# Higgs Compositeness from Top Dynamics and Extra Dimensions

Bogdan A. Dobrescu
*Fermi National Accelerator Laboratory*
*P.O. Box 500, Batavia, IL 60510, USA* *

March 18, 1999

The fundamental Higgs doublet may be replaced in the Standard Model by certain non-perturbative four-quark interactions, whose effect is to induce a composite Higgs sector responsible for electroweak symmetry breaking. A simple composite two-Higgs-doublet model is presented. The four-quark interactions arise naturally if there are either extra spatial dimensions or larger gauge symmetries at a multi-TeV scale. Some theoretical and phenomenological implications of these scenarios are discussed.

## I. MOTIVATION

The fundamental interactions observed so far in experiments are the $SU(3)_C \times SU(2)_W \times U(1)_Y$ gauge interactions acting on three generations of chiral fermions, measured at length scales larger than $10^{-16}$ cm, and the universal gravitational interactions, measured at length scales larger than about 1 mm. In addition, it is necessary that some new interactions break spontaneously the electroweak symmetry. In the Standard Model these new interactions involve a fundamental Higgs doublet, which is a convenient but rather ad-hoc assumption. It is troublesome that the Higgs mass is not controlled by gauge invariance. The same remains true for the Supersymmetric Standard Model, where the $\mu$ term is not determined by the supersymmetry breaking scale. Although this problem may be cured by the presence of a new gauge group whose breaking is linked to the transmission of supersymmetry breaking [1], it remains desirable to investigate other origins of the electroweak asymmetry of the vacuum.

It is remarkable that the electroweak symmetry may be spontaneously broken without the need for new fields. This is possible because the electroweak symmetry is embedded in the chiral symmetry of the quarks and leptons. Given that strongly coupled four-fermion interactions induce chiral symmetry breaking [2], it is possible to replace the fundamental Higgs doublet by certain four-quark operators [3,4]. However, a computation of the $W$ and $Z$ masses in the large $N_c$ limit indicates that the electroweak breaking quark mass should be of order 0.5 TeV in the absence of excessive fine-tuning. It would be tempting then to consider a fourth generation of fermions, but this is disfavored by the current electroweak data (at the 99.8% confidence level [5] in the case of degenerate fermions).

There is a simple solution to this puzzle: to introduce a seventh quark, $\chi$, whose left and right components transform as the right-handed top quark, $t_R$, under the Standard Model gauge group. In this case, there is mass mixing between $\chi$ and $t$. The electroweak breaking $\bar{\chi}_R t_L$ mass induced by four-quark interactions can be of order 0.5 TeV, while the electroweak preserving masses $\bar{\chi}_R \chi_L$ and $\bar{t}_R \chi_L$ may be chosen to yield the physical top mass at 175 GeV. This is the top condensation seesaw mechanism [6].

A consequence of this scenario is the existence of several composite scalars. If the four-quark operators involve only the $\chi_R$ and the $(t, b)_L$ doublet, then the low energy theory is the Standard Model with a heavy Higgs boson [7]. On the other hand, if the $t, b$ and $\chi$ participate in the four-quark interactions, the composite Higgs sector includes three weak doublets and three singlets [7]. Due to the mixing between the doublets, one of the neutral Higgs bosons may be light. In fact there is a second order phase transition from the electroweak asymmetric vacuum to a non-viable vacuum, in which the lightest Higgs mass vanishes. Therefore, in this case the only lower bound on the Higgs mass is set by direct searches.

*e-mail address: bdob@fnal.gov

By contrast to other mechanisms for dynamical electroweak symmetry breaking, this framework reduces to the Standard Model in a decoupling limit. For example, in the three-Higgs doublet model mentioned above, one can increase the masses of all the composite states other than the lightest Higgs boson, by increasing the scale of the four quark operators and the $\chi$ mass while keeping the VEVs close to the boundary of the second order phase transition. This guarantees that the models of this type are in agreement with the electroweak precision measurements (at least as long as the Standard Model is in agreement). A simpler model, with a composite Higgs sector involving only two doublets, which reduces to the Standard Model in the decoupling limit is presented in Section II.

These arguments show that the fundamental Higgs doublet from the Standard Model may be successfully replaced by a vectorlike quark and four-quark interactions. Such a theory is non-renormalizable, so that one has to identify a suitable origin of the four-quark operators. They may be generated by the dynamics of some spontaneously broken gauge group. The prototypical group of this sort is topcolor [8], which is an embedding of $SU(3)_C$ in $SU(3) \times SU(3)$. Models of this type are discussed in [6,7]. Other embeddings have been introduced in Ref. [9,10]. A complication of this approach is that the breaking of the additional gauge groups is non-trivial because of their strong couplings.

Another possibility is that the four-quark operators are induced by some gauge dynamics in compact spatial dimensions [11]. There are two immediate reasons that make this possibility attractive. First, from a four-dimensional point of view, the gauge bosons that have a non-zero momentum in the extra dimensions appear as massive, so that they induce four-fermion operators in the low energy theory without the need of breaking the gauge symmetry. Second, the gauge coupling is dimensionfull in more than four-dimensions such that the strength of the gauge interactions increases with the energy, giving rise to non-perturbative effects [12].

In addition, the compact spatial dimensions may provide a special bonus. To see this note that the existence of a fundamental $\chi$ quark may be seen as artificial as the Higgs sector in the Supersymmetric Standard Model because its mass is not controlled by gauge invariance. However, if the $t_R$ propagates in the extra dimensions, the four-dimensional fundamental $\chi$ may be replaced by the Kaluza-Klein excitations of $t_R$.

The possibility of electroweak symmetry breaking due to dynamics in extra dimensions is discussed in section III.

## II. A VIABLE COMPOSITE TWO-HIGGS-DOUBLET MODEL

Certain features of the models introduced in Ref. [7] may be combined to construct a minimal composite Higgs model that has the Standard Model as a decoupling limit. The only fundamental fields are the $SU(3)_C \times SU(2)_W \times U(1)_Y$ gauge bosons, the three generations of chiral fermions, and a weak-singlet vectorlike quark, $\chi$, of electric charge 2/3. Since the $\chi_L, \chi_R$ and $t_R$ transform in the same representation of the Standard Model gauge group, the Lagrangian includes the following gauge invariant mass terms

$$\mathcal{L}_{\rm mass} = -\mu_{\chi\chi}\overline{\chi}_L\chi_R - \mu_{\chi t}\overline{\chi}_L t_R + {\rm h.c.} \tag{1}$$

An attractive interaction between $\bar{\chi}_R$ and $\psi_L^3 \equiv (t,b)_L$ gives rise to a bound state with the quantum numbers of a Higgs field. In order to keep the $\psi_L^3$ field in the low energy theory, this interaction has to be non-confining. By increasing the strength of this interaction, the Higgs field becomes more deeply bound, until it develops a VEV.

The simplest interaction of this sort is a four-quark term in the Lagrangian. For breaking electroweak symmetry it is sufficient that the four-quark operator involves the $\chi_R$ and the $\psi_L^3$ [7]. In this case, the mass of the composite Higgs boson is determined by the electroweak scale, and is large, close to the upper bound allowed by unitarity and triviality. On the other hand, if the $t_R$ also participates in the four-quark interaction, then the relation between the Higgs mass and the electroweak scale disappears. This is the consequence of the mixing between the two composite Higgs doublets that are present in this situation. Therefore, the Standard Model with a Higgs boson mass constrained from below only by direct searches, may be obtained in a certain decoupling limit from a theory which includes the following four-quark operators:

$$\mathcal{L}_{\rm eff} = \frac{g^2_{\psi\chi}}{M^2}\left(\overline{\psi}^3_L\chi_R\right)\left(\overline{\chi}_R\psi^3_L\right) + \frac{g^2_{\psi t}}{M^2}\left(\overline{\psi}^3_L t_R\right)\left(\overline{t}_R\psi^3_L\right) \ . \tag{2}$$

These interactions give rise below the scale $M$ to two composite Higgs doublets: $H_t \equiv \overline{t}_R \psi_L^3$ and $H_\chi \equiv \overline{\chi}_R \psi_L^3$. The Lagrangian for the composite Higgs fields[1], valid below the scale $M$, includes kinetic terms involving the usual covariant derivative, Yukawa couplings of the scalars to their constituents, and a scalar potential:

$$\mathcal{L}_H = (D^\mu H_t^\dagger)(D_\mu H_t) + (D^\mu H_\chi^\dagger)(D_\mu H_\chi) + \left(\xi_t \overline{\psi}_L^3 t_R H_t + \xi_\chi \overline{\psi}_L^3 \chi_R H_\chi + \text{h.c.}\right) + V(H_t, H_\chi) \ . \tag{3}$$

The potential, $V$, and the Yukawa couplings, $\xi_t, \xi_\chi$, can be determined as an expansion in $1/N_c$. Although in practice the number of colors is only $N_c = 3$, there is no reason to believe that the corrections to the leading order in $1/N_c$ change the results qualitatively. Furthermore, certain important features such as the existence of a second-order phase transition are independent of the $1/N_c$ expansion [13]. In the large-$N_c$ limit, there are contributions (see Fig. 1) to the kinetic terms and to the following scalar terms:

$$V(H_t, H_\chi) = \frac{\lambda}{2}\left[(H_t^\dagger H_t)^2 + (H_\chi^\dagger H_\chi)^2 + 2|H_t^\dagger H_\chi|^2\right] + M_{H_t}^2 H_t^\dagger H_t + M_{H_\chi}^2 H_\chi^\dagger H_\chi + \mu_{\chi\chi}\mu_{\chi t}\left(H_t^\dagger H_\chi + \text{h.c.}\right) \ . \tag{4}$$


FIG. 1. Higgs self-energy and quartic couplings induced in the large-$N_c$ limit.

The quartic coupling $\lambda$, the scalar masses, $M_{H_t}, M_{H_\chi}$, and the Yukawa couplings are scale dependent parameters. Their values at a momentum scale $\mu$ are determined by the matching conditions at the scale $M$, namely that all of them blow up because the Higgs fields are no longer degrees of freedom above $M$. Also, the exchange of the Higgs fields should give rise to the four quark interactions (2), so that at the scale $M$

$$\frac{g_{\psi t}^2}{M^2} = \frac{\xi_t^2}{M_{H_t}^2} \quad , \quad \frac{g_{\psi\chi}^2}{M^2} = \frac{\xi_\chi^2}{M_{H_\chi}^2} \ . \tag{5}$$

Computing the diagrams of Fig. 1 with a momentum cut-off $M$, gives for $\mu \ll M$ the following results [7]:

$$\xi_t = \xi_\chi = \sqrt{\frac{\lambda}{2}} = \frac{4\pi}{\sqrt{N_c \ln\left(M^2/\mu^2\right)}} \quad , \tag{6}$$

$$M_{H_{t,\chi}}^2 = \frac{2M^2}{\ln\left(M^2/\mu^2\right)}\left(\frac{8\pi^2}{N_c g_{\psi t,\chi}^2} - 1\right) + 2\mu_{\chi\chi}^2 \ . \tag{7}$$

From the expression for the masses it follows that as the four-quark couplings $g_{\psi t,\chi}^2$ are varied, the vacuum suffers a second order phase transition (or at least a weakly first order one). Note that in addition to the terms displayed in Eq. (4), the potential includes higher dimensional terms and sub-leading terms in $\mu/M$, which may be neglected. The

---

[1]Some related composite two-Higgs-doublet models have been studied in [15].

$H_t^\dagger H_\chi$ term is appropriately included in $V$ because for $g^2_{\psi t,\chi}$ close to the critical value, the $|M^2_{H_{t,\chi}}|$ squared masses may be comparable or smaller than $\mu_{\chi\chi}\mu_{\chi t}$. In the presence of this term, the two doublets acquire aligned VEVs for a range of $g^2_{\psi t,\chi}$:

$$\langle H_t \rangle = \frac{1}{\sqrt{2}} \begin{pmatrix} v\cos\beta \\ 0 \end{pmatrix} \quad , \quad \langle H_\chi \rangle = \frac{1}{\sqrt{2}} \begin{pmatrix} v\sin\beta \\ 0 \end{pmatrix} . \tag{8}$$

where the electroweak scale, $v \approx 246$ GeV, and $\tan\beta$ are determined from the minimization conditions as functions of $M^2_{H_{t,\chi}}/\lambda$ and $\mu_{\chi\chi}\mu_{\chi t}/\lambda$. The top quark mass is given by the smaller eigenvalue of the $t-\chi$ mass matrix:

$$\begin{pmatrix} \frac{1}{\sqrt{2}}\xi_t v\cos\beta & \frac{1}{\sqrt{2}}\xi_\chi v\sin\beta \\ \mu_{\chi t} & \mu_{\chi\chi} \end{pmatrix} \tag{9}$$

For example, when $\tan\beta \gg 1$ the top mass is given by a seesaw relation:

$$m_t \approx \frac{\xi_\chi v}{\sqrt{2}} \frac{\mu_{\chi t}}{\mu_{\chi\chi}} . \tag{10}$$

Note that $\xi_\chi v/\sqrt{2} \sim 0.5$ TeV, as mentioned in Section I. The lightest CP-even neutral Higgs boson has a squared-mass

$$M^2_{h^0} = \lambda v^2 + M^2_{H_t} + M^2_{H_\chi} - \left[ \left( \frac{\lambda v^2}{2}\cos 2\beta + M^2_{H_t} - M^2_{H_\chi} \right)^2 + \left( \frac{\lambda v^2}{2}\sin 2\beta + 2\mu_{\chi t}\mu_{\chi\chi} \right)^2 \right]^{1/2} . \tag{11}$$

Consider the following decoupling limit

$$M_{H_t}, M_{H_\chi}, \mu_{\chi\chi} \gg \mu \sim v . \tag{12}$$

This limit implies a fine-tuning of the mass parameters in the potential such that the two VEVs remain fixed when the masses increase. Eqs. (6) and (7) remain valid provided $M \gg \mu^2_{\chi\chi}/\mu$. It is clear that a tuning of $M_{H_{t,\chi}}$, $\mu_{\chi\chi}$ and $\mu_{\chi t}$ allows the $h^0$ to remain light, while the other physical scalar masses decouple. Therefore, in this decoupling limit the composite Higgs sector reduces to the Standard Model Higgs. In actuality it is not necessary to tune the parameters more finely than a few percent in order to keep the corrections to the electroweak observables in agreement with the data (the main constraint comes from the $\rho$ parameter [6,7,14]).

With the electroweak symmetry broken correctly and the top quark mass accommodated as shown above, it remains to produce the masses and mixings of the other quarks and leptons. In the effective theory below the scale $M$, these are easily generated by perturbative four-fermion operators which induce Standard Model Yukawa couplings. For example, the operators

$$\frac{1}{M^2} \left(\overline{\chi}_R \psi_L^3\right) \left(\overline{l}_L^j i\sigma_2 e_R^k\right) \tag{13}$$

induce Yukawa couplings between $H_\chi$ and the charged lepton fields, $l_L^j$ and $e_R^i$, ($j,k$ are generational indices). These four-fermion operators and the non-perturbative four-quark operators that produce the Higgs bound states have to originate in physics above the scale $M$. Another possibility for fermion mass generation is to extend the seesaw mechanism to all the quarks and leptons [9].

## III. GAUGE DYNAMICS IN COMPACT SPATIAL DIMENSIONS

The non-perturbative four-quark operators introduced in Section II may arise naturally in the presence of compact spatial dimensions [11]. Although this is a rather general statement, it is instructive to discuss it from the perspective of string theory.

String (or M) theory predicts the existence of 16 or 32 conserved supercharges, and six (seven) extra spatial dimensions at the string scale $M_s$. It has been traditionally assumed that i) $M_s$ is close to the Planck scale, ii) the extra dimensions are compactified at $M_s$, and iii) $N = 1$ supersymmetry is preserved all the way down to the electroweak scale. All these assumptions are consistent with a perturbative string picture. However, the recent progress in understanding non-perturbative string dynamics raises questions about these three assumptions.

First, it is natural that the string scale $M_s$ is lowered down to the GUT scale [16], and in fact it can be placed anywhere below the Planck scale, with a lower bound set by phenomenology in the TeV range [17]. A remarkable realization of this idea [18] is based on the observations that there may be compact dimensions as large as 1 mm provided they are accessible only to gravitons, and that the scale where gravity becomes strong is suppressed by the volume of the compact space. The restriction of matter and gauge fields to a surface in extra dimensions is permitted in string and M theory by the existence of branes.

Second, extra dimensions accessible to both gravitons and gauge bosons may also have compactification scales as small as a few TeV [19], and in principle could be somewhat larger than $M_s^{-1}$.

Finally, non-perturbative string effects may break supersymmetry completely at the $M_s$ scale. For example, a general manifold does not preserve any supercharge, so that if some of the spatial dimensions have the compactification scale at $M_s$, the low energy theory may be a non-supersymmetric field theory.

Of course, string theory is still plagued with phenomenological disasters, such as the cosmological constant problem, the existence of potentially massless moduli, or the inability of predicting the Standard Model at low energy. However, these problems might be solved in the future. Until the non-perturbative string effects will be better understood, it is useful to adopt a phenomenological approach, and to investigate various scenarios for physics beyond the Standard Model inspired by the above considerations regarding string theory.

A commonly used assumption for allowing chiral fermions in the four-dimensional theory obtained upon dimensional reduction, is that a quantum field theory may be defined on an orbifold, which is a space with singularities. This assumption, made a long time ago [19] and reinvigorated by the $S^1/Z_2$ compactification of M theory [20], has been used to study supersymmetry breaking [21], gauge coupling unification [22], and other phenomenological and theoretical issues [23,24] regarding extra dimensions. It is not clear whether the string dynamics may be treated below $M_s$ as pure quantum field theory in a higher dimensional space, especially since the separation between the compactification scale and $M_s$ is not large, but one can assume that this is the case and investigate the consequences.

The above considerations have a striking connection with the composite Higgs models: the required four-quark interactions may be induced by gauge dynamics in extra dimensions. The physical picture employed here is that the eleven-dimensional spacetime of M theory includes three flat spatial dimensions and seven compact dimensions with a sufficiently large volume to allow $M_s$ much below the Planck scale. The gauge fields are restricted to a region of this space which includes the three-dimensional flat space and has a thickness larger than $M_s^{-1}$ in $\delta$ extra dimensions. In this case, the gluons that have zero momentum in the extra dimensions correspond to the massless QCD gluons, while the ones with non-zero momentum in the extra dimensions are massive gauge bosons from a four-dimensional point of view, with a spectrum given by

$$M^2_{n_1,\ldots,n_\delta} = \sum_{l=1}^{\delta} \frac{n_l^2}{R_l^2} \ , \tag{14}$$

where $R_l$ are the compactification radii of the $\delta$ extra dimensions, and $n_l$ are the Kaluza-Klein (KK) excitation numbers. If the quarks propagate only on the three-dimensional boundary of the $3+\delta$ space accessible to the gluons (this corresponds to the fixed point of an orbifold), then the couplings of the KK modes of the gauge bosons to the quarks and leptons are identical (up to an overall normalization) with those of the Standard Model gauge bosons. Therefore, the tree level exchange of the gluonic KK modes induces flavor universal four-quark operators in the low energy theory:

$$\mathcal{L}^c_{\rm eff} = -\sum_{l=1}^{\delta} \sum_{n_l \ge 0} \frac{g_s^2}{2M^2_{n_1,\ldots,n_\delta}} \left( \sum_q \overline{q} \gamma_\mu T^a q \right)^2 \ , \tag{15}$$

where $q$ are all the quark fields, $T^a$ are the $SU(3)_{\rm C}$ generators, and $g_s$ is the QCD gauge coupling. The sum over KK modes should be cut off at modes of mass $\sim M_s$. The left-right current-current part of $\mathcal{L}^c_{\rm eff}$ has the same form (due to a Firz transformation and the large-$N_c$ limit) as the operators that produce scalar bound states, while the other parts of $\mathcal{L}^c_{\rm eff}$ do not contribute in the large-$N_c$ limit to the scalar potential.

It is clear that for a sufficiently large number of extra dimensions, the number of KK modes, $N_{KK}$, grows to the point where the attractive interaction between the quarks is strong enough to trigger electroweak symmetry breaking. In practice, the sum (15) is super-critical for $\delta = 4$ dimensions of size $\sim 2/M_s$. This result [11] is derived at tree level, so that it may be modified when loops are included. However, the qualitative picture appears to be robust.

Since the gluon couplings are flavor universal, all the quarks participate as constituents in scalar bound states. The VEVs of these scalars may be controlled by additional interactions. For example, the KK modes of the hypercharge gauge boson give rise to four-quark interactions which are attractive for the up-type quarks and repulsive for the down-type ones. Therefore, the scalars made up of up-type quarks are the most strongly bound and they will be the only ones with non-zero VEVs provided the flavor universal interactions are close to criticality. Furthermore, some flavor non-universal interactions should prevent the bound states involving the up and charm quarks from developing large VEVs. Such interactions, as well as the perturbative four-fermion terms that should induce the light quark and lepton masses, could be accommodated by various flavor physics scenarios above the compactification scale [23].

It is convenient to assume that $M_s$ corresponds to the scale where the perturbative expansion in $\alpha_s N_{KK}$ breaks down, which is usually larger than the compactification scale by less than one order of magnitude. Given that the compactification scale corresponds to the scale $M$ of the four-quark operators, which is most likely in the $10-50$ TeV range, the string scale is expected to be $M_s \sim \mathcal{O}(100)$ TeV. The compact space accessible to the gravitons has in this case a seven-dimensional volume of order $(10 \text{ GeV})^{-7}$.

The framework discussed here can be extended in different directions. For example, instead of using the gluonic KK modes for binding the composite scalar states, one can use some new gauge symmetries. Most economical would be an anomalous $U(1)$ which leads to the formation of only two Higgs doublets, as in Section II. Another interesting possibility is to allow the $t_R$ to propagate in extra dimensions such that its KK modes play the role of $\chi$. Some of the above arguments change in this case because the couplings of the gluonic KK modes are modified.

In conclusion, a composite Higgs sector may be provided by quark-antiquark states bound by gluons propagating in compact dimensions. If the Higgs sector will be discovered in collider experiments, then the study of the spectrum will allow to distinguish the origin of the attractive four-quark interactions: new gauge bosons, or extra dimensions.

*Acknowledgements:* I would like to thank Chris Hill and Nima Arkani-Hamed for illuminating conversations, and the organizers of DPF'99 for a stimulating environment. Fermilab is operated by the URA under DOE contract DE-AC02-76CH03000.

---

[1] H. C. Cheng, B. A. Dobrescu, and K. T. Matchev, Phys. Lett. **B439**, 301 (1998), hep-ph/9807246, and hep-ph/9811316.
[2] Y. Nambu and G. Jona-Lasinio, Phys. Rev. **122**, 345 (1961).
[3] Y. Nambu, in the Proceedings of the *XI Warsaw Symposium on Elementary Particle Physics, May 1988*, ed. Z. Ajduk, *et al* (World Scientific, 1989); in the Proceedings of the *1988 International Workshop on New Trends in Strong Coupling Gauge Theories*, Nagoya, Japan, ed. Bando, Muta and Yamawaki (World Scientific, 1989); report EFI-89-08 (1989); V. A. Miransky, M. Tanabashi and K. Yamawaki, Mod. Phys. Lett. **A4**, 1043 (1989); Phys. Lett. **B221**, 177 (1989); W. J. Marciano, Phys. Rev. Lett. **62**, 2793 (1989).
[4] W.A. Bardeen, C.T. Hill and M. Lindner, Phys. Rev. **D41**, 1647 (1990).
[5] J. Erler and P. Langacker, hep-ph/9809352.
[6] B. A. Dobrescu and C. T. Hill, Phys. Rev. Lett. **81**, 2634 (1998), hep-ph/9712319.
[7] R. S. Chivukula, B. A. Dobrescu, H. Georgi, and C. T. Hill, Phys. Rev. **D59**, 075003 (1999), hep-ph/9809470.
[8] C. T. Hill, Phys. Lett. **B266**, 419 (1991).
[9] G. Burdman and N. Evans, hep-ph/9811357.

[10] M. Lindner and G. Triantaphyllou, Phys. Lett. **B430**, 303 (1998), hep-ph/9803383;
G. Triantaphyllou, hep-ph/9811250.
[11] B. A. Dobrescu, hep-ph/9812349.
[12] N. Arkani-Hamed and S. Dimopoulos, hep-ph/9811353.
[13] R. S. Chivukula and H. Georgi, Phys. Rev. **D58**, 075004 (1998), hep-ph/9805478.
[14] R. S. Chivukula, B. A. Dobrescu and J. Terning, Phys. Lett. **B353**, 289 (1995), hep-ph/9503203.
[15] M. Suzuki, Phys. Rev. **D41**, 3457 (1990);
M. A. Luty, Phys. Rev. **D41**, 2893 (1990).
[16] E. Witten, Nucl. Phys. **B471**, 135 (1996), hep-th/9602070.
[17] J. Lykken, Phys. Rev. **D54**, 3693 (1996), hep-th/9603133.
[18] N. Arkani-Hamed, S. Dimopoulos and G. Dvali, Phys. Lett. **B429**, 263 (1998), hep-ph/9803315.
[19] I. Antoniadis, Phys. Lett. **B246**, 377 (1990);
V. A. Kostelecky and S. Samuel, Phys. Lett. **B270**, 21 (1991);
I. Antoniadis and K. Benakli, Phys. Lett. **B326**, 69 (1994), hep-th/9310151;
I. Antoniadis, K. Benakli and M. Quiros, Phys. Lett. **B331**, 313 (1994), hep-ph/9403290.
[20] P. Horava and E. Witten, Nucl. Phys. **B475**, 94 (1996), hep-th/9603142; Nucl. Phys. **B460**, 506 (1996), hep-th/9510209.
[21] E. A. Mirabelli and M. E. Peskin, Phys. Rev. **D58**, 065002 (1998), hep-th/9712214;
L. Randall and R. Sundrum, hep-th/9810155.
[22] K. R. Dienes, E. Dudas, and T. Ghergheta, Phys. Lett. **B436**, 55 (1998), hep-ph/9803466;
D. Ghilencea and G. G. Ross, hep-ph/9809217;
S. A. Abel and S. F. King, hep-ph/9809467;
C. D. Carone, hep-ph/9902407.
[23] K. R. Dienes, E. Dudas, T. Ghergheta, Nucl. Phys. **B537**, 47 (1999), hep-ph/9806292, and hep-ph/9811428;
Z. Berezhiani and G. Dvali, hep-ph/9811378;
N. Arkani-Hamed, S. Dimopoulos, G. Dvali and J. March-Russell, hep-ph/9811448;
Z. Kakushadze, hep-th/9812163 and hep-th/9902080;
A. E. Faraggi and M. Pospelov, hep-ph/9901299.
[24] I. Antoniadis, N. Arkani-Hamed, S. Dimopoulos and G. Dvali, Phys. Lett. **B436**, 257 (1998), hep-ph/9804398;
H. Hatanaka, T. Inami and C. S. Lim, hep-th/9805067;
G. Shiu and S.-H. H. Tye, Phys. Rev. **D58**, 106007 (1998), hep-th/9805157;
A. Pomarol and M. Quiros, hep-ph/9806263;
Z. Kakushadze and S.-H. H. Tye, hep-th/9809147;
K. Benakli, hep-ph/9809582;
I. Antoniadis, S. Dimopoulos, A. Pomarol, and M. Quiros, hep-ph/9810410;
Z. Kakushadze, hep-th/9811193;
T. Kobayashi, J. kubo, M. Mondragon and G. Zoupanos, hep-ph/9812221;
A. Delgado, A. Pomarol, and M. Quiros, hep-ph/9812489;
A. Donini and S. Rigolin, hep-ph/9901443;
Y. Kawamura, hep-ph/9902423;
M. Masip and A. Pomarol, hep-ph/9902467;
P. Nath and M. Yamaguchi, hep-ph/9902323 and hep-ph/9903298.