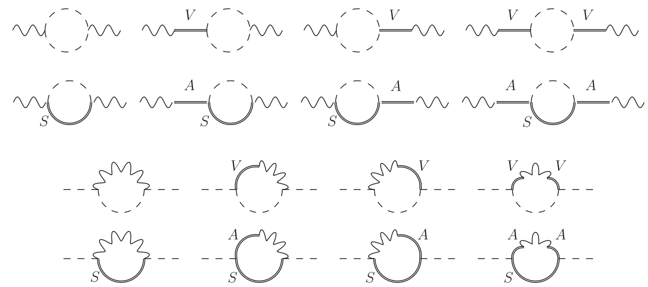

FIG. 1. NLO contributions to S (two first lines) and T (two last lines). A dashed (double) line stands for a Goldstone (resonance) boson, and a curved line represents a gauge boson.

$$T = \frac{4\pi}{g'^2\cos^2\theta_W}\int_0^\infty \frac{dt}{t^2}[\rho_T(t) - \rho_T(t)^{\rm SM}], \quad (5)$$

with $\rho_T(t)$ the spectral function of the difference of the neutral and charged Goldstone boson self-energies.

Short-distance constraints.—Fixing the scalar mass to $m_{S_1} = 126$ GeV, we have seven undetermined parameters: M_V, M_A, F_V, G_V, F_A, ω, and λ_1^{SA}. The number of unknown couplings can be reduced using short-distance information.

Assuming that weak isospin and parity are good symmetries of the strong dynamics, the W^3B correlator is proportional to the difference of the vector and axial-vector two-point Green's functions [5]. In asymptotically free gauge theories, this difference vanishes at $s\to\infty$ as $1/s^3$ [10], implying two superconvergent sum rules, known as the first and second Weinberg sum rules (WSRs) [11], which at LO give the relations

$$F_V^2 - F_A^2 = v^2, \qquad F_V^2 M_V^2 - F_A^2 M_A^2 = 0. \quad (6)$$

This determines F_V and F_A in terms of the resonance masses, leading to

$$S_{\rm LO} = \frac{4\pi v^2}{M_V^2}\left(1 + \frac{M_V^2}{M_A^2}\right). \quad (7)$$

Since the WSRs also imply $M_A > M_V$, this prediction turns out to be bounded by [4]

$$\frac{4\pi v^2}{M_V^2} < S_{\rm LO} < \frac{8\pi v^2}{M_V^2}. \quad (8)$$

It is likely that the first WSR is also true in gauge theories with nontrivial ultraviolet fixed points [7,12], whereas the second WSR is questionable in some scenarios. If only the first WSR is considered, but still assuming the hierarchy $M_A > M_V$, one obtains the lower bound [4]

$$S_{\rm LO} = 4\pi\left\{\frac{v^2}{M_V^2} + F_A^2\left(\frac{1}{M_V^2} - \frac{1}{M_A^2}\right)\right\} > \frac{4\pi v^2}{M_V^2}. \quad (9)$$

The possibility of an inverted mass ordering of the vector and axial-vector resonances [12] would turn this lower bound into the upper bound $S_{\rm LO} < 4\pi v^2/M_V^2$. Note that if the splitting of the vector and axial-vector resonances was small, the prediction of $S_{\rm LO}$ would be close to the bound.

At the next-to-leading order (NLO) the computed W^3B correlator should also satisfy the proper short-distance behavior. The $\pi\pi$ and $S\pi$ spectral functions would have an unphysical growth at large momentum transfer unless $F_V G_V = v^2$ and $F_A \lambda_1^{SA} = \omega v$. The first constraint guarantees a well-behaved vector form factor [3], whereas the second relates the axial and scalar couplings. Once these relations are enforced, the Goldstone boson self-energies are convergent enough to allow for an unambiguous determination of T in terms of masses and ω. Neglecting terms of $\mathcal{O}(m_{S_1}^2/M_{V,A}^2)$,

$$T = \frac{3}{16\pi \cos^2\theta_W}\left[1 + \log\frac{m_H^2}{M_V^2} - \omega^2\left(1 + \log\frac{m_{S_1}^2}{M_A^2}\right)\right], \quad (10)$$

where m_H is the SM reference Higgs mass adopted to define S and T. Notice that on taking $m_H = m_{S_1}$ and $\omega = 1$ (the SM value), T vanishes when $M_V = M_A$ as it should.

To enforce the second WSR at NLO, one needs the additional constraint $\omega = M_V^2/M_A^2$ (constrained to the range $0 \le \omega \le 1$). One can then obtain a NLO determination of S in terms of M_V and M_A

$$S = 4\pi v^2\left(\frac{1}{M_V^2} + \frac{1}{M_A^2}\right) + \frac{1}{12\pi} \times \left[\log\frac{M_V^2}{m_H^2} - \frac{11}{6} + \frac{M_V^2}{M_A^2}\log\frac{M_A^2}{M_V^2} - \frac{M_V^4}{M_A^4}\left(\log\frac{M_A^2}{m_{S_1}^2} - \frac{11}{6}\right)\right], \quad (11)$$

where terms of $\mathcal{O}(m_{S_1}^2/M_{V,A}^2)$ have been neglected. Taking $m_H = m_{S_1}$, the correction to the LO result vanishes when $M_V = M_A$ ($\omega = 1$); in this limit, the NLO prediction reaches the LO upper bound in Eq. (8).

If only the first WSR is considered, one can still obtain a lower bound at NLO in terms of M_V, M_A, and ω,

$$S \ge \frac{4\pi v^2}{M_V^2} + \frac{1}{12\pi}\left[\log\frac{M_V^2}{m_H^2} - \frac{11}{6} - \omega^2\left(\log\frac{M_A^2}{m_{S_1}^2} - \frac{17}{6} + \frac{M_A^2}{M_V^2}\right)\right], \quad (12)$$

where $M_V < M_A$ has been assumed and we have again neglected terms of $\mathcal{O}(m_{S_1}^2/M_{V,A}^2)$. With $m_H = m_{S_1}$, the NLO correction vanishes in the combined limit $\omega = 1$ and $M_V = M_A$, where the LO lower bound (9) is recovered.

Phenomenology.—Taking the SM reference point at $m_H = m_{S_1} = 126$ GeV, the global fit to precision electroweak data gives the results $S = 0.03 \pm 0.10$ and $T = 0.05 \pm 0.12$, with a correlation coefficient of 0.891 [13]. In Fig. 2 we show the compatibility between these "experimental" values and our NLO determinations imposing the two WSRs: Eq. (10) with $\omega = M_V^2/M_A^2$ and Eq. (11). Notice that the line with $\omega = M_V^2/M_A^2 = 1$ ($T=0$) coincides with the LO upper bound in Eq. (8), while the $\omega = M_V^2/M_A^2 \to 0$ curve reproduces the lower bound in Eq. (12) in the same limit. Thus, a vanishing scalar-Goldstone boson coupling ($\omega = 0$) would be incompatible with the data, independent of whether the second WSR has been assumed.

Figure 2 shows a very important result in the two-WSR scenario: with $m_{S_1} = 126$ GeV, the precision electroweak data require that the Higgs-like scalar should have a WW coupling very close to the SM one. At 68% (95%) C.L., one gets $\omega \in [0.97, 1]$ ([0.94,1]), in nice agreement with the present LHC evidence [1], but much more restrictive. Moreover, the vector and axial-vector states should be very heavy (and quite degenerate); one finds $M_V > 5$ TeV (4 TeV) at the 68% (95%) C.L.

This conclusion is softened when the second WSR is dropped and the lower bound in Eq. (12) is used instead. This is shown in Fig. 3, which gives the allowed 68% C.L. region in the space of parameters M_V and ω, varying M_V/M_A between 0 and 1. Note, however, that values of ω very different from the SM can only be obtained with a large splitting of the vector and axial-vector masses. In general, there is no solution for $\omega > 1.3$. Requiring $0.2 < M_V/M_A < 1$ leads to $1 - \omega < 0.4$ at 68% C.L., while the

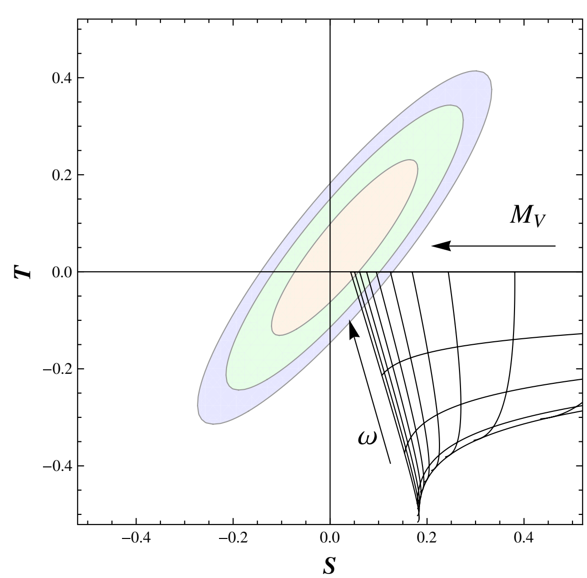

FIG. 2 (color online). NLO determinations of S and T, imposing the two WSRs. The approximately vertical curves correspond to constant values of M_V, from 1.5 to 6.0 TeV at intervals of 0.5 TeV. The approximately horizontal curves have constant values of ω: 0.00, 0.25, 0.50, 0.75, 1.00. The arrows indicate the directions of growing M_V and ω values. The shaded ellipses give the experimentally allowed regions at 68% (orange, center region), 95% (green, midregion), and 99% (blue, outer region) C.Ls.

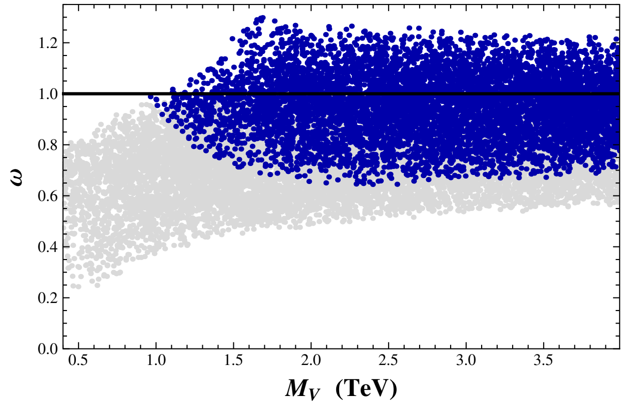

FIG. 3 (color online). Scatter plot for the 68% C.L. region, in the case when only the first WSR is assumed. The dark blue and light gray regions correspond to $0.2 < M_V/M_A < 1$ and $0.02 < M_V/M_A < 0.2$, respectively.

allowed vector mass stays above 1 TeV [14]. Taking instead $0.5 < M_V/M_A < 1$, one gets the stronger constraints $1 - \omega < 0.16$ and $M_V > 1.5$ TeV. In order to allow vector masses below the TeV scale, one needs a much larger resonance-mass splitting, so that the NLO term in Eq. (12) proportional to ω^2 compensates the growing of the LO vector contribution. The mass splitting gives also an additive contribution to T of the form $\delta T \sim \omega^2 \log(M_A^2/M_V^2)$, making lower values of ω possible for smaller M_V values. However, the limit $\omega \rightarrow 0$ can only be approached when $M_A/M_V \rightarrow \infty$.

In summary, strongly coupled electroweak models with massive resonance states are still allowed by the current experimental data. Nonetheless, the recently discovered Higgs-like boson with mass $m_{S_1} = 126$ GeV must have a WW coupling close to the SM one ($\omega = 1$). In those scenarios, such as asymptotically free theories, where the second WSR is satisfied, the S and T constraints force ω to be in the range [0.94,1] at the 95% C.L. Larger departures of the SM value can be accommodated when the second WSR does not apply, but one needs to introduce a correspondingly large mass splitting between the vector and axial-vector states.

Similar conclusions can be obtained within more specific models, particularizing our general framework. For instance, let us mention the recent phenomenological analyses of vector and axial-vector states within the $SO(5)/SO(4)$ minimal composite Higgs model [8,15]. In this context, our scalar coupling would be related to the $SO(4)$ vacuum angle θ and upper bounded in the form $\omega = \cos\theta \leq 1$ [15]. With this identification, the S and T constraints in Fig. 2 remain valid in this composite scenario.

This work was supported in part by Spanish government Grants No. FPA2007-60323, No. FPA2011-23778, No. CSD2007-00042 (Consolider Project CPAN), and MICINN-INFN No. AIC-D-2011-0818, Italian government Grant No. MIUR-PRIN-2009, National Nature Science Foundation of China Grant No. 0925522, and Universidad CEU Cardenal Herrera Grant No. PRCEU-UCH35/11. J. J. S.-C. thanks G. Cacciapaglia, H. Y. Cai, and H. Q. Zheng for useful discussions.

[1] CMS Collaboration, Phys. Lett. B **716**, 30 (2012); CMS Collaboration, CMS Physics Analysis Summaries, CERN Report No. CMS-PAS-HIG-12-045, 2012; G. Aad *et al.* (ATLAS Collaboration), Phys. Lett. B **716**, 1 (2012); CERN, ATLAS Note ATLAS-CONF-2012-170, 2012.
[2] T. Appelquist and C. Bernard, Phys. Rev. D **22**, 200 (1980).
[3] G. Ecker, J. Gasser, A. Pich, and E. de Rafael, Nucl. Phys. **B321**, 311 (1989); G. Ecker, J. Gasser, H. Leutwyler, A. Pich, and E. De Rafael, Phys. Lett. B **223**, 425 (1989); V. Cirigliano, G. Ecker, M. Eidemüller, R. Kaiser, A. Pich, and J. Portolés, Nucl. Phys. **B753**, 139 (2006).
[4] A. Pich, I. Rosell, and J. J. Sanz-Cillero, J. High Energy Phys. 08 (2012) 106.
[5] M. E. Peskin and T. Takeuchi, Phys. Rev. D **46**, 381 (1992); Phys. Rev. Lett. **65**, 964 (1990).
[6] S. Matsuzaki, R. S. Chivukula, E. H. Simmons, and M. Tanabashi, Phys. Rev. D **75**, 073002 (2007); R. Barbieri, G. Isidori, V. S. Rychkov, and E. Trincherini, Phys. Rev. D **78**, 036012 (2008); O. Cata and J. F. Kamenik, Phys. Rev. D **83**, 053010 (2011); R. Foadi and F. Sannino, Phys. Rev. D **87**, 015008 (2013).
[7] A. Orgogozo and S. Rychkov, J. High Energy Phys. 03 (2012) 046.
[8] A. Orgogozo and S. Rychkov, arXiv:1211.5543.
[9] R. Barbieri, M. Beccaria, P. Ciafaloni, G. Curci, and A. Viceré, Nucl. Phys. **B409**, 105 (1993).
[10] C. W. Bernard, A. Duncan, J. LoSecco, and S. Weinberg, Phys. Rev. D **12**, 792 (1975).
[11] S. Weinberg, Phys. Rev. Lett. **18**, 507 (1967).
[12] T. Appelquist and F. Sannino, Phys. Rev. D **59**, 067702 (1999).
[13] GFITTER, A Generic Fitter Project for HEP Model Testing, http://gfitter.desy.de/; LEP Electroweak Working Group, http://lepewwg.web.cern.ch/LEPEWWG/.
[14] A. Filipuzzi, J. Portoles, and P. Ruiz-Femenia, J. High Energy Phys. 08 (2012) 080.
[15] R. Contino, in *Physics of the Large and the Small: Tasi 2009, Proceedings of the 2009 Theoretical Advanced Study Institute in Elementary Particle Physics*, edited by C. Csaki and S. Dodelson (World Scientific, Singapore, 2011); R. Contino, D. Marzocca, D. Pappadopulo, and R. Rattazzi, J. High Energy Phys. 10 (2011) 081.