SUSY breaking scales in the gauge-Higgs unification

Hisaki Hatanaka and Yutaka Hosotani

Department of Physics, Osaka University, Toyonaka, Osaka 560-0043, Japan

(Dated: 23 May 2012)

Abstract

In the $SO(5) \times U(1)$ gauge-Higgs unification in the Randall-Sundrum (RS) warped space the Higgs boson naturally becomes stable. The model is consistent with the current collider signatures for a large warp factor $z_L > 10^{15}$ of the RS space. In order for stable Higgs bosons to explain the dark matter of the Universe the Higgs boson must have a mass $m_h = 70 \sim 75\,\text{GeV}$, which can be obtained in the non-SUSY model with $z_L \sim 10^5$. We show that this discrepancy is resolved in supersymmetric gauge-Higgs unification where a stop mass is about $300 \sim 320\,\text{GeV}$ and gauginos in the electroweak sector are light.

PACS numbers: 12.60.-i, 11.10.Kk, 11.30.Pb
The Higgs boson, necessary for inducing spontaneous symmetry breaking in the standard model (SM) of electroweak interactions, is yet to be discovered. It is not clear at all if the Higgs boson appears as described in the SM. New physics may be hiding behind it, the Higgs boson having properties quite different from those in the SM.

In the gauge-Higgs unification scenario the 4D Higgs boson becomes a part of the extra-dimensional component of gauge fields.[1]-[4] Many models have been proposed with predictions to be tested at colliders. Among them the $SO(5) \times U(1)$ gauge-Higgs unification in the Randall-Sundrum (RS) warped space is most promising.[5]-[14]

One of the most striking results in the model is that the 4D Higgs boson naturally becomes stable.[10] The vacuum expectation value of the Higgs boson corresponds to an Aharonov-Bohm (AB) phase $\theta_H$ in the fifth dimension. With bulk fermions introduced in the vector representation of $SO(5)$ the value of $\theta_H$ is dynamically determined to be $\frac{1}{2}\pi$, at which the Higgs boson becomes stable while giving masses to quarks, leptons, and weak bosons. There emerges $H$ parity ($P_H$) invariance at $\theta_H = \frac{1}{2}\pi$. All particles in the SM other than the Higgs boson are $P_H$-even, while the only $P_H$-odd particle at low energies is the Higgs boson, which in turn guarantees the stability of the Higgs boson.[12, 14] As a consequence the Higgs boson cannot be seen in the current collider experiments, since all experiments so far are designed to find decay products of the Higgs boson.

The model has one parameter to be determined, namely the warp factor $z_L$ of the RS spacetime. With $z_L$ given, the mass of the Higgs boson $m_h$ is predicted. It is found that $m_h = 72, 108$ and $135$ GeV for $z_L = 10^5, 10^{10}$ and $10^{15}$, respectively. We note that the LEP2 bound, $m_h > 114$ GeV, is evaded as the $ZZH$ coupling exactly vanishes as a result of the $P_H$ invariance.

There appears slight deviation in the gauge couplings of quarks and leptons from those in the SM. It turns out that the gauge-Higgs unification model gives a better fit to the forward-backward asymmetries in $e^+e^-$ collisions on the $Z$ pole than the SM. However, the branching fractions of $Z$ decay are fit well only for $z_L \gtrsim 10^{15}$. The gauge-Higgs unification model gives predictions for Kaluza-Klein (KK) excitation modes of various particles. In particular, the first KK $Z$ has a mass 1130 GeV and a width 422 GeV for $z_L = 10^{15}$. The current limit on the $Z'$ production at the Tevatron and LHC indicates $z_L > 10^{15}$. All of the collider data prefer a large warp factor in the gauge-Higgs unification model.[13] These analyses have been done at the tree level so far.
The fact that Higgs bosons become stable leads to another important consequence. They become the dark matter of the Universe. It has been shown that in order for stable Higgs bosons to account for the entire dark matter of the Universe observed by WMAP, \( m_h \) must be in the range \( 70 \sim 75 \) GeV, smaller than the \( W \) boson mass \( m_W \). If \( m_h > m_W \), the relic abundance of Higgs bosons would become very small. To have \( m_h = 70 \sim 75 \) GeV in the gauge-Higgs unification model we need \( z_L \sim 10^5 \), which is in conflict with the collider data.

Of course nothing is wrong with \( m_h \sim 135 \) GeV. It simply implies that Higgs bosons account for a tiny fraction of the dark matter of the Universe. Yet it is curious and fruitful to ask if there is a natural way in the gauge-Higgs unification scenario to satisfy the two requirements; (i) to be consistent with the collider data, and (ii) to explain the entire dark matter of the Universe.

In this paper we would like to show that the two requirements are naturally fulfilled if the model has softly broken supersymmetry (SUSY) such that SUSY partners of observed particles acquiring large masses. It will be found that the mass of a stop \( \tilde{t} \), a SUSY partner of a top quark \( t \), needs to be \( 300 \sim 320 \) GeV, when SUSY partners of \( W, Z \) and \( \gamma \) are light.

The key observation is that the nonvanishing Higgs boson mass \( m_h \) in the gauge-Higgs unification arises at the quantum-level, whereas the dominant part of collider experiments is governed by the structure at the tree-level. If SUSY is exact, the contributions of bosons and fermions to the effective potential \( V_{\text{eff}}(\theta_H) \) cancel so that \( V_{\text{eff}}(\theta_H) = 0 \), the Higgs boson remaining massless. As SUSY is broken, the cancellation becomes incomplete. If the SUSY breaking scale is much larger than the KK mass scale, the model is reduced to the non-SUSY model. In particular, \( m_h \) becomes \( \sim 135 \) GeV for \( z_L = 10^{15} \). Put differently, one can ask how large the SUSY breaking scale should be to have \( m_h = 70 \sim 75 \) GeV with \( z_L = 10^{15} \sim 10^{17} \) so that the relic abundance of Higgs bosons saturate the dark matter of the Universe.

The RS warped spacetime is given by \( ds^2 = e^{-2k y} dx_\mu dx^\mu + dy^2 \) for \( 0 \leq y \leq L \). The AdS curvature in \( 0 < y < L \) is \( -6k^2 \). The warp factor is \( z_L = e^{kL} \). In the \( SO(5) \times U(1) \) gauge-Higgs unification there appears an AB phase, or the Wilson line phase \( \theta_H \), in the fifth dimension, as the RS spacetime has topology of \( R^4 \times (S^1/Z_2) \). The 4D Higgs field appears as a zero mode in the \( SO(5)/SO(4) \) part of the fifth-dimensional component of the vector potential \( A_y(x,y) \). \( \langle A_y \rangle = \langle A_y^3/T^4 \rangle \neq 0 \) when the EW symmetry is spontaneously broken,
where the $SO(5)$ generator $T^4$ is defined by $(T^4)_{ab} = (i/\sqrt{2})(\delta_{a5}\delta_{b4} - \delta_{a4}\delta_{b5})$. The Wilson line phase $\theta_H$ is given by $\exp \{i\theta_H \sqrt{2} T^4 \} = P \exp \{ i g_A \int_0^L dy \langle A_y \rangle \}$.

The effective potential $V_{\text{eff}}(\theta_H)$ at the one-loop level is determined by the mass spectrum $\{m_n(\theta_H)\}$ in the presence of the phase $\theta_H \neq 0$. It is given in $d$ dimensions, after Wick rotation, by

$$V_{\text{eff}} = \pm \frac{1}{2} \int \frac{d^d p}{(2\pi)^d} \sum_n \ln(p^2 + m_n^2)$$

$$= \pm \frac{\Gamma\left(-\frac{1}{2}d\right)}{2(4\pi)^{d/2}} \sum_n m_n^d.$$  \hspace{1cm} (1)

The upper (lower) sign corresponds to bosons (fermions). The second equality is understood by analytic continuation for large $\text{Re} d$ in the complex $d$-plane.

In supersymmetric theory with SUSY breaking each KK tower with a spectrum $\{m_n\}$ is accompanied by its SUSY partner with a spectrum $\{\overline{m}_n\}$. With a SUSY breaking scale $\Lambda$ the latter is well mimicked by

$$\overline{m}_n = \sqrt{m_n^2 + \Lambda^2}$$  \hspace{1cm} (2)

which has a property that $\overline{m}_n \sim m_n$ for $m_n \gg \Lambda$. Suppose that the spectrum $\{m_n\}$ ($m_n > 0$) is determined by the zeros of an analytic function $\rho(z)$; $\rho(m_n) = 0$. Then the spectrum $\{\overline{m}_n\}$ is determined by the zeros of $\overline{\rho}(z) = \rho(\sqrt{z^2 - \Lambda^2})$. When $\rho(iy) = \rho(-iy)$ for real $y$ and $|\ln \rho| < |z|^q$ with some $q$ for $|z| \to \infty$, a convenient formula for $V_{\text{eff}}$ in [1] has been derived. In the present case the function $\overline{\rho}(z)$ has a branch cut between $\Lambda$ and $-\Lambda$ in the $z$-plane so that elaboration of the argument there is necessary.

In the dimensional regularization, [1] with $\{\overline{m}_n\}$ is transformed into

$$V_{\text{eff}} = \pm \frac{\Gamma\left(1 - \frac{1}{2}d\right)}{2\pi i (4\pi)^{d/2}} \int_C dz z^{d-1} \ln \overline{\rho}(z).$$  \hspace{1cm} (3)

Here the contour $C$ encircles the zeros of $\overline{\rho}(z)$, $\{\overline{m}_n\}$, clockwise. The contour is transformed, for sufficiently small $\text{Re} d$, to $C'$ which runs from $-i\infty$ to $+i\infty$, avoiding the cut, as shown in fig. 1. With $C'$ the formula is analytically continued to a larger $\text{Re} d$. The contribution coming from an infinitesimally small circle ($C_\epsilon$) around the branch point at $\Lambda$ vanishes. The integral just below the cut ($C_-^{\text{cut}}$) cancels the one just above the cut ($C_+^{\text{cut}}$) as $\rho(iy) = \rho(-iy)$.

The contributions coming from the integrals along the imaginary axis combine to give the expression involving an integral $\int_0^\infty dy y^{d-1} \ln \rho(i\sqrt{y^2 + \Lambda^2})$. Combining contributions from
the KK tower with a spectrum \( \{m_n\} \) and its SUSY partner with \( \{m'_n\} \), one finds that

\[
V_{\text{eff}} = \frac{\pm 1}{(4\pi)^{d/2}\Gamma(\frac{1}{2}d)} \int_0^\infty dy y^{d-1} \ln \frac{\rho(iy)}{\rho(i\sqrt{y^2 + \Lambda^2})}.
\]  

\[ (4) \]

Previously \( V_{\text{eff}}(\theta_H) \) in the gauge-Higgs unification in the RS spacetime has been evaluated by making use of the formula with only \( y^{d-1} \ln \rho(iy) \) in the integrand. \[8, 12, 17, 18\]

\[ \text{FIG. 1: Contours } C \text{ and } C' \text{ in the expression (3).} \]

The \( SO(5) \times U(1) \) gauge-Higgs unification model has been specified in Refs. \[8, 9\]. In the bulk five-dimensional spacetime, in addition to the \( SO(5) \) and \( U(1) \) gauge fields, four bulk fermion multiplets in the vector representation of \( SO(5) \) are introduced for each generation of quarks and leptons. On the Planck brane at \( y = 0 \), right-handed brane fermions \( \hat{\chi}^\alpha_R \) and one brane scalar \( \hat{\Phi} \) are introduced. The orbifold boundary conditions break \( SO(5) \) to \( SO(4) \cong SU(2)_1 \times SU(2)_2 \). The nonvanishing vacuum expectation value \( \langle \hat{\Phi} \rangle \) spontaneously breaks \( SO(4) \times U(1) \) to \( SU(2)_1 \times U(1)_1 \), and at the same time give large masses of \( O(m_{KK}) \) to exotic fermions. The resultant fermion spectrum at low energies is the same as in the SM. The \( SO(4) \times U(1) \) gauge anomalies are canceled.

The relevant contributions to \( V_{\text{eff}}(\theta_H) \) come from the \( W \) and \( Z \) towers of the four-dimensional components \( A_\mu(x,y) \), the Nambu-Goldstone towers of the fifth-dimensional components \( A_y(x,y) \), and the top quark tower. Contributions coming from other quark and lepton towers are negligible. \[8, 12\] Brane fields give no contribution to \( V_{\text{eff}}(\theta_H) \), as they do not couple to \( A_y \). We recall the way the quarks and leptons acquire masses is different from that in SM. \( A_y \) connects the left-handed and right-handed components of the up-type quarks and charged leptons directly, whereas those of the down-type quarks and neutrinos are intertwined through both gauge couplings and additional interactions with brane
fermions and scalars. Thus the effective Higgs couplings of the down-type quarks and neutrinos appear after integrating heavy brane fermions. It is notable that only the ratio of two large mass scales of the brane fermion couplings appears in the effective Higgs couplings.\cite{8, 9}

We comment that in the supersymmetric extension of the model two brane scalar fields, \( \hat{\Phi}_u \) and \( \hat{\Phi}_d \), need to be introduced. Further in 5D SUSY there appear 4D scalar fields, associated to the zero mode of \( A_\mu \), to form a 5D \( N = 1 \) (4D \( N = 2 \)) vector multiplet. In this paper we assume that such scalar fields acquire large SUSY breaking masses, giving little effect on the Wilson-line dynamics.

In the supersymmetric extension two SUSY breaking scales become important for \( V_{\text{eff}}(\theta_H) \): \( \Lambda_{\text{gh}} \) for the super partners of the \( W, Z \) and Nambu-Goldstone towers, and \( \Lambda_{\text{stop}} \) for super partner of the \( t \) quark tower. The stop (\( \tilde{t} \)) mass is given by \( m_{\tilde{t}} = \sqrt{m_t^2 + \Lambda_{\text{stop}}^2} \).

There arises no constraint to the masses of other squarks and sleptons. The masses of gluinos do not affect \( V_{\text{eff}}(\theta_H) \), being irrelevant in the present analysis.

It is most convenient to express the function \( \rho(z) \) in \( (4) \) in the form
\[
\rho(iy) = 1 + \frac{Q(q)}{1 + \frac{Q(q)}{\Lambda_{\text{gh}}}}\frac{1 + Q(q)}{\Lambda_{\text{gh}}}.
\]
where \( \Lambda \) is related to the SUSY breaking scale by \( \Lambda = k z_L^{-1} \tilde{\Lambda} \). We define
\[
Q_0(q; c) = \frac{z_L^{1/2}}{q^2 F_{c-1/2}(q) F_{c+1/2}(q)} ,
\]
\[
Q_0(q; c) = \frac{z_L}{(1/2) F_{c-1/2}(q) F_{c+1/2}(q)} ,
\]
\[
F_\alpha(q) = I_\alpha(q z_L^{-1}) K_\alpha(q) - K_\alpha(q z_L^{-1}) I_\alpha(q) \]
where \( I_\alpha(q) \) and \( K_\alpha(q) \) are the modified Bessel functions. Then \( V_{\text{eff}}(\theta_H) \) in the model is given by
\[
V_{\text{eff}}(\theta_H) \simeq 4 I \left[ \frac{1}{2} Q_0(q; 1/2), \tilde{\Lambda}_{\text{gh}} \right] + 2 I \left[ \frac{1}{2 \cos^2 \theta_W} Q_0(q; 1/2), \tilde{\Lambda}_{\text{gh}} \right] \\
+ 3 I \left[ Q_0(q; 1/2), \tilde{\Lambda}_{\text{gh}} \right] - 12 I \left[ \frac{1}{2(1 + r_t)} Q_0(q; c_t), \tilde{\Lambda}_{\text{stop}} \right]
\]
where \( r_t = (m_b/m_t)^2 \), and \( \theta_W \) and \( c_t \) are the Weinberg angle and the bulk mass parameter for the top multiplet, respectively. The \( \theta_H \)-dependence enters through \( Q_0(q; c) \). The values of the parameters with a given \( z_L \) are summarized in Table\[I\]. The effective potential has the global minima at \( \theta_H = \pm \frac{1}{2} \pi \) for \( z_L \gg 1 \).
TABLE I: The values of the parameters in the model employed in the evaluation of \( V_{\text{eff}} \) and \( m_h \). The bulk mass parameter \( c_t \) is determined from \( m_t = 171.17 \text{ GeV} \). \( \sin^2 \theta_W \) is determined by global fit of the forward-backward asymmetries in \( e^+e^- \) collisions on the Z pole and the branching fractions of Z decay.[13] \( k \) and \( m_{\text{KK}} \) are in units of GeV.

| \( z_L \) | \( k \) | \( m_{\text{KK}} \) | \( \sin^2 \theta_W \) | \( c_t \) |
|---|---|---|---|---|
| \( 10^{15} \) | \( 4.67 \times 10^{17} \) | 1466 | 0.2309 | 0.432 |
| \( 10^{17} \) | \( 4.97 \times 10^{19} \) | 1562 | 0.2310 | 0.440 |

The mass of the Higgs boson \( m_h \) is related to the effective potential \( V_{\text{eff}}(\theta_H) \) by \( m_h^2 = f_H^{-2}(d^2V_{\text{eff}}/d\theta_H^2)\big|_{\theta_H=\pi/2} \) where \( \frac{1}{2}g_w f_H = (k/L)^{1/2}(z_L^2 - 1)^{-1/2} \sim m_W \). \( g_w \) is the 4D weak \( SU(2)_L \) gauge coupling. Noting that the KK mass scale is given by \( m_{\text{KK}} \sim \pi k z_L^{-1} \), one finds

\[
m_h^2 = \frac{g_w^2 k L m_{\text{KK}}^2}{32\pi^4} \left\{ -4J\left[\frac{1}{2}Q_0(q; \frac{1}{2}), \tilde{\Lambda}_{gh}\right] + 2J\left[\frac{1}{2\cos^2 \theta_W} Q_0(q; \frac{1}{2}), \tilde{\Lambda}_{gh}\right] -3J\left[Q_0(q; \frac{1}{2}), \tilde{\Lambda}_{gh}\right] + 12J\left[\frac{1}{2(1 + r_t)} Q_0(q; c_t), \tilde{\Lambda}_{\text{stop}}\right] \right\},
\]

(8)

Given \( z_L, m_h \) is determined as a function of \( \Lambda_{gh} \) and \( \Lambda_{\text{stop}} \). The result is summarized in fig. 2.

For small \( \Lambda_{\text{stop}} \) (\( \lesssim 200 \text{ GeV} \)) with \( \Lambda_{gh} \gtrsim 600 \text{ GeV} \), \( V_{\text{eff}}(\theta_H) \) is minimized at \( \theta_H = 0 \) so that the EW symmetry remains unbroken. If both \( \Lambda_{gh} \) and \( \Lambda_{\text{stop}} \) are larger than 1 TeV, the model is reduced to the non-supersymmetric model.

For \( \Lambda_{gh} \gtrsim 1 \text{ TeV} \), the desired \( m_h = 70 \sim 75 \text{ GeV} \) is obtained with \( \Lambda_{\text{stop}} = 450 \sim 475 \text{ GeV} \) with tiny dependence on \( z_L \) in the range \( 10^{15} \sim 10^{17} \). With these values of \( \Lambda_{\text{stop}} \) one finds the mass of the stop to be \( m_{\tilde{t}} = 480 \sim 505 \text{ GeV} \). In the analysis we have not specified masses of the sfermions except for the stop. If these sfermions are sufficiently heavy, evading the current bounds by LHC data, the stop or gravitino would become the lightest SUSY particle.

For \( \Lambda_{gh} \lesssim 100 \text{ GeV} \) implying a light neutralino, the Higgs mass \( m_h = 70 \sim 75 \text{ GeV} \) is obtained with \( \Lambda_{\text{stop}} = 250 \sim 275 \text{ GeV} \), corresponding to \( m_{\tilde{t}} = 300 \sim 320 \text{ GeV} \). We stress that SUSY breaking scales for other quarks and leptons can be much larger (\( \gtrsim 1 \text{ TeV} \)), which does
FIG. 2: SUSY breaking scales. Solid (dashed) lines correspond to $z_L = 10^{15}$ ($10^{17}$). Below the bottom solid (dashed) line the global minimum of $V_{\text{eff}}(\theta_H)$ is located at $\theta_H = 0$ so that the EW symmetry remains unbroken.

not affect the above result. There arises no constraint for gluino masses from this analysis so that gluinos can be heavier than 1 TeV. $m_{\tilde{t}} = 300 \sim 320$ GeV with $\Lambda_{\text{gh}} < \sim 100$ GeV is in the range allowed by current experiments.[19]

In the above analysis we have supposed that the stop masses are degenerate. Though unnecessary in the current scheme, it may be of interest to see the effect of large stop mixing, which plays an important role in MSSM to obtain a desired Higgs mass[23]. In the presence of the left-right squark mixing, the stop masses become non-degenerate. The spectra of their KK towers are approximated by $m_{\text{stop},i}^{\text{top},i} = \sqrt{(m_{\text{top},i})^2 + \Lambda_{\text{stop},i}^2}$, $i = 1, 2$, and accordingly the last terms in Eq. (7) and in the first equation of (8) are separated into two parts. As an extremal case we consider the case where one of the stops is very heavy and decouple. In such a case the curves in Fig. 2 are shifted downward. For example we obtain $\Lambda_{\text{stop}} \sim 260$ GeV [600 GeV] for $\Lambda_{\text{gh}} = 100$ GeV [1000 GeV] to obtain $m_h = 120$ GeV. The Higgs mass $m_h$ can be lowered only to $110$ GeV [84 GeV] for $\Lambda_{\text{gh}} = 100$ GeV [1000 GeV]. To obtain $m_h = 70 \sim 75$ GeV, it is desirable to have approximately degenerate stop masses in the current scheme.
In the present analysis we adopted a mass spectrum of a SUSY KK partner in the form (2) for convenience. Depending on how SUSY is broken, the spectrum may deviate from (2). However, the detailed form of the spectrum is not relevant in the present analysis, provided that $\overline{m}_n > m_n$ and $\overline{m}_n \sim m_n$ for $m_n \gg \Lambda$. Only low lying modes in the KK towers give relevant contributions to the $\theta_H$-dependent part of $V_{\text{eff}}(\theta_H)$. Contributions coming from the modes with $m_n, \overline{m}_n \gg m_{\text{KK}}$ are irrelevant.

We need a further consideration of the consistency with the current electroweak precision measurements, especially with Peskin-Takeuchi $S$ and $T$ parameters. There are many studies on these parameter in the models of extra dimensions. The composite Higgs models, which are regarded as holographic duals of the five-dimensional gauge-Higgs unification models, are severely constrained by the precision measurements. On the other hand it has been shown that the gauge couplings of gauge bosons, leptons and quarks in the $SO(5) \times U(1)$ gauge-Higgs unification model deviate little from those in the standard model, which indicates subtle difference between the composite Higgs models and the $SO(5) \times U(1)$ gauge-Higgs unification model. The spontaneous breaking of $SO(4) \times U(1)$ to $SU(2) \times U(1)$ triggered by a brane scalar field is crucial to have a realistic model of the electroweak symmetry breaking, which has not been properly taken into account in the literature. Recently it has been noticed that the symmetry group of the standard model may rotate in the $SO(5)$ group space according to the value of $\theta_H$. This certainly necessitates reexamination of $S$ and $T$ in the model.

To conclude, the dark matter of the Universe can be explained by stable Higgs bosons in the supersymmetric extension of the $SO(5) \times U(1)$ gauge-Higgs unification in the RS spacetime without spoiling the consistency with collider data at low energies, if $m_{\tilde{t}} = 300 \sim 320$ GeV when gauginos in the electroweak sector are light. The masses of gluinos as well as other squarks and sleptons do not affect the result. It would be of extreme importance to find the stop $\tilde{t}$ at LHC to get insight into the structure of spacetime.

Besides studying the electroweak precision observables, it is necessary to complete the model by incorporating flavor mixings and implementing light neutrinos by the seesaw mechanism in the bulk-brane system. We hope to report on these in the near future.

**Acknowledgements:** The authors would like to thank K. Oda and M. Tanaka for valuable comments. This work was supported in part by scientific grants from the Ministry
References

[1] Y. Hosotani, *Phys. Lett.* B126, 309 (1983); *Ann. Phys. (N.Y.)* 190, 233 (1989).

[2] A. T. Davies and A. McLachlan, *Phys. Lett.* B200, 305 (1988) *Nucl. Phys.* B317, 237 (1989).

[3] H. Hatanaka, T. Inami and C.S. Lim, *Mod. Phys. Lett.* A13, 2601 (1998).

[4] C.A. Scrucca, M. Serone and L. Silvestrini, *Nucl. Phys.* B669, 128 (2003); G. Burdman and Y. Nomura, *Nucl. Phys.* B656, 3 (2003); C. Csaki, C. Grojean and H. Murayama, *Phys. Rev.* D67, 085012 (2003).

[5] K. Agashe, R. Contino and A. Pomarol, *Nucl. Phys.* B719, 165 (2005).

[6] A. D. Medina, N. R. Shah and C. E. M. Wagner, *Phys. Rev.* D76, 095010 (2007); Y. Hosotani and Y. Sakamura, *Prog. Theoret. Phys.* 118, 935 (2007); Y. Sakamura, *Phys. Rev.* D76, 065002 (2007). Y. Hosotani and Y. Kobayashi, *Phys. Lett.* B674, 192 (2009).

[7] G.F. Giudice, C. Grojean, A. Pomarol and R. Rattazzi, *JHEP* 0706, 045 (2007); K. Agashe, A. Azatov, T. Han, Y. Li, Z.-G. Si, and L. Zh, *Phys. Rev.* D81, 096002 (2010).

[8] Y. Hosotani, K. Oda, T. Ohnuma and Y. Sakamura, *Phys. Rev.* D78, 096002 (2008); *Erratum-ibid.* 79, 079902 (2009).

[9] Y. Hosotani, S. Noda and N. Uekusa, *Prog. Theoret. Phys.* 123, 757 (2010).

[10] Y. Hosotani, P. Ko and M. Tanaka, *Phys. Lett.* B680, 179 (2009).

[11] K. Cheung and J. Song, *Phys. Rev.* D81, 097703 (2010); *Erratum-ibid.* 81, 119905 (2010); A. Alves, *Phys. Rev.* D82, 115021 (2010).

[12] Y. Hosotani, M. Tanaka and N. Uekusa, *Phys. Rev.* D82, 115024 (2010).

[13] Y. Hosotani, M. Tanaka, and N. Uekusa, *Phys. Rev.* D84, 075014 (2011).

[14] R. Contino, D. Marzocca, D. Pappadopulo, R. Rattazzi, *JHEP* 1110, 081 (2011).

[15] E. Komatsu *et al.*, WMAP Collaboration, *Astrophys. J. Suppl.* 180, 330 (2009).

[16] L. Randall and R. Sundrum, *Phys. Rev. Lett.* 83, 3370 (1999).

[17] J. Garriga, O. Pujolas, and T. Tanaka, *Nucl. Phys.* B605, 192 (2001); A. Falkowski, *Phys. Rev.* D75, 025017 (2007).
[18] K. Oda and A. Weiler, Phys. Lett. B606, 408 (2005); H. Hatanaka, arXiv:0712.1334 [hep-th]; N. Haba, S. Matsumoto, N. Okada and T. Yamashita, Prog. Theoret. Phys. 120, 77 (2008); N. Maru and Y. Sakamura, JHEP 1004, 100 (2010); Y. Sakamura, Phys. Rev. D83, 036007 (2011).

[19] ATLAS Collaboration, Phys. Lett. B701, 398 (2011); ATLAS NOTE, ATLAS-CONF-2012-003.

[20] M.E. Peskin and T. Takeuchi, Phys. Rev. Lett. 65, 964 (1990); Phys. Rev. D46, 381 (1992).

[21] G. Cacciapaglia, C. Csaki, C. Grojean and J. Terning, Phys. Rev. D70, 075014 (2004).

[22] K. Agashe and R. Contino, Nucl. Phys. B742, 59 (2006); K. Agashe, R. Contino, L. Da Rold and A. Pomarol, Phys. Lett. B641, 62 (2006); M. S. Carena, E. Ponton, J. Santiago and C. E. M. Wagner, Phys. Rev. D76, 035006 (2007); G. Panico, M. Serone and A. Wulzer, Nucl. Phys. B762, 189 (2007); R. Barbieri, B. Bellazzini, V.S. Rychkov, A. Varagnolo, Phys. Rev. D76, 115008 (2007).

[23] P. Draper, P. Meade, M. Reece and D. Shih, arXiv:1112.3068 [hep-ph].

[24] Y. Adachi, N. Kurahashi, C.S. Lim, N. Maru, JHEP 1011, 150 (2010); Y. Adachi, N. Kurahashi, N. Maru, K. Tanabe, arXiv:1201.2290 [hep-ph].