A Negative S parameter from Holographic Technicolor

Johannes Hirn$^1$ and Verónica Sanz$^2$

$^1$IFIC - Universitat de València, Edifici d’Instituts de Paterna, Apt. Correus 22085, 46071 València, Spain
$^2$Departamento de Física Teórica y del Cosmos, Universidad de Granada, Campus de Fuentenueva, 18071 Granada, Spain

We present a new class of 5D models, Holographic Technicolor, which fulfills the basic requirements for a candidate of comprehensible 4D strong dynamics at the electroweak scale. It is the first Technicolor-like model able to provide a vanishing or even negative tree-level S parameter, avoiding any no-go theorem on its sign. The model is described in the large-N regime. $S$ is therefore computable: possible corrections coming from boundary terms follow the $1/N$ suppression, and generation of fermion masses and the $S$ parameter issue do split up. We investigate the model’s 4D dual, probably walking Technicolor-like with a large anomalous dimension.

**Introduction:** The idea that electroweak symmetry breaking (EWSB) could be due to the onset of the strong-coupling regime in an asymptotically-free gauge theory was first put forward to solve the hierarchy problem in [1]. Technicolor was based on the example of massless QCD with two flavors, where the global SU(2) $\times$ SU(2) symmetry is spontaneously broken to the diagonal subgroup. A similar theory with a mass scale of order 3000 larger would feed its three GBs to the SM SU(2)$_L$ $\times$ U(1)$_Y$ gauge fields, yielding masses for the $W^\pm$ and $Z$, without an associated Higgs boson. It was however shown that a simple rescaled version of QCD fails, since it leads to the famous $S$ parameter being too large and positive as compared to the value extracted from experiments, unless the number of techni-colors is small. This last possibility is however undesirable, as it signifies the loss of our last non-perturbative handle, namely the large-N expansion.

The recent developments in Holographic QCD [4, 5, 6] give us a computable way of departing from rescaled QCD. The models of Holographic QCD aim to describe the dynamics of the QCD bound states in terms of a 5D gauge theory: the input parameters in such a description can be identified with the number of colors, the confinement scale and the condensates. The present class of models for dynamical EWSB works in a similar spirit. For the first time, the tree-level $S$ parameter is negative. This has further consequences in the gauge boson spectrum.

**Holographic Technicolor:** Our starting point is a model in five-dimensions (5D) describing electroweak symmetry breaking via boundary conditions (BCs). The extra dimension we consider here is an interval. The two ends of the space are located at $l_0$ (the UV brane) and $l_1$ (the IR brane), with the names UV/IR implying $(2\pi/l_0)$ and $(2\pi/l_1)$, respectively. We focus on metrics that can be recast as $ds^2 = w(z)^2 (\eta_{\mu\nu} dz^\mu dz^\nu - dz^2)$. We only consider the dynamics of the bulk 5D symmetry SU(2)$_L$ $\times$ SU(2)$_R$ $\times$ U(1)$_{B-L}$ gauge symmetry. As in Higgsless models [4], the BCs are chosen to break the LR symmetry to the diagonal SU(2)$_D$ on the IR brane, while the breaking on the UV brane reduces SU(2)$_R$ $\times$ U(1)$_{B-L}$ to the hypercharge subgroup. The remaining 4D gauge symmetry is $U(1)_Q$.

An important ingredient of Holographic Technicolor comes from the lessons learned in Holographic QCD: breaking on the brane is too soft to account for all phenomena found in QCD, in particular power corrections at high energies due to condensates. Besides this breaking by BCs, we therefore introduce breaking in the bulk. In the following, the bulk source of breaking will be a crossed kinetic term between L and R gauge fields, just as in [7]. (The $z$-dependence of this term could be obtained from the profile of a scalar.) At the quadratic level, this well-defined procedure may effectively be summarized as yielding different metrics, $w_A(z) \neq w_V(z)$. This bulk breaking will play an important role in our description of strong dynamics at the TeV.

**The spectrum:** In terms of physical states, no massless mode survives except for the photon. The remainder will pick up masses via the compactification. For the class of metrics that decrease away from the UV as AdS or faster (gap metrics), the massive modes can be separated into two groups: ultra-light excitations and KK-modes. If we interpret the ultra-light modes as the $W$ and $Z$, the gap suppresses the KK contributions to the electroweak observables [6]: this can be seen clearly using Sum Rules (SRs).

For any gap metric, the KK modes are repelled from the UV brane, and the massive modes approximately split into separate towers of axial and vector fields (and B fields). Thus, $W'$, the first KK mode above the $W$ would a priori be a vector (the techni-rho), while the next one, $W''$, would be an axial resonance (techni-a1), etc... One can extract SRs involving KK-mode masses (excluding ultra-light modes)

$$
\sum_{n=1}^{\infty} \frac{1}{M_{X_n}^2} \simeq \int_{l_0}^{l_1} dz w_X(z) \alpha_X(z) \int_{l_0}^{z} \frac{dz'}{w_X(z')} .
$$

where $X = V, A, B$ and $\alpha_{V,B}(z) = 1$ and $\alpha_A(z) = \int_{l_0}^{l_1} \frac{dz'}{w_A(z')}/\int_{l_0}^{l_1} \frac{dz''}{w_A(z'')}$. The SR in Eq.(1) is exact at order $O(G^0)$, where $G$ is the gap between the ultra-light mode and the heavy modes: in AdS, $w(z) = l_0/z$ and...
metrics, it can be expanded to obtain the mass of the involving both heavy \( \alpha \) decreasing with BCs

\[ D \] Goldstone boson matrix, marking the 4D gauge coupling which can be shown to agree with the expression involv-

the gap is \( G = \log(l_1/l_0) \). As in Holographic QCD, the function \( \alpha_A(z) \) \( \square \) is the wavefunction of the “would-be” Goldstone boson matrix, \( D_{\mu U}(x) \): it is monotonously decreasing with BCs \( \alpha_A(l_0) = 1 \) and \( \alpha_A(l_1) = 0 \).

On the other hand, another exact SR can be obtained, involving both heavy and ultra-light modes. For gap metrics, it can be expanded to obtain the mass of the ultra-light mode: at order \( \mathcal{O}(1/G) \), we get

\[ M_W^2 \simeq \frac{1}{f} \left( \int_{l_0}^{l_1} dz (w_V(z) + w_A(z)) \int_{l_0}^{l_1} \frac{dz'}{w_A(z')} \right) \tag{2} \]

which can be shown to agree with the expression involving the 4D gauge coupling \( g \) and techni-pion decay constant \( f \)

\[ M_W^2 = \frac{g^2 f^2}{4} + \mathcal{O}(1/G^2) = \frac{1}{G l_1^2} + \mathcal{O}(1/G^2), \tag{3} \]

as expected from Technicolor. Eq. \( \square \) shows that, at leading order \( \mathcal{O}(1/G) \), \( M_W^2, Z \) do not feel any breaking of conformation in the bulk: their mass is dominated by the UV physics. On the other hand, Eq. \( \square \) showed that the KK masses do feel the effect of this bulk breaking at leading order \( \mathcal{O}(G^0) \). Also, since their wave-functions are repelled, the KK-modes have masses that are quite in-

FIG. 1: Masses at \( \mathcal{O}(G^0) \) divided by \( l_1 \) for the lightest vector and axial KK modes of the \( W \), as a function of the condensate in their respective channel \( \alpha_{V,A} \), for \( d = 2 \).

produced at low-energy: four-fermion interactions are generated by the exchange of KK states. It can be shown that the expression of the resulting Fermi constant in terms of the techni-pion decay constant is obtained from the SM by replacing \( v \rightarrow f \). Since the model is based on an \( SU(2) \times SU(2) \) (gauge) symmetry in the bulk, broken to the diagonal subgroup by the IR BCs, it possesses custodial symmetry. This implies that the low-energy rho parameter \( \rho_{\ast}(0) \) is strictly equal to one at tree level, as was found in the deconstructed case \( \square \) and indicated by \( \square \). Also, the KK modes, being repelled from the UV brane, are insensitive to the UV BCs. The KK spectrum is therefore isospin symmetric up to \( 1/G \) corrections: \( W_n \) degenerate with \( Z_n \). In addition, since the KK contribution is small due to the large masses of the KK modes, one concludes that the \( T \) parameter is suppressed in these models. Finally, one can also show from two SRs that the \( E^4 \) and \( E^2 \) contributions to the \( W_L W_L \) scattering vanish \( \square \).

The \( S \) Parameter: The tree-level contribution to the \( S \) parameter, being a low-energy effect due to strong dynamics responsible for spontaneous symmetry-breaking, can be expressed \( \square \) in terms of the \( L_{10} \) coupling of chiral lagrangians \( S_{\text{tree}} = -16 \pi L_{10} \). The value extracted from LEP physics is \( \square \) \( S = -0.13(0.07) \pm 0.10 \) with reference Higgs mass \( m_H = 117(150) \) GeV, where the value in parentheses is the most recent analysis of data at the Z pole (2005). A sizeable negative \( L_{10} \) would easily upset the experimental constraint (note that in \( N_c = 3 \) QCD, \(-16 \pi L_{10} \sim 0.3 \)). On the other hand, large-\( N \) models of strong dynamics predict the value of \( L_{10} \) in terms of contribution of spin-1 resonances \( L_{10} = -1/4 \sum_{n=1}^{\infty} f_{V,n}^2 - f_{A,n}^2 \), via their decay constants \( f_{X,n} \) according to \( \square \), whereas other contributions are down by \( 1/N \). Higgsless models thus face a serious challenge, a no-go theorem \( \square \): \( L_{10} \) is bounded to be negative. This is readily understood by using a SR: one can translate the sum over resonance contributions into a purely geometric factor

\[ L_{10} = -\frac{N}{48 \pi^2} \int_{l_0}^{l_1} \frac{dz}{l_0} w(z) \left( 1 - \alpha(z)^2 \right), \]

where we have defined \( N/12 \pi^2 = l_0/g^2 \). The bound \( \alpha(z) \ll 1 \) implies that \( L_{10} \) is negative and proportional to the loop expansion parameter, \( N \). The most natural value for \( L_{10} \) will thus drive a large positive \( S \) parameter, excluding the simplest realization of the model. For example, pure AdS yields \( S_{\text{tree}} = N/4 \pi \).

One possibility would be to consider these models in the low-\( N \) regime. This situation is most unwelcome, as has been stressed by many authors \( \square \). The main reason is that the value of \( N \) plays an important role: it sets the range of computability of the model. Low \( N \) implies strong coupling of the gauge KK modes. A way
of putting it is via the position-dependent cutoff\[12, 16\]: a cutoff \(\Lambda\) at the position where \(w(z)\) is normalized to unity will be redshifted for processes located near a position \(z\) as \(\Lambda(z) = \Lambda \sqrt{g_{00}} = \Lambda w(z)\). For example, in pure AdS, the 5D loop expansion breaks down when \(\Lambda(z) \sim 24\pi^2 l_0/g_0^2 = 2\pi N\). The other parameter playing an important role is the gap \(G\). Reproducing the Fermi constant and the W mass implies \(NG \sim 500\). Pushing to low values of \(N\) is thus asking for a bigger separation between the W and its KK modes, which would conflict with the premise of perturbativity: strong coupling would set in before the resonances tame the high-energy behavior of amplitudes.

Returning to the large-\(N\) regime, one is then cornered to hope for miraculous cancellations. Efficient possibilities would be: introducing IR localized kinetic terms proportional to SU(2)\(_D\) or hoping for cancellations against fermion contributions.\[8\] Both possibilities face again new challenges, difficult to resolve. Trying to add large localized kinetic terms with the "wrong" sign, which are of order \(1/N\) directs again towards the low-\(N\) problem. Besides it leads to a tachyon instability.\[13\] The way out with bulk fermions poses a problem of naturalness and dangerously ties the \(S\) parameter problem with the fermion mass hierarchy, and therefore with non-oblique corrections.\[6, 17\]

Here we propose a different point of view, which arises naturally in Holographic QCD and should therefore appear in a Technicolor-like model. Local order parameters of the symmetry-breaking imply a different behavior for the \(V\) and \(A\) combinations of bulk fields.\[3, 5\] In the simplest realization of this IR behavior,\[7\] \(L_{10}\) is modified from Eq.\[14\] to read

\[
L_{10} = -\frac{N}{48\pi^2} \int_{l_0}^{\ell_1} \frac{dz}{l_0} \left( w_V(z) - w_A(z) \alpha_A(z)^2 \right),
\]

where \(w_{V,A}\) are the metrics felt by the axial and vector combinations of fields.

\(L_{10}\) is still proportional to \(N\), but the integrand in Eq.\[14\] can reverse sign for \(z\) such that \(w_A(z) \alpha(z)^2 > w_V(z)\), and \(L_{10}\) may come out positive. The first consequence is quite clear: a large-\(N\) scenario is then preferred, extending the perturbativity regime. In particular, the bulk value of \(S\) will not receive sizeable corrections from the localized kinetic terms, since these are still suppressed by \(1/N\). The \(S\) parameter is therefore computable. Another important property of the bulk \(S\) parameter is its independence on the exact IR dynamics. Contrary to the spectrum, contributions to \(S\) come mainly from the bulk far from the branes.\[10\]

We now assume that the metrics behave as AdS near the UV brane and deviate from conformality in the bulk according to

\[
w_X(z) = \frac{l_0}{z} \exp\left( \nu_X \left( \frac{z - l_0}{l_1 - l_0} \right)^{2d} \right).
\]

As explained at the beginning of the paper, this parametrically simple form encodes effects of couplings with other background fields, whose dynamics we neglect here. At order \(O(G^d)\), one can obtain an analytic expression for \(S\) in the case \(\nu_A = 0\) and \(\nu_V < 0\).

\[
S_{\text{tree}} = \frac{N}{4\pi} \left( 1 - \frac{2}{3d} (\Gamma(1) - \log(-\nu_V)) + \log(-\nu_V) + \gamma_E \right).
\]

In \(\nu_V \equiv \frac{12\pi^2/3(3+1/2)}{2^5}\sqrt{d} \omega_V\), NDA sets \(\omega_V \sim O(1)\).\[17\]

In FIG.\[2\], we show the value of \(S_{\text{tree}}/N\) for different values of \(\omega_V\) as a function of the ratio of condensates in the two channels \(\omega_A/\omega_V\), and for the pure AdS case.

A refinement in the computation of the \(S\) parameter comes from taking into account the pion loop effects and subtracting the SM value with a reference Higgs mass

\[
S = -16\pi L_{10}(\mu) + \frac{1}{12\pi} \left( \log\left( \frac{\mu^2}{m_H^2} \right) - \frac{1}{6} \right).
\]

From the understanding of the QCD case, one expects the model to predict the value of \(L_{10}(\mu)\) at the matching scale of the model with a chiral lagrangian, i.e. \(\mu \sim \text{few}/l_1 \sim \text{few TeV}\), the mass scale of the resonances. The second term in Eq.\[15\] is then positive and of order 0.1, requiring a vanishing or slightly negative \(S_{\text{tree}}\), as provided by the present model.
Four-dimensional dual: Holographic models are inspired from the AdS/CFT correspondence \cite{21}. The precise form of this conjecture relates two highly symmetric theories and is, unfortunately, far from being of direct phenomenological relevance. After a pioneering work by Pomarol \cite{21}, authors in \cite{22} explored the audacious conjecture that more realistic models like Randall-Sundrum \cite{23} would somehow inherit properties of a strongly-coupled theory like QCD provides an bottom-up models of Holographic QCD \cite{3, 4, 5, 7, 24}. The success of these models in capturing the behavior of a strongly-coupled theory like QCD provides an incentive for applications to Technicolor. In this case, one starts off on a firmer footing: in the presence of condensates, the number of (techni-)colors can be made large since it no longer in conflict with the $S$ parameter.

Let us show the effect in the 4D two-point correlator of the current $X = V, A, B$ of a metric of the form given by Eq. (1). For large euclidean $Q^2$, the two-point function for this field $X$ reads \cite{2}

$$\Pi_X (-Q^2) \simeq -\frac{N}{12\pi^2} \left( \log \left( \frac{Q^2}{\mu^2} \right) + \lambda (\mu) \right) + \frac{\langle O_{2d_X} \rangle}{Q^{2d_X}} \tag{9}$$

where the parameter $\alpha_X \equiv \langle O_{2d_X} \rangle / (N l_0^{-2d_X}) \sim O(1)$. To have a chance of obtaining a positive value for $L_{10}$, we need $(O_{2d})_V < (O_{2d})_A$. This is in agreement with Witten’s positivity condition for $\Pi_A - \Pi_V$ \cite{22}, ensuring the stability of the selected vacuum \cite{24}. Holography tells us that this bulk field $X$ is dual to some operator $\mathcal{O}$ on the 4D side with the same quantum numbers: the correlators generated by $X$ and by $\mathcal{O}$ are the same. In this particular case we see that deviations from conformality with a given power of $z^{2d}$ in Eq. (1) mimick the effects of a condensate of dimension $2d$ in the 4D dual.

Generally speaking, non-perturbative effects in QCD-like Technicolor models make them unreliable. The same goes for the case of a flat extra-dimension, the cutoff of the theory is quite low, $\Lambda \sim 2\pi N / l_1$ and quantities like the $S$ parameter are no longer computable. On the other hand, extra-dimensional models in AdS behave in a similar fashion to walking Technicolor. The warping suppresses convolutions of wave-functions, as walking kills unwanted operators. But in pure AdS, one cannot choose which operators will be suppressed: their scaling is dictated by the warping, whereas gap metrics with violations of conformality like Eq. (1) do change the scaling.

The dual of Holographic Technicolor must be a strongly-coupled theory, with the running in the UV dictated by the one of a gap metric and with non-perturbative dynamics affecting the vector and axial channel in a similar way. If the 4D dual is going to yield small or negative $S$ parameter, the net effect of condensates in the vector and axial current must go in the direction of $w_A \alpha^2 > w_V$. For example, imagine that strong dynamics generate a techni-condensate $\langle Q\bar{Q} \rangle$ responsible of breaking the Technicolor gauge group SU($N$): this condensate is represented in the 5D dual as the rescaled vev of $\langle \Phi \rangle$. Assume now that the anomalous dimensions is large, for example, due to the running mass in the 5D picture. Then, there will be a difference between the canonical dimension of $\langle Q\bar{Q} \rangle$ and the running dimension of the operator. A way of modelling this anomalous dimension would be that the vector and axial fields couple to a scalar representing the techni-quark condensate, $\Phi$, via a running mass, such that $m_{\Phi}(l_0)^2 = -3/l_0^2$ and $m_{\Phi}(l_1)^2 = d(d - 4)/l_1^2$ with $d < 3$ ($d = 2$ for extreme walking).

Conclusions: In this paper we have shown quantitatively how technicolor models which depart from rescaled QCD can exhibit a negative tree-level $S$ parameter. This was done using a holographic model (i.e. using a 5D gauge theory) for the resonances created by a strongly-interacting theory such as technicolor. It is based on the recent successes of similar 5D models for the resonances of QCD. These successes themselves validated the idea of the duality between 4D strongly-coupled theories and 5D weakly-coupled ones at the quantitative level.

We have presented the first Technicolor-like model able to provide a small $S$ parameter, and to remain computable since it is defined in the large-$N$ limit. The 5D picture shows generic features of this class of models: 1) the metric has to fall off fast near the UV to generate a gap, 2) deviations from conformality must be introduced in the bulk, describing condensates, 3) a condensate of natural size can produce the desired effect if it has dimension close to 4 (as would happen for $\alpha_{TC}(Q\bar{Q})^2$ in walking Technicolor), 4) $W'$ and $Z'$ (vector resonances) then tend to become degenerate with the $W''$ and $Z''$ (axial resonances).

In the present paper, the fermions were located for simplicity on the UV brane. As soon as we let them live in the bulk, much more interesting phenomena should arise: one big advantage of the present models is that the fermion profiles are not constrained by the requirement of cancelling the $S$ parameter contributions. The issue of the $S$ parameter is therefore decoupled from that of fermion mass generation or from $Z \rightarrow \ell\bar{\nu}$, which can be addressed in a new view \cite{25}. In particular topcolor assisted models would be implemented as in \cite{24}, Acknowledgments: We acknowledge hospitality from Boston, Harvard and Yale Universities during the completion of this work. We also thank Tom Appelquist, Tony Gherghetta, Ami Katz, Ken Lane, John March-Russell, Toni Pich and Francesco Sannino for stimulating discussions. JH is supported by the EC RTN network HPRN-CT-2002-00311 and by the Generalitat Valenciana grant GV05/015.
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