Dynamical Electroweak Symmetry Breaking

F. Sannino

The Niels Bohr Institute, Blegdamsvej 17, Copenhagen Ø, DK-2100, Denmark.

Dynamical breaking of the electroweak theory, i.e. technicolor, is an intriguing extension of the Standard Model. Recently new models have been proposed featuring walking dynamics for a very low number of techniflavors. These technicolor extensions are not ruled out by current precision measurements. Here I first motivate the idea of dynamical electroweak symmetry breaking and then summarize some of the properties of the recent models and their possible cosmological implications.

1 Introduction

The Standard Model (SM) of particle interactions is a low energy effective theory valid up to a cutoff scale Λ. One of the reasons behind the phenomenological success of the SM is that most of the physical observables depend only logarithmically on Λ. There is only one operator in the SM depending strongly on this fundamental scale, i.e. the mass squared operator of the Higgs. There are two problems associated with such an operator: 1) It is unnaturally sensitive to the scale Λ which can be taken to be the highest scale in the game, i.e. the Plank mass. 2) Even if we set the Higgs mass operator to zero, at tree level, it will be regenerated by quantum corrections.

Different resolutions of these problem have been proposed. Here I will just mention the time-honored ones: i) supersymmetric generalizations of the SM ii) dynamical breaking of the electroweak theory.

Combining the problems above with the fact that the SM alone neither accounts for the experimentally observed dark matter in the Universe nor explains why the neutron electric dipole moment is so unnaturally small we arrive at the conclusion that it is not so standard after all.

---

*Here I am barring exotic solutions to the dark matter problem stemming from the SM augmented with Gravity.*
all. Here I will focus on the Higgs sector. Perhaps the first question to ask is: Have we observed a Higgs-type mechanism in Nature?

Ordinary Superconductivity (SC) is a noble example. In the first figure, made by two slides, I summarize the key features that SC and Electroweak Symmetry Breaking (ESB) have in common. SC and ESB are both an example of a screening effect. One can define a macroscopic wave function $\psi$ in SC with $|\psi|^2 = n_c = n_s/2$, $n_c$ the number of Cooper pairs and $n_s$ the number of SC electrons. This wave function can be mapped into the Higgs wave function whose square evaluated on the ground states $|\phi|^2 = v^2/2$ sets the scale of ESM, i.e. $v$ hundreds of GeV. For the few and even fewer attentive readers who spotted the fact that the SC wave function has different units than the Higgs wave function the reason is that $\psi$ emerges in a nonrelativistic framework while $\phi$ is for a relativistic one. However differences in units disappear when comparing the Meissner static mass of the photon $M$ with the typical SM gauge boson mass $M_W$. In the figure $q = -2e$ is the electric charge of the Cooper pair and $m = 2m_e$ is its mass. One can then compare the typical screening lengths in the two cases. Having constructed the superconductive material in a lab we know that the wave function $\psi$ is not a fundamental object but rather a low energy description of something more fundamental. In the SM we do not yet know the mechanism behind the ESB but it might very well be dynamical as it is in SC. We will pursue this idea here. By simply admitting that the Higgs wave function is not associated to a fundamental field but is a low energy effective description of a more fundamental theory valid up to a scale above which we will be able to resolve its constituents one solves the first two problems mentioned in the opening of this section. The highest scale in the game is no longer Planck but rather the TeV scale.

We hence postulate the presence of a new strong force driving ESB. Earlier attempts using QCD-like technicolor have been ruled out by precision measurements. Besides, one has also to face the problem of mass generation which is provided by extended technicolor (ETC) interactions and thus leads to large flavor changing neutral currents. Recently, it has been shown that one can construct viable theories explaining the breaking of the electroweak theory dynamically while not being at odds with electroweak precision measurements. In the recently proposed theories technimatter transforms according to a higher dimensional representation of the new gauge group. By direct comparison with data it turns out that the preferred representation is the two-index symmetric. The simplest theory of this kind is a two technicolor theory. In
this case the two-index symmetric representation coincides with the adjoint\(^b\). Remarkably, these theories are already near conformal for a very small number of techniflavors. Further properties of higher dimensional representations have also been explored in\(^7\). In\(^\text{15}\) the reader can find a summary of a number of salient properties of the new technicolor theories as well as a comprehensive review of the walking properties with an exhaustive list of important references. We also note that near the conformal window\(^8\) one of the relevant electroweak parameters (\(S\)) is smaller than expected in perturbation theory. This observation is further supported by other very recent analyses\(^\text{10,11}\).

2 The Minimal Walking Model

The new dynamical sector underlying the Higgs mechanism we consider is an \(SU(2)\) technicolor gauge group with two adjoint technifermions. The theory is asymptotically free if the number of flavors \(N_f < 2.75\). The critical value of the number of flavors needed to reach the infrared fixed point value is \(N_f^c \simeq 2.075\)\(^\text{12,13}\). We expect that the theory will enter a conformal regime before the coupling rises above the critical value triggering the formation of a fermion condensate. Hence, a \(N_f = 2\) theory is sufficiently close to the critical number of flavors \(N_f^c\). This makes it a perfect candidate for a walking technicolor theory.

The two adjoint fermions may be written as

\[ T_L^a = \begin{pmatrix} U_L^a \\ D_L^a \end{pmatrix}, \quad U_R^a, \quad D_R^a, \quad a = 1, 2, 3, \quad (1) \]

with \(a\) the adjoint color index of \(SU(2)\). The left fields are arranged in three doublets of the \(SU(2)_L\) weak interactions in the standard fashion. The condensate is \(\langle \bar{U}U + \bar{D}D \rangle\) which breaks correctly the electroweak symmetry.

The model described so far suffers from the Witten topological anomaly.\(^\text{12}\) \(^\text{13}\) We can avoid this problem by adding a new weak doublet uncharged under technicolor.\(^\text{4}\) Our additional matter content resembles a copy of a SM fermion family with quarks (here transforming in the adjoint of \(SU(2)\)) and the following lepton doublet

\[ L_L = \begin{pmatrix} N_L \\ E_L \end{pmatrix}, \quad N_R, \quad E_R. \quad (2) \]

Gauge anomaly cancellations do not fix uniquely the hypercharge for the additional matter. In\(^\text{5}\) the SM-like hypercharge has been investigated in the context of an extended technicolor theory. Another interesting choice for the hypercharge has been investigated from the point of view of the electroweak precision measurements, in\(^\text{14}\)\(^\text{15}\). In that case

\[ Q(U) = 1, \quad Q(D) = 0, \quad Q(N) = -1, \quad \text{and} \quad Q(E) = -2. \quad (3) \]

Notice that with this particular hypercharge assignment, the technidown \(D\) is electrically neutral. Independently on the hypercharge assignment, after spontaneous symmetry breaking we have nine Goldstone bosons. Three will become the longitudinal components of the massive vector bosons and the other six will acquire mass via new interactions and carry technibaryon number. The low energy effective theories have been constructed in\(^\text{14}\).

A possible feature of these theories is that the resulting composite Higgs can be light with a mass of the order of 150 GeV. The phenomenology of these theories leads to interesting signatures\(^\text{13}\). It is instructive to compare the present model with the new precision measurements\(^\text{4}\). In figure\(^\text{2}\) the ellipse corresponds to the one sigma contour in the \(T-S\) plane. The black

\(^b\)The use of higher dimensional representations for walking technicolor theories was suggested first in\(^\text{6}\).
area bounded by parabolas corresponds to the region in the $T$–$S$ plane obtained when varying the Dirac masses of the two new leptons. The point at $T=0$ where the inner parabola meets the $S$ axis corresponds to the contribution due solely to the technicolor theory. The electroweak parameters are computed perturbatively. Fortunately for walking technicolor theories the non-perturbative corrections further reduce the $S$ parameter contribution and hence our estimates are expected to be rather conservative. The figure clearly shows that the walking technicolor type theories are viable models for dynamical breaking of the electroweak symmetry. The central experimental values for $S$ and $T$ are respectively $S = +0.07 \pm 0.10$ and $T = +0.13 \pm 0.10$.

2.1 The Dark-Side

According to the choice of the hypercharge there are various possibilities for providing a cold dark matter component. Here we choose the hypercharge assignment in such a way that one of the pseudo Goldstone bosons (i.e. $D^D$) does not carry electric charge. The dynamics providing masses for the pseudo Goldstone bosons may be arranged in a way that the neutral pseudo Goldstone boson is the lightest technibaryon (LTB). If conserved by ETC interactions the technibaryon number protects the lightest baryon from decaying. Since the mass of the technibaryons are of the order of the electroweak scale they may constitute interesting sources of dark matter. Some time ago in a pioneering work Nussinov suggested that, in analogy with the ordinary baryon asymmetry in the Universe, a technibaryon asymmetry is a natural possibility. A new contribution to the mass of the Universe then emerges due to the presence of the LTB. It is useful to compare the fraction of the total technibaryon mass $\Omega_{TB}$ to the total baryon mass $\Omega_B$ in the universe

$$\frac{\Omega_{TB}}{\Omega_B} = \frac{T B}{B} \frac{m_{TB}}{m_p},$$

where $m_p$ is the proton mass, $m_{TB}$ is the mass of the LTB. $TB$ and $B$ are the technibaryon and baryon number densities, respectively. In order to determine few features of our LTB particle, we made in the oversimplified approximation in which our LTB constitutes the whole dark matter of the Universe. In this limit the previous ratio should be around 5. The fact that it is charged under $SU(2)_L$ makes it detectable in Ge detectors. The basic results are shown in Fig. 3. The desired value of the dark matter fraction in the Universe can be obtained for a LTB mass of the order of a TeV for quite a wide range of values of $T^*$ (which is the temperature below which the electroweak sphaleron processes cease to be relevant). The only free parameter in our analysis is the mass of the LTB which is ultimately provided by ETC interactions.
Acknowledgments

I am delighted to thank D.D. Dietrich, N. Evans, S.B. Gudnason, D.K. Hong, S.D. Hsu, C. Kouvaris, and K. Tuominen for having shared the work and fun on which this brief summary is based. I am indebted to T. Appelquist for introducing me to this fascinating subject and M. Shifman for getting me interested in the problem of higher dimensional representations. Many scientists have contributed to this subject and deserve to be cited. I hence refer to the beautiful review of Hill and Simmons, “Strong dynamics and electroweak symmetry breaking,” Phys. Rept. 381, 235 (2003) [Erratum-ibid. 390, 553 (2004)] for a complete list of relevant references and a better description of some key features of dynamical electroweak symmetry breaking. I am supported by the Marie Curie Excellence Grant as team leader under contract MEXT-CT-2004-013510 and by the Danish Research Agency.

References

1. S. Weinberg, Phys. Rev. D 19, 1277 (1979); L. Susskind, Phys. Rev. D 20, 2619 (1979).
2. M. E. Peskin and T. Takeuchi, Phys. Rev. Lett. 65, 964 (1990).
3. F. Sannino and K. Tuominen, Phys. Rev. D 71, 051901 (2005).
4. D. D. Dietrich, F. Sannino and K. Tuominen, Phys. Rev. D 72, 055001 (2005); Phys. Rev. D 73, 037701 (2006)
5. N. Evans and F. Sannino, arXiv:hep-ph/0512080
6. K. D. Lane and E. Eichten, Phys. Lett. B 222, 274 (1989).
7. N. D. Christensen and R. Shrock, Phys. Lett. B 632, 92 (2006)
8. T. Appelquist and F. Sannino, Phys. Rev. D 59, 067702 (1999) T. Appelquist, P. S. Rodrigues da Silva and F. Sannino, Phys. Rev. D 60, 116007 (1999). Z. y. Duan, P. S. Rodrigues da Silva and F. Sannino, Nucl. Phys. B 592, 371 (2001).
9. R. Sundrum and S. D. H. Hsu, Nucl. Phys. B 391, 127 (1993)
10. D. K. Hong and H. U. Yee, arXiv:hep-ph/0602177
11. M. Harada, M. Kurachi and K. Yamawaki, Prog. Theor. Phys. 115, 765 (2006).
12. S. Nussinov, Phys. Lett. B 165, 55 (1985).
13. E. Witten, Phys. Lett. B 117, 324 (1982).
14. S. B. Gudnason, C. Kouvaris and F. Sannino, hep-ph/0603014
15. A. R. Zerwekh, arXiv:hep-ph/0512261
16. S. M. Barr, R. S. Chivukula and E. Farhi, Phys. Lett. B 241, 387 (1990).
17. J. Bagnasco, M. Dine and S. D. Thomas, Phys. Lett. B 320, 99 (1994)