Higgs Mass in the Gauge-Higgs Unification Theory *

N. Haba
University of Tokushima, Tokushima 770-8502, JAPAN

The gauge-Higgs unification theory identifies the zero mode of the extra dimensional component of the gauge field as the usual Higgs doublet. Since this degree of freedom is the Wilson line phase, the Higgs does not have the mass term nor quartic coupling at the tree level. Through quantum corrections, the Higgs can take a vacuum expectation value, and its mass is induced. The radiatively induced mass tends to be small, although it can be lifted to $\mathcal{O}(100)$ GeV by introducing extra bulk fields. We analyze the Higgs mass in SUSY and non-SUSY $SU(3)_c \times SU(3)_W$ and $SU(6)$ models. We also show the case of introducing the soft SUSY breaking scalar masses in addition to the Scherk-Schwarz SUSY breaking.

1. GAUGE-HIGGS UNIFICATION

One of the most fascinating motivations for the higher dimensional gauge theory is that gauge and Higgs fields can be unified[2]. The higher dimensional components of gauge fields become scalar fields bellow the compactification scale, and these scalar fields are identified with the Higgs fields in the gauge-Higgs unification theory. The Higgs doublet fields can appear through the orbifold compactification such as S^1/Z_2 from 5D theory[3]-[7]. In the gauge-Higgs unification theory, the Higgs fields have only finite masses of order the compactification scale because the masses of the Higgs fields are forbidden by the higher dimensional gauge invariance. Yukawa interactions can be induced from the 5D gauge interactions when quarks and leptons are in the bulk. In order to obtain the Higgs doublets from the gauge fields in higher dimensions, the gauge group must be lager than the standard model (SM) gauge group. The Higgs doublets are identified as the zero modes of the extra dimensional component of the 5D gauge field, A_5. Here, we study $SU(3)_c \times SU(3)_W$ and $SU(6)$ models, and check whether the dynamical electro-weak symmetry breaking is possible or not by calculating one loop effective potential of the Higgs doublets[1, 8, 9].

We consider the 5D theory in which the 5th dimensional coordinate is compactified on an S^1/Z_2 orbifold. Denoting y as the coordinate of the 5th dimension, parity operator, P (P') are defined according to the Z_2 transformation, $y \to -y$ $(\pi R - y \to \pi R + y)$. In the $SU(3)_c \times SU(3)_W$ model[4–6], we take $P = P' = \mathrm{diag}(1, -1, -1)$ $(P = P' = I)$ in the base of $SU(3)_W (SU(3)_c)$. Then, there appears Higgs doublet as the zero mode of A_5[1],

$$H = \sqrt{2\pi R}\, A_5. \tag{1}$$

The 4D gauge coupling constant is defined as $g_4 = g/\sqrt{2\pi R}$[2]. The vacuum expectation value (VEV) of A_5 is parameterized by $a/(2gR)\mathbf{E}_3$, where $\mathbf{E}_3$ is the 3×3 matrix having 1 at (1,3) and (3,1) elements, while the other elements being zero[8, 9]. The relation between the VEV and electro-weak scale is given by

$$\sqrt{2\pi R}\,\langle A_5^4 \rangle = \frac{a_0}{g_4 R} = \, v \sim 246\,\mathrm{GeV}. \tag{2}$$

Here the component gauge field A_5^4 is defined by $A_5 = \sum_a A_5^a T^a$ through the generators T^a, where $T^4 = \frac{1}{2}\mathbf{E}_3$. The compactification scale must be above the weak scale, and when we take it as a few TeV, for examples, a_0 should

*This talk is based on Ref.[1] collaborated with T. Yamashita and K. Takenaga.

[1] Taking account of the scalar degrees of freedom in the gauge super multiplet, we can easily show that there appear two Higgs doublets in the SUSY theory.

[2] We should take $g_4 \gtrsim 1$ for the wall-localized kinetic terms being the main part of the MSSM kinetic terms[8]. Hence, we take $g_4 = \mathcal{O}(1)$ in this paper.

be a parameter of $\mathcal{O}(10^{-1\sim-2})$. Since the Higgs is essentially the Wilson line degree of freedom, the mass term nor quartic coupling does not exist in the Higgs potential at the tree level. Through quantum corrections, the Higgs can develop a vacuum expectation value, which means the dynamical electro-weak symmetry breaking is realized and accordingly its mass is induced.

Let us concentrate on SUSY theory with Scherk-Schwarz (SS) SUSY breaking[10]. We define

$$J^{(+)}[a,\beta,n] \equiv \frac{1}{n^5}\left(1-\cos(2\pi n\beta)\right)\cos(\pi n a), \quad J^{(-)}[a,\beta,n] \equiv \frac{1}{n^5}\left(1-\cos(2\pi n\beta)\right)\cos(\pi n(a-1)),$$

where $\beta (0 \leq \beta \leq 0.5)$ parameterizes the magnitude of the SS SUSY breaking. Then, the soft mass parameters become $\mathcal{O}(\beta/R)$[8, 9]. The contribution of the gauge multiplet to the one loop effective potential is written as

$$V_{\rm eff}^{\rm gauge} = -2C\sum_{n=1}^{\infty}\left(J^{(+)}[2a,\beta,n] + 2J^{(+)}[a,\beta,n]\right), \tag{3}$$

where $C \equiv 3/(64\pi^7 R^5)$. The VEV of σ, which forms the real part of scalar component of $N=1$ chiral multiplet at low-energies, becomes zero by calculation of the effective potential for $\langle\sigma\rangle$[11]. The minimum of the effective potential (3) is located at $a_0 = 1$ (mod 2), which means that the suitable electro-weak scale VEV, $(0 <) a_0 \ll 1$ and electro-weak symmetry breaking are not realized. Thus, for the desirable dynamical electro-weak symmetry breaking, one needs to introduce the extra bulk fields, which are $N_{fnd.}^{(\pm)}$ and $N_{adj.}^{(\pm)}$ species of hypermultiplets of fundamental and adjoint representations, respectively. Here the index, $(\pm)$, denotes the sign of the *intrinsic* parity of PP' defined in Refs.[8, 9].

The effective potential from the bulk fields is given by

$$\begin{aligned} V_{\rm eff}^{\rm matter} \; = \; 2C\sum_{n=1}^{\infty}\Big\{ & N_{adj.}^{(+)}\left(J^{(+)}[2a,\beta,n]+2J^{(+)}[a,\beta,n]\right) + N_{adj.}^{(-)}\left(J^{(-)}[2a,\beta,n]+2J^{(-)}[a,\beta,n]\right) \\ & + N_{fnd.}^{(+)}J^{(+)}[a,\beta,n] + N_{fnd.}^{(-)}J^{(-)}[a,\beta,n]\Big\}. \end{aligned} \tag{4}$$

Reference [8] shows one example, $N_{adj.}^{(+)} = N_{adj.}^{(-)} = 2$, $N_{fnd.}^{(-)} = 4$, $N_{fnd.}^{(+)} = 0$ with $\beta = 0.1$ and R^{-1} of order a few TeV, in which the suitable electro-weak symmetry breaking is realized by the small VEV, $a_0 = 0.047$. The Higgs mass is calculate by the second derivative of the effective potential, $V_{\rm eff} \equiv C\bar{V}_{\rm eff} = V_{\rm eff}^{\rm gauge} + V_{\rm eff}^{\rm matter}$ with respect to a at the minimum, $a = a_0$,

$$m_H \sim \frac{\sqrt{3}}{4\pi^3}\left(\frac{\partial^2 \bar{V}_{\rm eff}}{\partial a^2}\right)^{1/2}_{a=a_0} \times \frac{v g_4^2}{a_0}, \tag{5}$$

where we have used (2). In this case Higgs mass is calculated as

$$m_H^2 \sim \left(\frac{0.025\; g_4}{R}\right)^2 \sim (118\; g_4^2\;{\rm GeV})^2, \tag{6}$$

where $g_4 = \mathcal{O}(1)$, as explained above. The Higgs mass is likely to be smaller than the weak scale, 246 GeV (Eq.(2)) since it is zero at the tree level and is induced through the radiative corrections (Coleman-Weinberg mechanism[12]).

As for the $SU(6)$ model[5, 6], we take the parities, $P = {\rm diag}(1,1,1,1,-1,-1)$ and $P' = {\rm diag}(1,-1,-1,-1,-1,-1)$, which induces Higgs doublet in A_5 as the zero mode. The VEV of A_5 is written as $a/(2gR)\mathbf{E}_6$, where $\mathbf{E}_6$ is the 6×6 matrix having 1 at (1,6) and (6,1) elements while the other elements being zero[8, 9]. The gauge part of the effective potential is given by

$$V_{\rm eff}^{\rm gauge} = -2C\sum_{n=1}^{\infty}\left(J^{(+)}[2a,\beta,n] + 2J^{(+)}[a,\beta,n] + 6J^{(-)}[a,\beta,n]\right). \tag{7}$$

As in the $SU(3)_c \times SU(3)_W$ model, the suitable symmetry breaking can not be realized only by the gauge sector. This situation can be changed by introducing the extra bulk fields, which induce the effective potential,

$$V_{\rm eff}^{\rm matter} = 2C \sum_{n=1}^{\infty} \Big\{ N_{adj.}^{(+)} \left(J^{(+)}[2a,\beta,n] + 2J^{(+)}[a,\beta,n] + 6J^{(-)}[a,\beta,n] \right) \\ + N_{adj.}^{(-)} \left(J^{(-)}[2a,\beta,n] + 2J^{(-)}[a,\beta,n] + 6J^{(+)}[a,\beta,n] \right) + N_{fnd.}^{(+)} J^{(+)}[a,\beta,n] + N_{fnd.}^{(-)} J^{(-)}[a,\beta,n] \Big\}. \quad (8)$$

We show one example of the suitable symmetry breaking in Ref.[8], which is the case of $N_{adj.}^{(+)} = 2$, $N_{fnd.}^{(-)} = 10$, $N_{adj.}^{(-)} = N_{fnd.}^{(+)} = 0$ with $\beta = 0.1$ and R^{-1} of order a few TeV. In this case, the minimum exists at $a_0 = 0.047$, and the Higgs mass squared is calculated as

$$m_H^2 \sim \left(\frac{0.024\ g_4}{R} \right)^2 \sim (120\ g_4^2\ \mathrm{GeV})^2. \quad (9)$$

In the above two examples $\mathcal{O}(10)$ numbers of bulk fields are required for the suitable symmetry breaking and Higgs mass. However, is it inevitable situation? To see it, let us estimate the coefficients of a^2 and a^4 in the effective potential.

2. HIGGS MASS

The Higgs mass is defined by the 2nd derivative of the effective potential. For $a \ll \beta^3$, the coefficients of a^2 and a^4 in $J^{(+)}[a/\pi, \beta/\pi, n]$ are given by[1]

$$-\frac{\beta^2}{288}(432 - 144\ln(4\beta^2))\ \ (<0), \quad \text{and} \quad \frac{25 - 6\ln(a_0^2/4\beta^2)}{288}\ \ (>0), \quad (10)$$

respectively. On the other hand, the coefficients of a^2 and a^4 in $J^{(-)}[a/\pi, \beta/\pi, n]$ become

$$\beta^2 \ln 2\ \ (>0), \quad \text{and} \quad -\beta^2/48\ \ (<0), \quad (11)$$

respectively. For realizing the suitable heavy Higgs mass, the quartic coupling should be large and positive. On the other hand, the VEV (W and Z boson masses) should be maintained small ($a_0 \ll 1$) comparing to the compactification scale. For this purpose, large negative contribution in the first term in Eq.(10) must be almost canceled by introducing $N_{fnd.}^{(-)} = \mathcal{O}(\ln 4\beta^2)$ numbers bulk fields acting on the first term in Eq.(11). This means that the less (more) bulk fields are needed when β becomes large (small). Eqs.(10) and (11) shows that even in the case of this cancellation, the coefficient of a^4 is still positive and large enough when $a_0 \ll \beta$. Thus, the heaviness of Higgs mass is mainly controlled by the factor, $-\ln(a_0^2/\beta^2)$, in the effective quartic coupling, which implies that the smaller (larger) a_0^2/β^2 becomes, the larger (smaller) the Higgs mass becomes. For the $SU(3)_c \times SU(3)_W$ model, the Higgs mass is calculated as

$$\frac{m_H}{g_4^2} \simeq v\,\frac{\sqrt{3}}{4\pi}\sqrt{4B\,\ln\left(\frac{a_0^2}{4\beta^2}\right) + \text{const.}}, \qquad B \equiv \frac{-1}{24}\left(18(N_{adj.}^{(+)} - 1) + N_{fnd.}^{(+)}\right), \quad (12)$$

and the constant term depends on β and the number of flavors. Equation (12) shows that a few adjoint bulk fields are enough and essential for the large quartic coupling. The contribution from the adjoint bulk field overcome the loop factor $\sim 1/4\pi$ to enhance the magnitude of the Higgs mass. Hence, $\mathcal{O}(1)$ numbers of bulk fields can realize the

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[3] In the usual scenario, $a < \beta$ should be satisfied since the SUSY breaking mass, $\mathcal{O}(\beta/R)$ must be larger than the electro-weak scale, $\mathcal{O}(a/R)$.

suitable symmetry breaking and Higgs mass of $\mathcal{O}(100)$ GeV as shown in Ref.[1]. We should note that the dependence of the Higgs mass on the supersymmetry breaking parameter is logarithmic, as expected.

In non-SUSY models, it is also possible to cancel the coefficient of a^2 terms between the bulk fields with even parity and the one with odd parity, keeping the positive and large quartic coupling, by an appropriate choice of the matter content. This case has the similar situation as the SUSY case studied above, and we can have the desirable size of the Higgs mass. The non-SUSY model with the appropriate matter content, which realize the suitable dynamical electro-weak symmetry breaking, is presented in Ref.[8].

3. SUSY GAUGE-HIGGS WITH BULK MASS

Next, we show another example for realizing the dynamical electro-weak symmetry breaking with the small number of extra bulk fields. We introduce explicit soft SUSY breaking scalar mass in addition to the SS parameter for the bulk superfields. We do not introduce the soft gaugino masses because the mass terms are odd under the Z_2 operation.

For $SU(3)_c \times SU(3)_W$ model, the contribution of the gauge multiplet to the effective potential is the same as Eq.(3). We introduce the soft SUSY breaking mass, m for the bulk hypermultiplets and define a dimensionless parameter, $z \equiv mR$ (< 1). We denote

$$I^{(+)}[a,\beta,z,n] \equiv \frac{1}{n^5}\left(1-\left(1+2\pi zn+\frac{(2\pi zn)^2}{3}\right)e^{-2\pi zn}\cos(2\pi n\beta)\right)\times\cos(\pi na), \tag{13}$$

$$I^{(-)}[a,\beta,z,n] \equiv \frac{1}{n^5}\left(1-\left(1+2\pi zn+\frac{(2\pi zn)^2}{3}\right)e^{-2\pi zn}\cos(2\pi n\beta)\right)\times\cos(\pi n(a-1)), \tag{14}$$

in which $I^{(\pm)}[a,\beta,z,n]$ is reduced to $J^{(\pm)}[a,\beta,n]$ in the limit of $z\to 0$ $(m\to 0)$. The contribution of the matter hypermultiplet to the effective potential is given by[1]

$$\begin{aligned} V_{\rm eff}^{\rm matter} = 2C\sum_{n=1}^{\infty}\Big\{ & N_{adj.}^{(+)}\left(I^{(+)}[2a,\beta,z_{adj.}^{(+)},n]+2I^{(+)}[a,\beta,z_{adj.}^{(+)},n]\right)+N_{adj.}^{(-)}\left(I^{(-)}[2a,\beta,z_{adj.}^{(-)},n]+2I^{(-)}[a,\beta,z_{adj.}^{(-)},n]\right) \\ & +N_{fnd.}^{(+)}I^{(+)}[a,\beta,z_{fnd.}^{(+)},n]+N_{fnd.}^{(-)}I^{(-)}[a,\beta,z_{fnd.}^{(-)},n]\Big\}, \end{aligned} \tag{15}$$

where $z_{rep.}^{(\pm)}$ stands for the explicit soft mass defined by $z_{rep.}^{(\pm)} \equiv m_{rep.}^{(\pm)}R$ (< 1) for each representation field. Eq.(15) becomes Eq.(4) in the limit of the vanishing soft scalar mass, $m\to 0$.

We find some examples of extra matter contents and SUSY breaking parameters, for which the suitable VEV and Higgs mass are realized, and we summarize them in the following table[1].

	$N_{adj.}^{(+)}$	$N_{adj.}^{(-)}$	$N_{fnd.}^{(+)}$	$N_{fnd.}^{(-)}$	β	$z_{adj.}^{(+)}$	$z_{adj.}^{(-)}$	$z_{fnd.}^{(+)}$	$z_{fnd.}^{(-)}$	a_0	m_H/g_4^2
(1)	2	3	0	4	0.05	0.01	0.01	-	0.045	0.0040	164
(2)	2	4	2	6	0.05	0	0	0.05	0.05	0.0037	176
(3)	2	4	0	6	0.025	0.025	0.025	-	0.025	0.0066	129
(4)	2	1	0	2	0.1	0.1	0.1	-	1	0.0097	150
(5)	1	1	0	2	0.01	1	1	-	1	0.0196	125
(6)	2	2	0	2	0.14	0	0	-	0	0.0379	130

The Higgs mass m_H/g_4^2 is measured in GeV unit. This table shows that even small number of extra bulk fields can realize the suitable dynamical electro-weak symmetry breaking with the heavy Higgs mass. The effect of the bulk masses increases not only the degrees of freedom of parameter space, but also induces a similar effect of large β, which is necessary for the symmetry breaking with a small number bulk fields, as explained in the previous section.

In $SU(6)$ model, the contribution of the gauge multiplet to the effective potential is the same as Eq.(7). On the other hand, the contribution of the matter hypermultiplet to the effective potential is given by

$$V_{\rm eff}^{\rm matter} = 2C\sum_{n=1}^{\infty}\Big\{N_{adj.}^{(+)}\left(I^{(+)}[2a,\beta,z_{adj.}^{(+)},n]+2I^{(+)}[a,\beta,z_{adj.}^{(+)},n]+6I^{(-)}[a,\beta,z_{adj.}^{(+)},n]\right)$$

$$+N^{(-)}_{adj.}\left(I^{(-)}[2a,\beta,z^{(-)}_{adj.},n]+2I^{(-)}[a,\beta,z^{(-)}_{adj.},n]+6I^{(+)}[a,\beta,z^{(-)}_{adj.},n]\right)$$
$$+N^{(+)}_{fnd.}I^{(+)}[a,\beta,z^{(+)}_{fnd.},n]+N^{(-)}_{fnd.}I^{(-)}[a,\beta,z^{(-)}_{fnd.},n]\Big\}, \quad (16)$$

which becomes Eq.(8) in the zero limit of explicit soft scalar masses. Some examples for realizing the suitable dynamical electro-weak symmetry breaking are shown in the following table.

	$N^{(+)}_{adj.}$	$N^{(-)}_{adj.}$	$N^{(+)}_{fnd.}$	$N^{(-)}_{fnd.}$	β	$z^{(+)}_{adj.}$	$z^{(-)}_{adj.}$	$z^{(+)}_{fnd.}$	$z^{(-)}_{fnd.}$	a_0	m_H/g_4^2
(7)	2	0	0	10	0.1	0.05	-	-	0.05	0.0207	139
(8)	2	0	0	6	0.15	0.1	-	-	0.1	0.0268	139
(9)	2	0	0	16	0.04	0	-	-	0.03	0.0021	173
(10)	2	0	0	4	0.07	0.5	-	-	0.5	0.0366	138
(11)	2	0	0	2	0.32	0	-	-	0	0.0594	135

We would like to comment on some phenomenological issues relating to Higgs self interactions. For the higher order operators of self interactions, we see that the effective potential contains a^n interactions by the expansion of the cosine function, which implies $a^n = (g_4RH)^n$ from Eqs.(1) and (2). When g_4R is of order a few TeV, higher order operators, H^n ($n \geq 6$) have the dimensionful suppression of order a few TeV. This means that the contributions from the higher order operators are not so significant.

Study of the effective 3-point self coupling of H is important for the search of the new physics in the future linear colliders[13]. The coupling of the effective λH^3 interaction is given by $\lambda \equiv \frac{3g_4^3}{32\pi^6 R}\left.\frac{\partial^3(V/C)}{\partial a^3}\right|_{a_0}$, and the deviation from the tree level SM coupling, $\lambda_{SM} = 3m_h^2/v$, is estimated by $\Delta\lambda = (\lambda - \lambda_{SM})/\lambda_{SM}$ [13]. The value of $\Delta\lambda$ becomes -17.4% for the example of $SU(3)_c \times SU(3)_W$ model ($N^{(+)}_{adj.} = N^{(-)}_{adj.} = 2$, $N^{(-)}_{fnd.} = 4$, $N^{(+)}_{fnd.} = 0$ with $\beta = 0.1$) and -16.6% for $SU(6)$ model ($N^{(+)}_{adj.} = 2$, $N^{(-)}_{fnd.} = 10$, $N^{(-)}_{adj.} = N^{(+)}_{fnd.} = 0$ with $\beta = 0.1$). As for the above examples, we show it in the following table.

	(1)	(2)	(3)	(4)	(5)	(6)	(7)	(8)	(9)	(10)	(11)
$\Delta\lambda(\%)$	-8.6	-8.3	-14.0	-10.2	-3.1	-13.7	-12.0	-12.0	-7.6	-11.2	-12.7

It suggests that the effective 3-point self couplings tend to be small comparing to that of the SM.

We should notice that the Higgs field in our model has VEV in A_5 not σ. The VEV of A_5 should be distinguished from that of σ in the dynamically induced effective potential in the gauge-Higgs unification theory[11]. Since H is the field of the D-flat direction, which is massless at tree level, it corresponds to the lighter Higgs scalar in the MSSM, h^0. Since h^0 becomes the SM-like Higgs in the large soft SUSY breaking masses, we have compared the effective 3-point self coupling to the SM one in the above estimation of $\Delta\lambda$. Other masses of Higgs eigenstates, charged Higgs, neutral scalar, and heavier pseudo-scalar, can be calculated by the effective potential of these directions[14]. Regarding the large quantum corrections as shown above, the Higgs particles may have the quite different mass spectrum than the SM. The value of $\tan\beta$ can be also calculated once the bulk matter content is fixed, which is expected to be small ($\mathcal{O}(1)$) in the gauge-Higgs unification.

The study of gaugino-higgsino mass spectrum at the weak scale may be also interesting, since they have the same mass at the compactification scale (higgsinos is originated from gauginos as Higgs from gauge field).

4. SUMMARY AND DISCUSSIONS

We have studied the possibility of the dynamical electro-weak symmetry breaking in two gauge-Higgs unified models, $SU(3)_c \times SU(3)_W$ and $SU(6)$ models. We calculated the one loop effective potential of Higgs doublets and analyze the vacuum structure of the models. We found that the introduction of the appropriate numbers and representation of extra bulk fields are required for the desirable symmetry breaking. Since the Higgs is essentially the

Wilson line degree of freedom, the mass term nor quartic coupling does not exist in the Higgs potential at the tree level. Through quantum corrections, the Higgs can develop a vacuum expectation value, which means the dynamical electro-weak symmetry breaking is realized and accordingly its mass is induced. The induced Higgs mass tends to be small, less than the weak scale, reflecting the nature of the Coleman-Weinberg mechanism. We have also studied the case of introducing the soft SUSY breaking scalar masses in addition to the SS SUSY breaking. In this case the suitable electro-weak symmetry breaking and the $\mathcal{O}(100)$ GeV Higgs mass can be realized by $\mathcal{O}(1)$ numbers of bulk fields. We have also show the effective 3-point self coupling of H in some matter contents.

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References

[1] N. Haba, K. Takenaga and T. Yamashita, Phys. Lett. B **615** (2005), 247.

[2] N. S. Manton, Nucl. Phys. B **158**, (1979), 141;
Y. Hosotani, Phys. Lett. **B126** (1983), 309.

[3] N. V. Krasnikov, Phys. Lett. B **273**, (1991), 246;
H. Hatanaka, T. Inami and C. S. Lim, Mod. Phys. Lett. A **13**, (1998), 2601;
G. R. Dvali, S. Randjbar-Daemi and R. Tabbash, Phys. Rev. D **65**, (2002), 064021;
N. Arkani-Hamed, A. G. Cohen and H. Georgi, Phys. Lett. B **513**, (2001), 232;
I. Antoniadis, K. Benakli and M. Quiros, New J. Phys. **3**, (2001), 20.

[4] M. Kubo, C. S. Lim and H. Yamashita, Mod. Phys. Lett. A **17** (2002), 2249.

[5] L. J. Hall, Y. Nomura and D. R. Smith, Nucl. Phys. B **639** (2002), 307.

[6] G. Burdman and Y. Nomura, Nucl. Phys. B **656** (2003), 3.

[7] N. Haba and Y. Shimizu, Phys. Rev. D **67** (2003), 095001;
C. A. Scrucca, M. Serone and L. Silvestrini, Nucl. Phys. B **669** (2003), 128;
L. Gogoladze, Y. Mimura and S. Nandi, Phys. Lett. **B560** (2003), 204;
C. Csaki, C. Grojean, H. Murayama, L. Pilo and J. Terning, Phys. Rev. D **69** (2004), 055006.

[8] N. Haba, Y. Hosotani, Y. Kawamura and T. Yamashita, Phys. Rev. D **70** (2004), 015010.

[9] N. Haba and T. Yamashita, JHEP **0402** (2004), 059; JHEP **0404** (2004), 016.

[10] J. Scherk and J. H. Schwarz, Phys. Lett. B **82** (1979), 60; Nucl. Phys. B **153** (1979), 61.

[11] N. Haba, K. Takenaga and T. Yamashita, Phys. Rev. D **71** (2005), 025006.

[12] S. Coleman and E. Weinberg, Phys. Rev. D **7** (1973), 1888.

[13] S. Kanemura, S. Kiyoura, Y. Okada, E. Senaha and C. P. Yuan, Phys. Lett. B **558** (2003), 157;
S. Kanemura, Y. Okada, E. Senaha and C. P. Yuan, Phys. Rev. D **70** (2004), 115002.

[14] N. Haba, K. Takenaga and T. Yamashita, *in preparation.*