WWZ, WWH, and ZZH Couplings
in the Dynamical Gauge-Higgs Unification
in the Warped Spacetime

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Abstract

In the dynamical gauge-Higgs unification in the Randall-Sundrum warped spacetime, where the 4D Higgs field is unified with gauge fields and the electroweak symmetry is dynamically broken by the Hosotani mechanism, the trilinear couplings for WWZ, WWH, and ZZH, where H stands for the Higgs field, are evaluated. The latter two couplings are suppressed by a factor of cos θH where θH is the Yang-Mills Aharonov-Bohm phase in the extra dimension, while the WWZ couplings remain the same as in the standard model to good accuracy.

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The Higgs field in the standard model of electroweak interactions plays a vital role in the electroweak symmetry breaking and in giving masses to $W$, $Z$, quarks and leptons. The Higgs boson is expected to be discovered at LHC in the near future. We are entering in the era when the structure of the electroweak symmetry breaking is disclosed.

It is not clear, however, if the Higgs sector in the standard model remains valid at the fundamental level. It has been argued that the Higgs boson mass suffers from quadratically divergent radiative corrections unless protected by symmetry, which requires unnatural fine tuning of parameters in the theory. The leading candidate for overcoming this theoretical unnaturalness is supersymmetry. The minimal supersymmetric standard model (MSSM) predicts the mass of the Higgs boson, $m_H$, to be less than 130 GeV.\[1\] The experimental lower bound for $m_H$ is 114 GeV.\[2\]

Many alternative scenarios for the Higgs sector in the electroweak interactions have been proposed, including the little Higgs model,\[3\] the Higgsless model,\[4, 5\] and the gauge-Higgs unification models.\[6\]-\[30\] Among them the dynamical gauge-Higgs unification predicts various properties in the Higgs field couplings and gauge field couplings which differ from those in the standard model and can be tested experimentally at LHC and future linear colliders.

In the gauge-Higgs unification scenario the Higgs field in four dimensions is unified with gauge fields within the framework of higher dimensional gauge theory. Low energy modes of extra-dimensional components of gauge potentials are 4D Higgs scalar fields. Fairlie and Manton proposed gauge-Higgs unification in six dimensions with ad hoc symmetry ansatz.\[6, 7\]. Justification for the ansatz was attempted by making use of quantum dynamics, but was afflicted with the cut-off dependence.\[8\] More attractive scheme is obtained when the extra-dimensional space is non-simply connected.\[9, 10\] There appear Yang-Mills Aharonov-Bohm phases, $\theta_H$, associated with the gauge field holonomy, or the phases of Wilson line integrals along noncontractible loops. Although classical vacua are degenerate with respect to the values of $\theta_H$, quantum dynamics of $\theta_H$ lifts the degeneracy and the non-Abelian gauge symmetry is dynamically broken. It is called the Hosotani mechanism. Fluctuations of $\theta_H$ in four dimensions correspond to the 4D Higgs field. With dynamical gauge symmetry breaking, the dynamical gauge-Higgs unification is achieved.

In recent years the dynamical gauge-Higgs unification has been applied to the electroweak interactions. Chiral fermions are naturally accommodated in the scheme by considering an orbifold as extra-dimensional space.\[13, 14\] To have an $SU(2)_L$ doublet Higgs
field, one has to start with a gauge group larger than \( SU(2)_L \times U(1)_Y \). As the Higgs field is a part of gauge fields, most of the couplings associated with the Higgs field are tightly constrained by the gauge principle. In flat space \( m_H \) typically turns out to be
\[
\sim \left( \frac{g_{SU(2)}^2}{4\pi} \right)^{1/2} m_W,
\]
which contradicts with the observation. It is nontrivial to obtain quark-lepton mass matrix naturally.\[15\]-[22] For instance, one needs fermions in higher dimensional representation of the group.\[20\]

These problems can be resolved in the dynamical gauge-Higgs unification in the Randall-Sundrum warped spacetime.\[23\]-[30] \( m_H \) is predicted in the range 120 GeV \( \sim \) 290 GeV. The hierarchical mass spectrum of quarks and leptons is naturally explained in terms of bulk kink masses, with which couplings of quarks and leptons to gauge bosons and their Kaluza-Klein excited states are determined. It was pointed out that the universality in the weak gauge couplings is slightly broken, and Yukawa couplings of quarks and leptons are substantially reduced compared with those in the standard model.\[27, 28\]

The previous model based on the gauge group \( SU(3) \) is unsatisfactory in many respects. It gives the incorrect Weinberg angle \( \theta_W \) and the neutral current sector is unrealistic. It is also difficult to have a realistic fermion mass matrix. It has been argued that dynamics at the boundaries (fixed points) of the orbifold such as brane kinetic terms of gauge fields can reproduce the observed \( \theta_W \).\[31\]

More promising approach is to adapt a gauge group \( SO(5) \times U(1)_{B-L} \) to start with, as advocated by Agashe, Contino and Pomarol.\[24\] The custodial symmetry in the 4D Higgs field sector is contained in \( SO(5) \) and the correct \( \theta_W \) is reproduced so that phenomenology in the neutral currents can be reliably discussed. In this paper we focus mainly on the gauge couplings among the gauge bosons and the Higgs boson, to find substantial deviation from those in the standard model. We briefly describe how interactions of fermion multiplets should be introduced to have realistic gauge couplings and mass matrix, but the detailed discussions are reserved for future work.

We add that the dynamical gauge-Higgs unification is defined not only at the tree and one-loop levels, but also beyond one loop. It has been argued recently that the Higgs boson mass \( m_H \), for instance, may be finite to all order in five dimensions, indicating that the properties of the Higgs boson can be determined independent of physics at the cutoff scale.\[32\]-[37]

The model we consider is \( SO(5) \times U(1)_{B-L} \) gauge theory in the Randall-Sundrum (RS) geometry in five dimensions.\[39\] We use \( M, N, \cdots = 0, 1, 2, 3, 4 \) for the 5D curved indices,
\( A, B, \cdots = 0, 1, 2, 3, 4 \) for the 5D flat indices in tetrads, and \( \mu, \nu, \cdots = 0, 1, 2, 3 \) for 4D indices. The background metric is given by

\[
 ds^2 = G_{MN} dx^M dx^N = e^{-2\sigma(y)} \eta_{\mu\nu} dx^\mu dx^\nu + dy^2 ,
\]

(1)

where \( \eta_{\mu\nu} = \text{diag} (-1, 1, 1, 1) \), \( \sigma(y) = \sigma(y + 2\pi R) \), and \( \sigma(y) \equiv k |y| \) for \( |y| \leq \pi R \). The cosmological constant in the bulk 5D spacetime is given by \( \Lambda = -k^2 \). \( (x^\mu, -y) \) and \( (x^\mu, y + 2\pi R) \) are identified with \( (x^\mu, y) \). The spacetime is equivalent to the interval in the fifth dimension \( y \) with two boundaries at \( y = 0 \) and \( y = \pi R \), which we refer to as the Planck brane and the TeV brane, respectively.

There are the \( SO(5) \) gauge field \( A_M \) and the \( U(1)_{B-L} \) gauge field \( B_M \), the former of which is decomposed as

\[
 A_M = \sum_{l=1}^{10} A^l_M T^l = \sum_{a_L=1}^{3} A^{a_L}_M T^{a_L} + \sum_{a_R=1}^{3} A^{a_R}_M T^{a_R} + \sum_{\hat{a}=1}^{4} A_{\hat{a}} M T_{\hat{a}},
\]

(2)

where \( T^{a_L, a_R} \) \((a_L, a_R = 1, 2, 3)\) and \( T_{\hat{a}} \) \((\hat{a} = 1, 2, 3, 4)\) are the generators of \( SO(4) \sim SU(2)_L \times SU(2)_R \) and \( SO(5)/SO(4) \), respectively. As a matter field, we introduce a spinor field \( \Psi \), belonging to the spinorial representation of \( SO(5) \) \((i.e., 4 \text{ of } SO(5))\). Extension to multi-spinor case is straightforward. We will argue later that multiple spinor fields are necessary to have phenomenologically acceptable fermion content even in the one generation case.

The relevant part of the action is

\[
 S = \int d^5 x \sqrt{-G} \left[ -\text{tr} \left( \frac{1}{2} F^{(A)MN} F^{(A)}_{MN} + \frac{1}{\xi} (f^{(A)}_{gf})^2 + L_{gh}^{(A)} \right) \right.
\]

\[
 \left. - \left( \frac{1}{4} F^{(B)MN} F^{(B)}_{MN} + \frac{1}{\xi} (f^{(B)}_{gf})^2 + L_{gh}^{(B)} \right) + i \bar{\Psi} \Gamma^N D_N \Psi - i M_\Psi \varepsilon \bar{\Psi} \Psi \right] ,
\]

(3)

where \( G \equiv \det(G_{MN}) \), \( \Gamma^N \equiv e_A^N \Gamma^A \). \( \Gamma^A \) is a 5D \( \gamma \)-matrix. \( f^{(A,B)}_{gf} \) are the gauge-fixing functions, \( L_{gh}^{(A,B)} \) are the associated ghost Lagrangians, and \( M_\Psi \) is a bulk mass parameter. \( [40] \)

Since the operator \( \bar{\Psi} \Psi \) is \( Z_2 \)-odd, we need the periodic sign function \( \varepsilon(y) = \sigma'(y)/k \) satisfying \( \varepsilon(y) = \pm 1 \). The field strengths and the covariant derivatives are defined by

\[
 F^{(A)}_{MN} \equiv \partial_M A_N - \partial_N A_M - i g_A [A_M, A_N],
\]

\[
 F^{(B)}_{MN} \equiv \partial_M B_N - \partial_N B_M ,
\]

\[
 D_M \Psi \equiv \left\{ \partial_M - \frac{1}{4} \omega_M^{AB} \Gamma_{AB} - i g_A A_M - i \frac{g_B}{2} q_{B-L} B_M \right\} \Psi,
\]

(4)
where $g_A$ ($g_B$) is the 5D gauge coupling for $A_M$ ($B_M$), $\Gamma^{AB} \equiv \frac{1}{2}[\Gamma^A, \Gamma^B]$, and $q_{B-L}$ is a charge of $U(1)_{B-L}$. The spin connection 1-form $\omega^{AB} = \omega^{AB}_M dx^M$ determined from the metric (1) is $\omega^{\nu A} = -\sigma^\nu dx^\nu$ with all other components vanishing.

The boundary conditions consistent with the orbifold structure are written as [11]

\[
\begin{align*}
(A^\mu_A)(x, -y) &= P_0 \begin{pmatrix} A^\mu_A \\ -A_y \end{pmatrix}(x, y) P_0^{-1}, \\
(A^\mu_A)(x, \pi R - y) &= P_\pi \begin{pmatrix} A^\mu_A \\ -A_y \end{pmatrix}(x, \pi R + y) P_\pi^{-1}, \\
(B^\mu_B)(x, -y) &= \begin{pmatrix} B^\mu_A \\ -B_y \end{pmatrix}(x, y), \quad \begin{pmatrix} B^\mu_B \\ -B_y \end{pmatrix}(x, \pi R - y) = \begin{pmatrix} B^\mu_B \\ -B_y \end{pmatrix}(x, \pi R + y), \\
\Psi(x, -y) &= \eta_0 P_0 \gamma_5 \Psi(x, y), \quad \Psi(x, \pi R - y) = \eta_\pi P_\pi \gamma_5 \Psi(x, \pi R + y),
\end{align*}
\tag{5}
\]

where $\gamma_5 \equiv \Gamma^4$ is the 4D chiral operator, $\eta_j = \pm 1$, $P_j \in SO(5)$ and $P_j^2 = 1$. In the present paper, we take $P_0$ and $P_\pi$ given by

\[
P_0 = P_\pi = \begin{pmatrix} 12 \\ -12 \end{pmatrix}
\tag{6}
\]

in the spinorial representation, or equivalently $P_0 = P_\pi = \text{diag} (-1, -1, -1, -1, 1)$ in the vectorial representation. The boundary condition (5) breaks the gauge symmetry to $SO(4) \times U(1)_{B-L} \sim SU(2)_L \times SU(2)_R \times U(1)_{B-L}$ at both boundaries. (The broken generators are $T^\hat{a}$; $\hat{a} = 1, 2, 3, 4$.) It is convenient to decompose $\Psi$ as

\[
\Psi = \begin{pmatrix} q \\ Q \end{pmatrix},
\tag{7}
\]

where $q$ and $Q$ belong to $(2, 1)$ and $(1, 2)$ of $SU(2)_L \times SU(2)_R$, respectively.

First notice that with (5) and (6) there arise zero modes for $A^\hat{a}_y$ ($\hat{a} = 1, \ldots, 4$), which correspond to the $SU(2)_L$ doublet Higgs field in the standard model; $\Phi \propto (A^1_y + iA^2_y, A^4_y - iA^3_y)$. They also give rise to Yang-Mills Aharonov-Bohm phases, $\theta_H$, in the fifth dimension. Making use of the residual symmetry, one can suppose that the zero mode of $A^4_y$ develops a nonvanishing expectation value;

\[
A^4_y = \frac{2\sqrt{2k} e^{2ky}}{g_A(z_2^2 - 1)} \theta_H
\tag{8}
\]

\footnote{For an explicit representation of the generators $T^I$, see the appendix in the first reference in Ref. [24].}
where \( z_\pi = e^{\pi k R} \). Although \( \theta_H \neq 0 \) gives vanishing field strengths, it affects physics at the quantum level. The effective potential for \( \theta_H \) becomes non-trivial at the one loop level, whose global minimum determines the quantum vacuum. It is this nonvanishing \( \theta_H \) that induces dynamical electroweak gauge symmetry breaking.

There are residual gauge transformations which maintain the boundary condition (5). Among them there is a large gauge transformation given by

\[
\Omega^{\text{large}}(y) = \exp \left\{ i n \pi \frac{e^{2k y} - 1}{z_\pi^2 - 1} \right\} 2\sqrt{2} T^4 \tag{9}
\]

for \( 0 \leq y \leq \pi R \) where \( n \) is an integer. Under the transformation (9), \( \theta_H \) is transformed to \( \theta_H + 2\pi n \), which implies that all physical quantities are periodic functions of \( \theta_H \). This large gauge invariance has to be maintained at all stages. It is vital to guarantee the finiteness of the Higgs boson mass. The \( \theta_H \)-dependent part of the effective potential for \( \theta_H \) diverges without the large gauge invariance.

It is important to recognize that the even-odd property in (5) does not completely fix boundary conditions of the fields. If there are no additional dynamics on the two branes, fields which are odd under parity at \( y = 0 \) or \( \pi R \) obey the Dirichlet boundary condition (D) so that they vanish there. On the other hand, fields which are even under parity obey the Neumann boundary conditions (N). For gauge fields the Neumann boundary condition is given by \( dA_\mu / dy = 0 \) or \( d(e^{-2k y} A_y) / dy = 0 \). As a result of additional dynamics on the branes, however, a field with even parity, for instance, can obey the Dirichelet boundary condition, provided the large gauge invariance is maintained. We argue below that this, indeed, happens on the Planck brane.

Let us define new fields \( A^3_R^M \) and \( A_Y^M \) by

\[
\begin{pmatrix}
A^3_R^M \\
A_Y^M
\end{pmatrix}
= \begin{pmatrix}
c_\phi & -s_\phi \\
s_\phi & c_\phi
\end{pmatrix}
\begin{pmatrix}
A^3_R^M \\
B_M
\end{pmatrix},
\]

\[
c_\phi \equiv \frac{g_A}{\sqrt{g_A^2 + g_B^2}}, \quad s_\phi \equiv \frac{g_B}{\sqrt{g_A^2 + g_B^2}}. \tag{10}
\]

\( A^{3R}_\mu \) and \( B_\mu \) are even under parity, whereas \( A^{3Y}_\mu \) and \( B_y \) are odd. It is our contention that the even fields \( A^{1R}_\mu \), \( A^{2R}_\mu \), and \( A^3_R^\mu \) obey the Dirichlet (D) boundary condition on the Planck brane as a result of additional dynamics there. The boundary conditions for gauge fields are tabulated in Table I. It is straightforward to confirm that the boundary conditions in Table I preserve the large gauge invariance, that is, new gauge potentials obtained by (9)
obey the same boundary conditions as the original fields. We note that the Neumann (N) boundary condition on the Planck brane cannot be imposed on $A_{y}^{1R}$, $A_{y}^{2R}$, and $A_{y}^{3R}$, as it does not preserve the large gauge invariance. A similar conclusion has been obtained in Ref. [25].

With the condition in Table II, the gauge symmetry $SO(5) \times U(1)_{B-L}$ in the bulk is reduced to $SO(4) \times U(1)_{B-L}$ at the TeV brane and to $SU(2)_{L} \times U(1)_{Y}$ at the Planck brane. The resultant symmetry of the theory is $SU(2)_{L} \times U(1)_{B-L}$ at the TeV brane and to $SU(2)_{L} \times U(1)_{Y}$ at the Planck brane. The weak hypercharge $Y$ is given by $Y = T_{3R} + q_{B-L}/2$.

One way to achieve the change of the boundary conditions of $A_{y}^{1R}$, $A_{y}^{2R}$, and $A_{y}^{3R}$ from N to D on the Planck brane is to have additional fields and dynamics on the Planck brane such that $SU(2)_{R} \times U(1)_{B-L}$ is spontaneously broken to $U(1)_{Y}$ at relatively high energy scale $M$, say, near the Planck scale $M_{Pl}$. Below the scale $M$, the mass terms

$$L_{\text{mass}} = -\left\{M_{1}^{2}(A_{\mu}^{1R}A_{\mu}^{1R} + A_{\mu}^{2R}A_{\mu}^{2R}) + M_{2}^{2}(A_{\mu}^{3R}A_{\mu}^{3R})\right\}\delta(y), \quad (11)$$

where $M_{1}, M_{2} = O(M)$, are induced on the Planck brane. Below the TeV scale, the mass terms (11) strongly suppress the boundary values of $(A_{\mu}^{1R}, A_{\mu}^{2R}, A_{\mu}^{3R})$ on the Planck brane, changing the boundary conditions from N to D at the Planck brane. We note that when masses are induced by spontaneous symmetry breaking on the Planck brane, well-controlled ultra-violet behavior of gauge bosons is not spoiled so that the finiteness of the 4D Higgs boson mass at the one loop level, for instance, is expected to be maintained.

It has been shown recently that the requirement of the tree level unitarity constrains boundary conditions satisfied by gauge bosons. [42, 43] With the underlying mechanism of spontaneous symmetry breaking, the effective boundary conditions in Table I are expected to preserve the tree level unitarity.

Now we expand the bulk fields in 4D KK modes. It is convenient to use the conformal coordinate $z \equiv e^{\sigma(y)}$ for the fifth dimension, in which the boundaries are located at $z = 1$ and $z_{\pi} = e^{k\pi R}$. As in Ref. [28], we split $A_{M}$ into the classical and quantum parts: $A_{M} = A_{M}^{c} + A_{M}^{q}$. We take $A_{\mu}^{c} = 0$ and $A_{y}^{c} = (dz/dy)A_{z}^{c}$ given by (8). Further we move to a new basis by a gauge transformation[4]

$$\tilde{A}_{M} = \Omega A_{M}^{q} \Omega^{-1}, \quad \tilde{B}_{M} = B_{M}^{q}, \quad \begin{pmatrix} \tilde{q} \\ \tilde{Q} \end{pmatrix} = z^{-2}\Omega \begin{pmatrix} q \\ Q \end{pmatrix},$$

[2] There arises a type II defect on the Planck brane in the terminology of Ref. [25].

[3] $\Omega(z)$ is defined here such that $\Omega(1) = 1$ whereas in our previous work [28] $\Omega(z_{\pi}) = 1$. 

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| $A_{\mu}^{\alpha L}$ | $A_{\mu}^{1,2R}$ | $A_{\mu}^{3R}$ | $A_{\mu}^{Y}$ | $A_{\mu}^{\hat{a}}$ |
|-----------------|-----------------|-----------------|-----------------|-----------------|
| (N,N)           | (D,N)           | (N,N)           | (D,D)           | (N,N)           |
| $A_{y}^{\alpha L}$ | $A_{y}^{1,2R}$ | $A_{y}^{3R}$ | $A_{y}^{Y}$ | $A_{y}^{\hat{a}}$ |
| (D,D)           | (D,D)           | (D,D)           | (D,D)           | (N,N)           |

Table I: Boundary conditions for the gauge fields. $a_{L}, a_{R} = 1, 2, 3$ and $\hat{a} = 1, 2, 3, 4$. The notation (D,N), for example, denotes the Dirichlet boundary condition at $y = 0$ and the Neumann boundary condition at $y = \pi R$.

$$\Omega(z) = \exp \left\{ ig_{A} \int_{z}^{z'} dz' \ A_{z}'(z') \right\} .$$ (12)

In the new basis the classical background of the gauge fields vanishes so that the linearized equations of motion reduce to the simple forms, while the boundary conditions become more involved. The 5D fields are expanded into the 4D modes as

$$\tilde{A}_{\mu}(x, z) = \sum_{n} \tilde{h}_{A,n}(z) A_{\mu}^{(n)}(x), \quad \tilde{A}_{z}(x, z) = \sum_{n} \tilde{h}_{\varphi,n}(z) \varphi^{(n)}(x),$$

$$\tilde{B}_{\mu}(x, z) = \sum_{n} \tilde{h}_{B,n}(z) A_{\mu}^{(n)}(x), \quad \tilde{B}_{z}(x, z) = \sum_{n} \tilde{h}_{\varphi,n}(z) \varphi^{(n)}(x),$$

$$\tilde{q}_{R}(x, z) = \sum_{n} \tilde{f}_{R,n}(z) \psi_{R}^{(n)}(x), \quad \tilde{Q}_{R}(x, z) = \sum_{n} \tilde{f}_{R,n}(z) \psi_{R}^{(n)}(x),$$

$$\tilde{q}_{L}(x, z) = \sum_{n} \tilde{f}_{L,n}(z) \psi_{L}^{(n)}(x), \quad \tilde{Q}_{L}(x, z) = \sum_{n} \tilde{f}_{L,n}(z) \psi_{L}^{(n)}(x).$$ (13)

From the action (3) and the boundary conditions, the mass spectrum and analytic expressions of the mode functions are obtained in terms of the Bessel functions in the same manner as in Ref. [28].

Gauge fields are classified in three sectors. There are the charged sector (or the $W$ boson sector)

$$(A_{3}^{\pm L}, A_{3}^{\pm R}, A_{M}^{\hat{a}}) \equiv \frac{1}{\sqrt{2}} (A_{M}^{1L} \pm i A_{M}^{2L}, A_{M}^{1R} \pm i A_{M}^{2R}, A_{M}^{1} \pm i A_{M}^{2}) ,$$ (14)

the neutral sector

$$(A_{3}^{\alpha L}, A_{3}^{\alpha R}, A_{M}^{\hat{a}}, B_{M}) ,$$ (15)

and the “Higgs” sector $A_{M}^{\hat{a}}$, which is also neutral. The neutral sector (15) consists of the $Z$ boson sector and the photon sector.

The mass spectrum ($m = k \lambda$) in the charged sector is determined by

$$F_{0,1} \left\{ \lambda^{2} z_{\pi} F_{0,0} F_{1,1} - \frac{2}{\pi^{2}} \sin^{2} \theta_{H} \right\} = 0$$ (16)

8
where \( F_{\alpha,\beta} = J_\alpha(\lambda z_\pi)Y_\beta(\lambda) - Y_\alpha(\lambda z_\pi)J_\beta(\lambda) \). We note that the \( \theta_H \)-dependence appears in the form \( \sin^2 \theta_H \) in the \( SO(5) \times U(1)_{B-L} \) model, whereas it appeared in the form \( \sin^2 \frac{\theta_H}{2} \) in the \( SU(3) \) model in [28].

The mass of the lightest mode, the \( W \) boson \( W^{(0)}_\mu(x) \), is approximately given by

\[
m_W \simeq \frac{m_{KK}}{\pi} \sqrt{\frac{1}{k\pi R}} |\sin \theta_H|, \tag{17}
\]

for \( z_\pi \gg 1 \) where \( m_{KK} \equiv k\pi/(z_\pi - 1) \) is the KK mass scale. Its mode functions are approximately given by

\[
\tilde{h}^{\pm L}_{A,0}(z) \simeq \frac{1 + \cos \theta_H}{2\sqrt{\pi R}}, \quad \tilde{h}^{\pm R}_{A,0}(z) \simeq \frac{1 - \cos \theta_H}{2\sqrt{\pi R}}, \quad \tilde{h}^A_{A,0}(z) \simeq \sin \theta_H \left( \frac{z^2}{z_\pi^2} - 1 \right). \tag{18}
\]

In the photon sector, there is a massless mode, namely the photon mode \( A^{(0)}_\mu(x) \), whose mode functions are constants.

\[
\tilde{h}^{3L}_{A,0}(z) = \tilde{h}^{3R}_{A,0}(z) = \frac{s_\phi}{\sqrt{(1 + s_\phi^2)\pi R}},
\]

\[
\tilde{h}^3_{A,0}(z) = 0, \quad \tilde{h}^B_{A,0}(z) = \frac{c_\phi}{\sqrt{(1 + s_\phi^2)\pi R}}. \tag{19}
\]

Here \( s_\phi \) and \( c_\phi \) are defined in [10]. The mass spectrum in the \( Z \)-boson sector is determined by

\[
F_{0,1} \left\{ \lambda^2 z_\pi F_{0,0}F_{1,1} - \frac{2}{\pi^2} \left( 1 + s_\phi^2 \right) \sin \theta_H \right\} = 0. \tag{20}
\]

The mass \( m_Z \) and the mode functions of the \( Z \) boson \( Z^{(0)}_\mu(x) \), which is the second lightest mode in the neutral sector [15], are approximately given by

\[
m_Z \simeq \frac{m_{KK}}{\pi} \sqrt{\frac{1 + s_\phi^2}{k\pi R}} |\sin \theta_H|, \tag{21}
\]

and

\[
\tilde{h}^{3L}_{A,1}(z) \simeq \frac{c_\phi^2 + \cos \theta_H (1 + s_\phi^2)}{2\sqrt{(1 + s_\phi^2)\pi R}}, \quad \tilde{h}^{3R}_{A,1}(z) \simeq \frac{c_\phi^2 - \cos \theta_H (1 + s_\phi^2)}{2\sqrt{(1 + s_\phi^2)\pi R}},
\]

\[
\tilde{h}^3_{A,1}(z) \simeq \sin \theta_H \sqrt{\frac{1 + s_\phi^2}{2\pi R} \left( \frac{z^2}{z_\pi^2} - 1 \right)}, \quad \tilde{h}^B_{A,1}(z) \simeq -\frac{s_\phi c_\phi}{\sqrt{(1 + s_\phi^2)\pi R}}. \tag{22}
\]

The lightest mode \( \varphi^{(0)}(x) \) in the Higgs sector corresponds to the 4D Higgs field, whose mode function is given by

\[
\tilde{h}^4_{\varphi,0}(z) = \sqrt{\frac{2}{k(z_\pi^2 - 1)}} z. \tag{23}
\]
At the classical level the potential for $\phi(0)$ is flat. Quantum effects yield nontrivial, finite corrections to the effective potential for $\phi(0)$, giving the Higgs boson a finite mass $m_H = \mathcal{O}(m_{\text{KK}}kR/(4\pi))$ \cite{27, 28}.

Some comments are in order regarding the mass spectrum determined by (16) and (20). First, $m_W$ and $m_Z$ are not proportional to the VEV of the “Higgs field”, or $\theta_H$, in contrast to the ordinary Higgs mechanism. This is because the Higgs mechanism, or the mechanism of mass generation, does not complete within each KK level and the lowest mode in each KK tower necessarily mixes with heavy KK modes when $\theta_H$ acquires a nonzero value. The similar property is seen in the $SU(3)$ model in \cite{28}. In the $SU(3)$ model, the linear dependence of the mass spectrum on $\theta_H$ recovers in the flat limit, $kR \to 0$. (See Sec. 5.1 in Ref. \cite{28}.) In the $SO(5) \times U(1)_{B-L}$ model, however, the mass spectrum deviates from the linear dependence even in the flat limit. This is due to the fact that the numerical factor in front of $\sin^2 \theta_H$ is a half of that in the $SU(3)$ model in \cite{28}. It is one of the distinctive properties of the $SO(5) \times U(1)_{B-L}$ model. Secondly, the mass spectrum of the modes corresponding to $F_{0,1} = 0$ in (16) or (20) is independent of $\theta_H$. Their mode functions, however, have nontrivial $\theta_H$-dependence. We note that these modes do not have definite $Z_2$-parities when $\theta_H$ acquires a nonzero value, although the condition $F_{0,1} = 0$ is the same as that for the modes which have the boundary condition $(D,N)$ at $\theta_H = 0$. The existence of such modes is one of the characteristics of the $SO(5) \times U(1)_{B-L}$ model.

From (17) and (21), the Weinberg angle $\theta_W$ determined from $m_W$ and $m_Z$ becomes

$$\sin^2 \theta_W \equiv 1 - \frac{m_W^2}{m_Z^2} \approx \frac{s^2_\phi}{1 + s^2_\phi} = \frac{g_B^2}{g_A^2 + 2g_B^2} = \frac{g_Y^2}{g_A^2 + g_Y^2}.$$ \hspace{1cm} (24)

The approximate equality in the second line is valid to the $\mathcal{O}(0.1\%)$ accuracy for $m_{\text{KK}} = \mathcal{O}(\text{TeV})$. In the last equality the relation $g_Y = g_Ag_B/\sqrt{g_A^2 + g_B^2}$ has been made use of. We note that $s_\phi \approx \tan \theta_W$. The Weinberg angle $\theta_W$ may be determined from the vertices in the neutral current interactions. As we will see in Eqs. (30)-(32) below, $\theta_W$ in this definition coincides with that in (24) to good accuracy. Thus the rho parameter is approximately one in our model.

The mass spectrum of a fermion multiplet (7) is determined by

$$\chi^2 z_\pi F_{c-\frac{1}{2},c-\frac{1}{2}}F_{c+\frac{1}{2},c+\frac{1}{2}} - \frac{4}{\pi^2} \left( \frac{\sin^2 \frac{1}{2} \theta_H}{\cos^2 \frac{1}{2} \theta_H} \right) \left( \eta_0 \eta_\pi \right) = 0 \quad \text{for} \quad \eta_0 \eta_\pi = \begin{pmatrix} +1 \\ -1 \end{pmatrix}. \hspace{1cm} (25) \quad \text{(25)}$$
Here $c = M_\Psi/k$. The lightest mass eigenvalue $m_f$ is approximately given by

$$m_f \simeq k \left\{ \frac{e^2 - \frac{1}{4}}{z_c \sinh \left( (c + \frac{1}{2}) k \pi R \right) \sinh \left( (c - \frac{1}{2}) k \pi R \right)} \right\}^{1/2} \left( \left| \sin \frac{1}{2} \theta_H \right| \right). \quad (26)$$

For $c > \frac{1}{2}$ and $(\eta_0, \eta_n) = (1, 1)$, the corresponding mode functions are approximately given by

$$\tilde{f}_{L,0}^q(z) \simeq i p_{H/2} \cos \frac{\theta_H}{2} \sqrt{2k(c - \frac{1}{2})} z^{-c}, \quad \tilde{f}_{L,0}^Q(z) \simeq -i \sin \frac{\theta_H}{2} \sqrt{2k(c - \frac{1}{2})} z^{-c},$$

$$\tilde{f}_{R,0}^q(z) \simeq -i \sin \theta_H \frac{\sqrt{k(c + \frac{1}{2})}}{\sqrt{2z_c^{c+\frac{1}{2}}}} z^{1-c}, \quad \tilde{f}_{R,0}^Q(z) \simeq -\frac{\sqrt{2k(c + \frac{1}{2})}}{z_c^{c+\frac{1}{2}}} z^c, \quad (27)$$

where $p_{H/2} \equiv \text{sgn}(\sin \frac{1}{2} \theta_H)$. As shown in [28], the hierarchical mass spectrum for fermions is obtained from (26) by varying the dimensionless parameter $c$ in an $O(1)$ range.

The $\theta_H$-dependence of $m_f$ differs from that of $m_W$ and $m_Z$. Consequently the ratios $m_f/m_W$, $m_f/m_Z$ depend on $\theta_H$ in the $SO(5) \times U(1)_{B-L}$ model in contrast to those in the ordinary Higgs mechanism$^4$.

Let us turn to the various coupling constants in the 4D effective theory. We first look at the 4D gauge coupling constants of fermions, which are obtained as overlap integrals of the mode functions. The result is

$$L_{gc}^{(4)} = \sum_n W_{\mu}^{(n)} \left\{ \frac{g_L}{\sqrt{2}} W_{L2}^{(0)} \gamma^\mu \psi_{L1}^{(0)} + \frac{g_R}{\sqrt{2}} W_{R2}^{(0)} \gamma^\mu \psi_{R1}^{(0)} + \text{h.c.} \right\}$$

$$+ \sum_n Z_{\mu}^{(n)} \sum_{i=1}^2 \left\{ g_L Z_{Li}^{(n)} \psi_{Li}^{(0)} \gamma^\mu \psi_{Li}^{(0)} + g_R Z_{Ri}^{(n)} \psi_{Ri}^{(0)} \gamma^\mu \psi_{Ri}^{(0)} \right\}$$

$$+ \sum_n A_{\mu}^{(n)} \sum_{i=1}^2 \left\{ \psi_{Li}^{(0)} \gamma^\mu \psi_{Li}^{(0)} + \psi_{Ri}^{(0)} \gamma^\mu \psi_{Ri}^{(0)} \right\} + \cdots, \quad (28)$$

where the index $i = 1, 2$ denotes the upper or lower components of $\psi_{L,R}^{(0)}$, and the ellipsis denotes terms involving the massive KK modes. From the approximate expressions of the mode functions (18), (19), (22), and (27) for $c > \frac{1}{2}$, the 4D gauge couplings are found to be

$$g_L^{W(0)} \simeq \frac{g_A}{\sqrt{2} \pi R} \equiv g, \quad (29)$$

$^4$ In the $SU(3)$ model in [28], the $\theta_H$-dependences of $m_f$ and $m_W$ are the same so that the ratio $m_f/m_W$ becomes independent of $\theta_H$ even in the warped spacetime.
\[ g^{Z(0)}_L \approx \frac{(-1)^{i-1} g_A - g_B q_{B-L} s_\phi c_\phi}{2 \sqrt{(1 + s_\phi^2) \pi R}} \approx \frac{g}{\cos \theta_W} \left\{ \left( -1 \right)^{i-1} \frac{1}{2} - q_{EM} \sin^2 \theta_W \right\}, \quad (30) \]

\[ g^{\gamma(0)}_i = e q_{EM}, \quad (31) \]

where \( e \equiv g_A \sin \theta_W / \sqrt{\pi R} = g \sin \theta_W \) is the \( U(1) \) gauge coupling constant and \( q_{EM} \equiv \{( -1 )^{i-1} + q_{B-L} \} / 2 \) is the electromagnetic charge. The relation (24) has been made use of in the second equality in (30). Note that Eqs. (29), (30) and (31) agree with the counterparts in the standard model, and are consistent with the experimental results.

Rigorously speaking, the couplings \( g^{W(0)}_L \) and \( g^{Z(0)}_L \) have small dependence on the parameter \( c \), which results in slight violation of the universality in weak interactions as discussed in Ref. [28]. It was found that there is violation of the \( \mu-e \) universality of \( O(10^{-8}) \).

However, \( g^{W(0)}_R \) and \( g^{Z(0)}_R \) evaluated in a similar manner for the same multiplet \( \Psi \) substantially deviate from the standard model values. For instance, one finds \( g^{W(0)}_R = g(1 - \cos \theta_W) / 2 \), which is unacceptable. This is because the mode functions of the right-handed fermions are localized near the TeV brane for \( c > \frac{1}{2} \). Since KK excited states are also localized near the TeV brane, the mixing with KK excited states becomes strong, causing the deviation.

This implies that left-handed quarks \( (u_L, d_L) \) and right-handed quarks \( (u_R, d_R) \), for instance, cannot be in one single multiplet \( \Psi = (q, Q) \) in (7). Instead one should suppose that \( (u_L, d_L) \) is in \( q_L \) of \( \Psi = (q, Q) \) with \( c > \frac{1}{2} \), whereas \( (u_R, d_R) \) is in \( Q'_R \) of a distinct multiplet \( \Psi' = (q', Q') \) with \( c < -\frac{1}{2} \). Indeed, for \( c < -\frac{1}{2} \), the right-handed fermions are localized on the Planck brane and the mixing effect mentioned above becomes negligible.

The couplings become

\[ g^{W(0)}_R \approx 0, \]
\[ g^{Z(0)}_R \approx -\frac{(-1)^{i-1} g_A s_\phi^2 + g_B q_{B-L} s_\phi c_\phi}{2 \sqrt{(1 + s_\phi^2) \pi R}} \approx -\frac{g}{\cos \theta_W} q_{EM} \sin^2 \theta_W, \quad (32) \]

which agree with those in the standard model. This assignment solves another serious problem associated with fermions localized near the TeV brane. Those fermions have too large couplings to the KK gauge bosons, which may contradict with the current precision measurements.[44]-[46]

Of course there remain additional \( Q_R \) and \( q'_L \) which have light modes. These modes must be made substantially heavy, which can be achieved by having boundary mass terms.
connecting $Q_R$, $q'_L$ and additional boundary fields on the TeV brane. This issue will be discussed in more detail in the next paper [47].

Next let us consider the trilinear couplings among the 4D gauge bosons. From the self-interactions of the 5D gauge fields, the following couplings are induced in the 4D effective theory.

$$L^{(4)} = \left\{ i g^{(1)}_{WWZ} (\partial_\mu W^{(0)}_\nu - \partial_\nu W^{(0)}_\mu)^\dagger W^{(0)}_\mu Z^{(0)}_\nu + \text{h.c.} \right\} + i g^{(2)}_{WWZ} W^{(0)}_\mu W^{(0)}_\nu \left( \partial^{\mu} Z^{(0)}_\nu - \partial^{\nu} Z^{(0)}_\mu \right) + \cdots. \tag{33}$$

The couplings $g^{(1)}_{WWZ}$ and $g^{(2)}_{WWZ}$ are expressed by the overlap integrals of the mode functions as

$$g_{WWZ}^{(1)} = g_{WWZ}^{(2)} = g_A \int_1^{2\pi} \frac{dz}{k_z} \left[ \tilde{h}^{3L}_{A,0} \left( \tilde{h}^{+L}_{A,0} \right)^2 + \frac{1}{2} \left( \tilde{h}^{+L}_{A,0} \right)^2 \right] + \tilde{h}^{3R}_{A,0} \left( \tilde{h}^{+R}_{A,0} \right)^2 + \tilde{h}^{3}_{A,0} \tilde{h}^{+}_{A,0} \left( \tilde{h}^{+L}_{A,0} + \tilde{h}^{+R}_{A,0} \right) \right]. \tag{34}$$

Making use of (18) and (22), one finds that

$$g_{WWZ}^{(1)} = g_{WWZ}^{(2)} \simeq \frac{g_A}{\sqrt{(1 + s^2)\pi R}} \simeq g \cos \theta_W. \tag{35}$$

In the last equality, (21) and (29) have been made use of. These couplings have the same values as those in the standard model. The result is consistent with the data of $e^+e^- \rightarrow W^+W^-$ at LEP2 which indicates the validity of the $WWZ$ coupling in the standard model. In deriving (35), we have neglected corrections suppressed by a factor of $(k\pi R)^{-1} \simeq 1/35$ in conformity with the approximation employed in deriving Eqs. (17) - (24).

The $WWZ$ couplings in the model under investigation agree with those in the standard model within this approximation. Small deviation from the standard model may arise beyond this approximation, which needs to be evaluated numerically. For the process $e^+e^- \rightarrow W^+W^-$ contributions from KK excited states also need to be incorporated.

As the 4D Higgs field is a part of 5D gauge fields, the self-interactions of the 5D gauge fields also determine the couplings of the Higgs boson $\varphi^{(0)}$ to the $W$ or $Z$ bosons in the 4D effective theory

$$L^{(4)} = -\lambda_{WWH} \varphi^{(0)} W^{(0)\mu} W^{(0)}_\mu - \frac{1}{2} \lambda_{ZZH} \varphi^{(0)} Z^{(0)\mu} Z^{(0)}_\mu + \cdots. \tag{36}$$
where

\[ \lambda_{WWH} = g_A k \int_1^{z_\pi} \frac{dz}{z} \bar{h}_{\psi,0}^4 \left\{ \bar{h}_{A,0}^+ \partial_z \left( \bar{h}_{A,0}^+ - \tilde{h}_{A,0}^+ \right) - \partial_z \bar{h}_{A,0}^+ \left( \tilde{h}_{A,0}^+ - \tilde{h}_{A,0}^+ \right) \right\}, \]

\[ \lambda_{ZZH} = g_A k \int_1^{z_\pi} \frac{dz}{z} \bar{h}_{\psi,0}^4 \left\{ \bar{h}_{A,0}^3 \partial_z \left( \bar{h}_{A,0}^3 - \tilde{h}_{A,0}^3 \right) - \partial_z \bar{h}_{A,0}^3 \left( \tilde{h}_{A,0}^3 - \tilde{h}_{A,0}^3 \right) \right\}. \] (37)

With the aid of (17)-(24), these couplings are evaluated to be

\[ \lambda_{WWH} \simeq \frac{g_A \sqrt{k}}{\pi R z_\pi} \sin \theta_H \cos \theta_H \simeq g m_W \cdot p_H \cos \theta_H, \]

\[ \lambda_{ZZH} \simeq \frac{g_A \sqrt{k} (1 + s^2_\phi)}{\pi R z_\pi} \sin \theta_H \cos \theta_H \simeq \frac{g m_Z}{\cos \theta_W} \cdot p_H \cos \theta_H, \] (38)

where \( p_H \equiv \text{sgn} (\sin \theta_H) \). It is seen that these couplings are suppressed, compared with those in the standard model, by a factor \( \cos \theta_H \).

So far we have neglected the \( SU(2)_R \)-breaking in the fermion sector. Since \( SU(2)_R \) is broken at the Planck brane, it is natural to have brane-localized mass terms with \( SU(2)_R \) breaking, which, in turn, alter mass eigenvalues and mode functions of the fermions. We would like to emphasize that the predictions for the gauge couplings (29), (30), (31) and (32) remain robust after such an \( SU(2)_R \) breaking effect is incorporated. The dependence of the gauge couplings on the fermion mode functions are exponentially suppressed as long as they are localized near the Planck brane.\(^5\)

Implications of brane-localized mass terms will be analysed in the separate paper [47]. The presence of \( SU(2)_L \times SU(2)_R \) symmetry in the bulk leads not only to the custodial symmetry in the 4D Higgs interactions but also to right-handed neutrino states. It would be interesting to implement the see-saw mechanism in this gauge-Higgs unification scenario.

The main result of this paper is the prediction of the suppression factor \( \cos \theta_H \) for \( \lambda_{WWH} \) and \( \lambda_{ZZH} \). The \( WWZ \) couplings remain the standard model values. All of these couplings can be measured at LHC and future linear colliders. They will certainly give crucial information about the mechanism of the symmetry breaking in the electroweak interactions.

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\(^5\) In this regard quantum corrections from the top quark may be important, and need further investigation.
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References

[1] Y. Okada, M. Yamaguchi, and T. Yanagida, Prog. Theoret. Phys. 85 (1991) 1.

[2] LEP Electroweak Working Group, http://lepewwg.web.cern.ch/lepewwg; The Higgs Working Group at Snowmass ’05, hep-ph/0511332.

[3] N. Arkani-Hamed, A.G. Cohen and H. Georgi, Phys. Lett. B513 (2001) 232; N. Arkani-Hamed, A.G. Cohen, E. Katz and A.E. Nelson, JHEP 0207 (2002) 034; M. Schmaltz and D. Tucker-Smith, Ann. Rev. Nucl. Part. Sci.55 (2005) 229; M. Perelstein, hep-ph/0512128.

[4] C. Csaki, C. Grojean, H. Murayama, L. Pilo, and J. Terning, Phys. Rev. D69 (2004) 055006.

[5] C. Csaki, C. Grojean, L. Pilo, and J. Terning, Phys. Rev. Lett. 92 (2004) 101802; R.S. Chivukula, E.H. Simmons, H.J. He, M. Kurachi and M. Tanabashi, Phys. Rev. D70 (2004) 075008.

[6] D.B. Fairlie, Phys. Lett. B82 (1979) 97; J. Phys. G5 (1979) L55.

[7] N. Manton, Nucl. Phys. B158 (1979) 141; P. Forgacs and N. Manton, Comm. Math. Phys. 72 (1980) 15.

[8] Y. Hosotani, Phys. Lett. B129 (1984) 193; Phys. Rev. D29 (1984) 731.

[9] Y. Hosotani, Phys. Lett. B126 (1983) 309.

[10] Y. Hosotani, Ann. Phys. (N.Y.) 190 (1989) 233.

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

[12] M. Kubo, C.S. Lim and H. Yamashita, Mod. Phys. Lett. A17 (2002) 2249.

[13] A. Pomarol and M. Quiros, Phys. Lett. B438 (1998) 255;

[14] I. Antoniadis, K. Benakli and M. Quiros, New. J. Phys. 3 (2001) 20.

[15] C. Csaki, C. Grojean and H. Murayama, Phys. Rev. D67 (2003) 085012; C.A. Scrucca, M. Serone and L. Silverstrini, Nucl. Phys. B669 (2003) 128.

[16] L.J. Hall, Y. Nomura and D. Smith, Nucl. Phys. B639 (2002) 307; L. Hall, H. Murayama, and Y. Nomura, Nucl. Phys. B645 (2002) 85; G. Burdman and Y. Nomura, Nucl. Phys. B656 (2003) 3; C.A. Scrucca, M. Serone, L. Silvestrini and A. Wulzer, JHEP 0402 (2004) 49; G. Panico and M. Serone, hep-ph/0502255.

[17] N. Haba, Y. Hosotani, Y. Kawamura and T. Yamashita, Phys. Rev. D70 (2004) 015010; N. Haba, K. Takenaga, and T. Yamashita, Phys. Lett. B605 (2005) 355.

[18] N. Haba, K. Takenaga, and T. Yamashita, Phys. Lett. B615 (2005) 247.

[19] Y. Hosotani, S. Noda and K. Takenaga, Phys. Lett. B607 (2005) 276.

[20] G. Cacciapaglia, C. Csaki and S.C. Park, hep-ph/0510366.

[21] G. Panico, M. Serone and A. Wulzer, hep-ph/0510373.

[22] B. Grzadkowski and J. Wudka, hep-ph/0604225.

[23] R. Contino, Y. Nomura and A. Pomarol, Nucl. Phys. B671 (2003) 148.

[24] K. Agashe, R. Contino and A. Pomarol, Nucl. Phys. B719 (2005) 165;
K. Agashe and R. Contino, hep-ph/0510164.

[25] L. Hall, Y. Nomura, T. Okui and S. Oliver, Nucl. Phys. B677 (2004) 87.
[26] K. Oda and A. Weiler, Phys. Lett. B606 (2005) 408.
[27] Y. Hosotani and M. Mabe, Phys. Lett. B615 (2005) 257.
[28] Y. Hosotani, S. Noda, Y. Sakamura and S. Shimasaki, Phys. Rev. D73 (2006) 096006.
[29] T. Gherghetta, hep-ph/0601213.
[30] M. Carena, E. Ponton, J. Santiago and C.E.M. Wagner, hep-ph/0607106.
[31] A. Aranda and J.L. Diaz-Cruz, Phys. Lett. B633 (2006) 591.
[32] G.v. Gersdorff, N. Irges and M. Quiros, Nucl. Phys. B635 (2002) 127.
[33] Y. Hosotani, in the Proceedings of “Dynamical Symmetry Breaking”, ed. M. Harada and K. Yamawaki (Nagoya University, 2004), p. 17. (hep-ph/0504272).
[34] T. Morris, JHEP 0501 (2005) 002.
[35] N. Irges and F. Knechtli, Nucl. Phys. B719 (2005) 121; hep-lat/0604006.
[36] N. Maru and T. Yamashita, hep-ph/0603237.
[37] Y. Hosotani, hep-ph/0607064.
[38] Y. Hosotani, K. Oda and Y. Sakamura, in preparation.
[39] L. Randall and R. Sundrum, Phys. Rev. Lett. 83 (1999) 3370.
[40] T. Gherghetta and A. Pomarol, Nucl. Phys. B586 (2000) 141.
[41] N. Haba, M. Harada, Y. Hosotani and Y. Kawamura, Nucl. Phys. B657 (2003) 169; Erratum, ibid. B669 (2003) 381.
[42] N. Sakai and N. Uekusa, hep-th/0604121.
[43] R.S. Chivukula, D.A. Dicus, and H.-J. He, Phys. Lett. B525 (2002) 175.
[44] S. Chang, J. Hisano, H. Nakano, N. Okada and M. Yamaguchi, Phys. Rev. D62 (2000) 084025.
[45] K. Agashe, A. Delgado, M.J. May and R. Sundrum, JHEP 0308 (2003) 050.
[46] S.J. Huber, Nucl. Phys. B666 (2003) 269; K. Agashe, G. Perez and A. Soni, Phys. Rev. D71 (2005) 016002.
[47] Y. Sakamura and Y. Hosotani, in preparation.