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Radiative electroweak symmetry breaking in the extra dimensions scenarios

X.-S. Li[1], Q.-S. Yan[2,a]

[1] Physics Department of Tsinghua University, 100084 Beijing, P.R. China
[2] IHEP of CAS, P.O. Box 918(4), 100039 Beijing, P.R. China

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Abstract. We study radiative spontaneous electroweak symmetry breaking in the non-supersymmetric extra dimension scenarios of the standard model extension proposed by Antoniadis et al., Dienes et al. and Pomarol et al. In the framework of the multi-scale effective theory, by using the renormalization group method with an up-to-down viewpoint, we find that the effects of Kaluza–Klein excitations of bosons of the standard model can change the sign of the Higgs mass term of the standard model from positive to negative and break the electroweak symmetry. The critical scale for the electroweak phase transition to occur depends on the compactification scale (say 1.6 (2.0) TeV if the compactification scale is assumed to be 0.8 (1.5) TeV or so), and is insensitive to the mass of the Higgs particle. This radiative spontaneous symmetry breaking mechanism can work naturally in the extra dimension scenarios, and neither new particle contents beyond the standard model from the supersymmetry nor technicolor are necessary.

The extra dimension (ED) scenario of the standard model (SM) extension provides a new theoretical paradigm to solve the hierarchy problem of the SM beyond the two traditional solutions: technicolor (TC) and supersymmetry (SUSY). More phenomenological implications of the idea are under vigorous investigation (refer to [1] for a brief review). As one of its recent important developments, much work has been done [2] on model construction in the ED. However, a realistic ED model should explain, in its own framework, how the electroweak symmetry is broken, a problem under intensive research at present.

In the SM, to trigger the spontaneous electroweak symmetry breaking (SEWSB), the square mass of the Higgs in the potential is set by hand to be negative. As we know, the Higgs mechanism is quite successful in explaining the masses of bosons and fermions and their mixing, but it is to some degree unnatural in this respect. The underlying mechanism for the SEWSB of the SM is still an open question [3].

There are two main ways to explain the SEWSB of the SM: the TC and the SUSY. In TC models, the Higgs field is regarded as a composite field of fundamental fields. The SEWSB occurs due to the fermion condensate. The SEWSB of the TC version in the ED has been intensively investigated (refer to [4] for a review). In SUSY models, the radiative symmetry breaking mechanism [5] is well known [6]. In an up-to-down viewpoint, due to the large values of the top Yukawa couplings, the square mass of one of the Higgs doublets which couples to u-type quarks can be driven from positive to negative and the SEWSB can be radiatively triggered in a natural way. However, we would like to emphasize that it is those massive bosonic superpartners (especially the superpartners of the top quark) that contribute constructively to the Higgs potential and trigger the SEWSB. The SEWSB of the SUSY version in the ED has also been studied [7].

In the non-supersymmetric (NSUSY) ED scenarios, there are plentiful massive bosonic states in the reduced 4D effective theory (for example, the Kaluza–Klein (KK) excitations of both vector and Higgs boson could be quite heavy). Then a natural question arises: is it possible for these massive bosons to induce the desired SEWSB at a few TeV, just as those massive bosonic superpartners do in SUSY models? In this work, we will answer this question and study the radiative mechanism in the NSUSY ED extension of the SM proposed by Antoniadis et al. [8], Dienes et al. [9] and Pomarol et al. [10], where both gauge and Higgs bosons are assumed to propagate in the bulk. We find that in this NSUSY extra dimension models, with an appropriate assumption on the mass of the Higgs at the low energy scale (M_Z), the radiative correction of KK excitations to the Higgs potential can naturally trigger the desired electroweak symmetry breaking at a few TeV. Neither new particle contents beyond the SM due to SUSY nor TC are necessary.

In order to deal with the effects of those massive KK modes to the problem we study here, we will use the idea of the multi-mass-scale effective potential method (MEPM) given in [11,12] to derive the RGEs. In a theory with more than one scale, MEPM can avoid large logarithms and preserves the validity of perturbation theory. The basic ingredients of the MEPM include the renormalization group equation (RGE) method and the decoupling theorem [14]. One of the advantages of the MEPM is that by using dif-

[a] e-mail: yanqs@mail.ihep.ac.cn

ferent effective field theories [13] which are dependent on the field scales, it is convenient to solve the RGEs of the dynamic system.

To understand how those massive bosonic KK excitations can trigger the spontaneous symmetry breaking (SSB), we examine how this happens in a two-real-scalar system. The full Lagrangian of the two-real-scalar system defined at the high energy scale $\Lambda_{\rm UV}$, is assumed to have the form

$$L_{\rm UV} = \frac{1}{2}\sum_{i=1}^{2}(\partial^\nu H_i)^\dagger \partial_\nu H_i - V(H_1, H_2), \quad (1)$$

where H_i, $i = 1, 2$, are fields of two real scalars.

Motivated by the effective 4D Lagrangian reduced from the 5D one, the potential $V(H_1, H_2)$ is simply assumed to have the form

$$V(H_1, H_2) = \sum_{i=1}^{2} m^2_{H_i}(\mu_{\rm UV}) H_i^\dagger H_i + \frac{\lambda(\mu_{\rm UV})}{4!}(H_1^4 + 6H_1^2 H_2^2 + H_2^4); \quad (2)$$

$m^2_{H_1}$ is assumed to be much lighter than $m^2_{H_2}$, while $\lambda(\Lambda_{\rm UV})$ is positive in order to have a stable vacuum. This Lagrangian is invariant under a $Z_2 \times Z_2$ transformation $H_1 \to -H_1$ and $H_2 \to -H_2$. At $\Lambda_{\rm UV}$, $m^2_{H_i}(\Lambda_{\rm UV})$, $i = 1, 2$, and $\lambda(\Lambda_{\rm UV})$ are the three free parameters for the system, and $m^2_{H_i}(\mu_{\rm UV})$ are assumed to be positive and the symmetry $Z_2 \times Z_2$ is assumed to be unbroken[1].

At the infrared boundary of the system $\Lambda_{\rm IR}$, the heavy degree of freedom H_2 should decouple[2], the low energy effective Lagrangian should contain only the light degree of freedom of the theory and can be written

$$L_{\rm eff} = \frac{1}{2}(\partial^\nu H_1)^\dagger \partial_\nu H_1 - V(H_1) + \ldots, \quad (3)$$

where the dots represent the omitted irrelevant terms, and the effective potential $V(H_1)$ can simply be expressed as

$$V(H_1) = m^2_{H_1}(\Lambda_{\rm IR}) H_1^\dagger H_1 + \frac{\lambda_1(\Lambda_{\rm IR})}{4!} H_1^4. \quad (4)$$

The low energy effective Lagrangian has a Z_2 reflection symmetry.

We can derive the one-loop RGEs of the system directly with the method given in [11]. The one-loop RGEs of this system are listed below:

$$\frac{{\rm d}m^2_{H_1}(t)}{{\rm d}t} = \frac{\alpha_\lambda}{4\pi}[m^2_{H_1}(t) + m^2_{H_2}(t)\theta(\mu - m_{H_2})], \quad (5)$$

$$\frac{{\rm d}m^2_{H_2}(t)}{{\rm d}t} = \frac{\alpha_\lambda}{4\pi}[m^2_{H_1}(t) + m^2_{H_2}(t)]\theta(\mu - m_{H_2}), \quad (6)$$

$$\frac{{\rm d}\alpha_\lambda}{{\rm d}t} = \frac{\alpha_\lambda}{4\pi}[3 + 3\theta(\mu - m_{H_2})], \quad (7)$$

where $t = \ln(\mu/\Lambda_{\rm IR})$, $\alpha_\lambda = \lambda/(4\pi)$.

At the low energy region near $\Lambda_{\rm IR}$, we just assume that $m^2_{H_1}(\Lambda_{\rm IR})$ is negative and the Z_2 symmetry is spontaneously broken, to simulate the case in the SM. At the low energy regions, v (i.e. the vacuum expectation value (VEV) of the field H_1), $\alpha_\lambda(\Lambda_{\rm IR})$ and $m^2_{H_2}$ can be chosen as the free parameters of the system. The $m^2_{H_2}$ is not defined at $\Lambda_{\rm IR}$ and should be larger than $\Lambda_{\rm IR}$; however, it is regarded as one of the low energy free parameters to control when the degree freedom of the field H_2 should be considered. These parameters defined at the low energy region will be used as the lower boundary conditions to solve the RGEs. The initial value of $m^2_{H_1}(\Lambda_{\rm IR})$ is fixed by the relation $m^2_{H_1}(\Lambda_{\rm IR}) = -2\pi/3\alpha_\lambda(\Lambda_{\rm IR})v^2$.

Now we study the behavior of the RGE of $m^2_{H_1}$. At the energy scale region below the threshold value of m_{H_2}, the running of the RGEs is quite simple: the self-coupling λ of the Higgs potential will become stronger with the increasing of the energy scale μ, and the value of the mass term $m^2_{H_1}$ will simply be driven to larger values and there is no hope to change its sign. However, when the running scale runs across the threshold value of m_{H_2} and the degree of freedom of the field H_2 is activated, something interesting happens. We know from (5) that if the condition $|m^2_{H_1}(\mu)| < |m^2_{H_2}(\mu)|$ is satisfied for $\mu \geq m_{H_2}$, it is sufficient for the existence of a scale $\mu_{\rm cri}$ where $m^2_{H_1}$ vanishes. Above the scale $\mu_{\rm cri}$, $m^2_{H_1}$ changes from negative to positive and the VEV of the H_1 vanishes; then we see the broken Z_2 symmetry is restored. This sufficient condition[3] can be solved from (5)–(7):

$$\left|\frac{2\pi}{3}\alpha_\lambda(\Lambda_{\rm IR})v^2\right| < m^2_{H_2}\left(1 - \frac{\alpha_\lambda(\Lambda_{\rm IR})}{4\pi}\ln\frac{m^2_{H_2}}{\Lambda^2_{\rm IR}}\right)^{1/3}. \quad (8)$$

We can view this phenomenon in a more instructive way: With the fixed v, and m_{H_2} and α_λ which satisfy the condition given in (8) as the lower boundary condition of the RGEs, we can get the corresponding upper boundary value defined at the high energy $\Lambda_{\rm UV}$ and parameterized with ($m^2_{H_i}(\Lambda_{\rm UV})$, $i = 1, 2$ and $\alpha_\lambda(\Lambda_{\rm UV})$) by solving the RGEs. With this determined upper boundary condition of the RGEs and running these parameters from $\Lambda_{\rm UV}$ down to the lower energy region, we can definitely get the desired radiative spontaneous symmetry breaking at the same critical scale $\mu_{\rm cri}$, since the RGEs are ordinary differential equations and are uniquely solvable with a specified boundary condition, either the lower or upper one.

Equipped with the experience we got from the above toy model, we will examine the case in the ED scenarios. We consider a simple extension of the SM to 5D [10] (it is straightforward to generalize our discuss to cases with a high number of extra dimensions). The fifth space-like dimension x_5 is assumed to compactify on the orbifold S^1/Z_2. The 5D Lagrangian is defined by

$$\mathcal{L}_{\rm 5D} = -\frac{1}{4}F^2_{MN} + |D_M H|^2 - V(H)$$

[1] The reason for a positive value of the Higgs square mass at the ultraviolet cut-off scale might be in some more fundamental theory, say supersymmetry breaking

[2] Although H_2 should decouple at $M_{H_2} = (m^2_{H_2} + \lambda H_1^2/2)^{1/2}$ as stated in [12], due to the large mass of KK excitations, the difference between m_{H_2} and M_{H_2} is omitted here

[3] The mass of the heavier Higgs should not be too large to break down the perturbation theory

$$+L_{\mathrm{GF}}+\left[\mathrm{i}\bar{\psi}_i\sigma^\mu D_\mu\psi_i+Y_F\bar{\psi}_{\mathrm{L}}H\psi_{\mathrm{R}}\right]\delta(x_5), \quad (9)$$

where $F_{MN}=\partial_M A_N-\partial_N A_M+g_5[A_M,A_N]$, which is the gauge field tensor defined in 5D, and $N,M=0,1,2,3,5$. The group generators' index is omitted in the above convention. Gauge fields A_M and the Higgs weak doublet field H have mass dimension 3/2. g_5 and Y_F are the gauge coupling and the Yukawa coupling, respectively, and have mass dimension $-1/2$. $\mathcal{L}_{\mathrm{GF}}$ is the gauge fixed term. $V(H)$ is the usual Higgs potential and has the form

$$V(H)=\mu^2H^\dagger H+\frac{\lambda}{4}(H^\dagger H)^2, \quad (10)$$

μ^2 and λ are the mass term and self-coupling of the H boson, respectively. λ has mass dimension -1. Here μ^2 is assumed to be positive and the $SU(2)\times U(1)$ symmetry in 5D is assumed to be preserved even when compactification occurs.

According to the power law, the gauge theory defined in 5D is non-renormalizable, due to the fact that the couplings have a negative mass dimension. So the Lagrangian in (9) can only be valid below an explicit ultraviolet cut-off Λ_{UV}. Besides its ultraviolet cut-off, it has a natural infrared cut-off $\Lambda_{\mathrm{IR}}^{5\mathrm{D}}$ ($\Lambda_{\mathrm{IR}}^{5\mathrm{D}}$ is near the compactification scale of the fifth dimension), below which the 5D description breaks down and above which the 5D description is proper.

In order to study the effective 5D theory with minimal interaction terms near its infrared cut-off $\Lambda_{\mathrm{IR}}^{5\mathrm{D}}$, we expect to find its 4D effective description. Dimension reduction, field and parameter rescaling and matching procedure are invoked to get the 4D effective theory. These procedures are quite standard in the literature.

The fields living in the bulk are defined to be even under the Z_2-parity transformation and then can be Fourier-expanded as

$$A_M(H)(x_\mu,x_5)=\sum_{n=0}^{\infty}\cos\frac{nx_5}{R_c}A^n_{M(5\mathrm{D})}(H^n_{(5\mathrm{D})})(x_\mu), \quad (11)$$

where R_c is the compactification size of the fifth dimension, and $A^{(n)}_{M(5\mathrm{D})}(H^n_{(5\mathrm{D})}), n\neq 0$, are KK excitations. Zero modes are localized on the 3-brane and are fields defined in the SM. Substituting (11) into (9), and rescaling fields and parameters with the following relations:

$$\begin{aligned}\lambda_{5\mathrm{D}}&=2\pi R\lambda_{4\mathrm{D}},\\ g_{5\mathrm{D}}(Y_{u5\mathrm{D}})&=\sqrt{2\pi R}g_{4\mathrm{D}}(Y_{u4\mathrm{D}}),\\ A^0_{\mu5\mathrm{D}}(H^0_{5\mathrm{D}})&=A^0_{\mu4\mathrm{D}}(H^0_{4\mathrm{D}})/\sqrt{2\pi R},\\ A^n_{\mu5\mathrm{D}}(H^n_{5\mathrm{D}},A^n_{55\mathrm{D}})&=A^n_{\mu4\mathrm{D}}(H^n_{4\mathrm{D}},A^n_{54\mathrm{D}})/\sqrt{\pi R}\quad(n\neq0),\end{aligned}$$

then we get the effective Lagrangian in 4D defined at the ultraviolet cut-off Λ_{UV}, which has the form

$$\mathcal{L}^{\mathrm{eff}}=\mathcal{L}_{\mathrm{SM}}+\delta\mathcal{L}^{\mathrm{ED}}, \quad (12)$$

where $\mathcal{L}_{\mathrm{SM}}$ is just the Lagrangian of the SM in 4D, and $\delta\mathcal{L}^{\mathrm{ED}}$ contains all interactions of KK excitations with zero modes on the 3-brane. The effective theory has a $SU_c(3)\times SU_{\mathrm{L}}(2)\times U_Y(1)$ symmetry. In order to be compatible with the SM, A_5 (the fifth component of the vector boson field) is assumed to have no zero mode.

Two features of the model are remarkable [9]. The first one is the universal KK excitation spectra for fields propagating in the bulk, i.e. the masses of KK excitations are independent of other quantum numbers (say, spin and charge) and can be universally expressed as $m_n=n/R_c$. The second one is that, after rescaling and redefining fields and couplings, there is no coupling with negative mass dimension; the Lagrangian could be called a renormalizable one. But it is the presence of infinite towers of KK states that makes the effective theory non-renormalizable.

The matching procedure transforms the non-renormalizable effective Lagrangian to the renormalizable effective Lagrangian L^{eff} by introducing a natural truncation on the infinite KK towers. The matching condition requires that the matching scale $\Lambda^{4\mathrm{D}}_{\mathrm{UV}}$ of the L^{eff} and $L_{4\mathrm{D}}$ with infinite KK towers should fall into the region $\Lambda^{5\mathrm{D}}_{\mathrm{IR}}<\Lambda^{4\mathrm{D}}_{\mathrm{UV}}<\Lambda^{5\mathrm{D}}_{\mathrm{UV}}$. Those KK excitations that have masses larger than $\Lambda^{4\mathrm{D}}_{\mathrm{UV}}$ are excluded, and only those with masses lighter than $\Lambda^{4\mathrm{D}}_{\mathrm{UV}}$ are left in L^{eff}. This Lagrangian, L^{eff}, having finite particles and finite interactions, and couplings with non-negative mass dimension, according to the power law, is renormalizable.

Considering that there are several mass scales in the renormalizable L^{eff} and that in the dimension regularization symmetries are preserved, we can use the idea of the MEPM to calculate RGEs.

There are two procedures to study the 4D effective theory we get: the up-to-down and down-to-up ones. For the up-to-down procedure, by using the decoupling theorem and integrating out those heavy KK modes successively, we can obtain a descending sequence of effective theories, each one with fewer KK modes. At the low energy scale M_Z, all KK excitations decouple and only the SM is left. For the down-to-up procedure, from the infrared scale M_Z up to the ultraviolet scale Λ_{UV}, we can obtain an ascending sequence of effective theories, each one with more activated KK modes. However, for both procedures, the RGEs are just the same and the dynamic behavior of the system (say, the symmetry breaking) revealed by the RGEs is independent of which viewpoint is taken. Below we will take the down-to-up viewpoint, due to the reason that it is easy to determine the lower boundary condition of RGEs directly from experiments at M_Z and a few free parameters of KK excitations.

Up to one-loop level, the RGEs of gauge couplings take the form $\mathrm{d}\alpha_g/\mathrm{d}t=2b_g\alpha_g/(2\pi)$, and b_g is defined by

$$\begin{aligned}b_g=&-\frac{11}{3}C_g(G)+\frac{2}{3}\sum_f T_g(\Psi_f)+\frac{1}{3}T_g(H) \quad (13)\\ &+\left[-\frac{11}{3}C_g(G)+\frac{1}{6}C_g(A_5)+\frac{1}{3}T_g(H)\right]N_{\mathrm{KK}},\end{aligned}$$

where $\alpha_g=g^2/(4\pi)$ and $N_{\mathrm{KK}}=\sum_{i=1}\theta(\mu-m_i)$ which counts the number of activated KK excitations. To understand the RGEs of the gauge couplings, it is noticeable that the KK excitation states of the vector fields A_μ and scalar fields A_5 are the adjoint representations of gauge groups, while the complex weak doublet $H^{(n)}$ particles are the fundamental representations of $SU(2)$.

It is remarkable that, in the extra dimension scenarios we consider, KK excitations always drive the 4D gauge couplings of the $SU(3)$ and $SU(2)$ groups to their weak coupling limits. This is just the typical feature of non-Abelian gauge theory, i.e. asymptotical freedom. Even for the case that all fermions in the SM (the universal case) are assumed to live in the bulk, the couplings of $SU(3) \times SU(2)$ are driven to their weak coupling limit. To simplify our analysis, below we will omit the contributions of the $U(1)$ group, due to its small effect on the problem we consider here.

The Yukawa coupling terms of quarks in the SM are assumed to be contact terms and have the form

$$L_{QUH} = Y_t \bar{Q}_{\mathrm{L}} H U_{\mathrm{R}} + \mathrm{h.c.}; \tag{14}$$

here $\bar{Q}_{\mathrm{L}} = (\bar{u}_{\mathrm{L}}, \bar{d}_{\mathrm{L}})$, $H^{\mathrm{T}} = (H^0, H^-)$, and Y_F is for the Yukawa couplings. Most of the Yukawa couplings have only negligible effects on the problem we consider here; it is reasonable to neglect them. Nevertheless, due to its large effect, the contribution of the Yukawa coupling of the top quarks will be considered. Up to one-loop level, its RGE can be written as $\mathrm{d}\alpha_h/\mathrm{d}t = b_h\alpha_h/(2\pi)$, and b_h is defined by

$$\begin{aligned} b_h = & -6C_c(1+2N_{\mathrm{KK}})\alpha_3 - 3C_w(1+2N_{\mathrm{KK}})\alpha_2 \\ & + \left[N_c + \frac{3}{2}(1+2N_{\mathrm{KK}})\right]\alpha_h, \end{aligned} \tag{15}$$

where $\alpha_h = Y_t^2/(4\pi)$, and N_c is the number of colors. The 2 in front of N_{KK} is due to the different normalization of the zero modes and KK excitations. C_c and C_w are quadratic Casimir operators for the fundamental representations of the $SU(3)$ and $SU(2)$ groups, respectively.

The one-loop RGE of the self-interaction coupling of the Higgs zero mode fields λ can be expressed as $\mathrm{d}\alpha_\lambda/\mathrm{d}t = b_\lambda\alpha_\lambda/(2\pi)$ and b_λ is defined by

$$\begin{aligned} b_\lambda = & \, 3(1+N_{\mathrm{KK}})\alpha_\lambda + \Bigg[(2C_w+4)(1+N_{\mathrm{KK}}) \\ & + \left(4C_w^2 + (2D_B+8)c'\right)\frac{\alpha_2}{\alpha_\lambda}(1+N_{\mathrm{KK}})\Bigg]\alpha_2 \\ & -2N_c\left(2\frac{\alpha_h}{\alpha_\lambda} - 1\right)\alpha_h, \end{aligned} \tag{16}$$

where $D_B = 5$, $\alpha_\lambda = \lambda/(4\pi)$, and $c' = 1 + N_w - 2/N_w + 1/N_w^2$. D_B counts the degrees of freedom of the gauge vector bosons. For $SU(2)$, $N_w = 2$. The first two terms in (16) tend to increase α_l, and the last term tends to decrease it, while KK excitations of bosons always tend to drive α_λ to large values.

The one-loop RGE of the Higgs mass term is expressed in the form below:

$$\begin{aligned} \frac{\mathrm{d}m_H^2}{\mathrm{d}t} = \frac{1}{2\pi}\Big\{ & (\alpha_\lambda - 3C_w\alpha_2 + N_c\alpha_h)m_H^2 \\ & + (\alpha_\lambda - 3C_w\alpha_2)N_{\mathrm{KK}}m_H^2 \\ & + \left(\alpha_\lambda + (D_B-2)C_w\alpha_2\right)M_{\mathrm{KK}}^2\Big\}, \end{aligned} \tag{17}$$

where $M_{\mathrm{KK}}^2 = \sum_{i=1} m_i^2\theta(\mu - m_i)$, which counts the contributions of the KK excitations to m_H^2.

A comment is in order on (16) and (17). One might think that non-standard divergences (cubic and quintic ones) might appear in the effective potential; then (16) and (17) might have an ambiguity. To calculate the one-loop effective potential of KK excitations, we note that the Higgs field in the L^{eff} is a weak doublet complex scalar field. The L^{eff} preserves the symmetries of the SM, the KK excitations of the vector bosons are adjoint representations of the gauge group and those of the scalars are fundamental representations. So when the computation is conducted in the dimensional regularization which preserves the gauge symmetry of the L^{eff}, after taking into account the contributions of the KK excitations to the Higgs potential, no cubic or quintic counter-term will appear in the Higgs potential. In more detail, considering our problem with the path integral method, the contribution of scalars and vector bosons to the effective potential will not contribute the non-standard divergences. The contribution of quarks is the possible origin for the appearance of the non-standard divergence. Then let us study the quark's contribution more carefully. After integrating out the weak left-hand doublet $(t, b)_{\mathrm{L}}$ and the weak singlet b_{R} and t_{R}, we will get the determinants of these fields, which are dependent upon the background field of the Higgs field. To extract the effective potential of the Higgs field from these determinants, we can use perturbation theory to do the calculations. Then again, due to the symmetry of the SM, the diagrams with only one external leg of the Higgs field, those with three legs, with five legs, and so on, vanish, at least in one-loop level and in the dimensional regularization scheme. Therefore, in the MEPM we can get (16) and (17) without any ambiguity, even in the universal case where the fermion fields have their KK excitations.

At the low energy scale regions, only two free parameters needed to be specified:
(1) α_l and
(2) M_c for the 4D effective Lagrangian, while α_3, α_2, α_h, and m_H^2 can be determined from experiment and from these two free parameters. The value of $\alpha_\lambda(M_Z)$ is constrained by the triviality and stability conditions of the Higgs potential in the SM [15], while $M_c = 1/R_c$ is to determine where KK excitations should be counted and should be larger than M_Z. According to the present estimated constraints from the literature [1], it should be near a few TeV.

We will focus on the analysis of the behavior of m_H^2 from the infrared boundary M_Z up to the ultraviolet boundary Λ_{UV}. The initial value of the m_H^2 is assumed to be negative as required by the SM, and its initial value at M_Z can be fixed by v (the VEV of H) and the free parameter α_λ. For the energy scale region below the threshold of the first KK excitations, only the first term in (17) contributes and tends to further decrease the negative m_H^2.

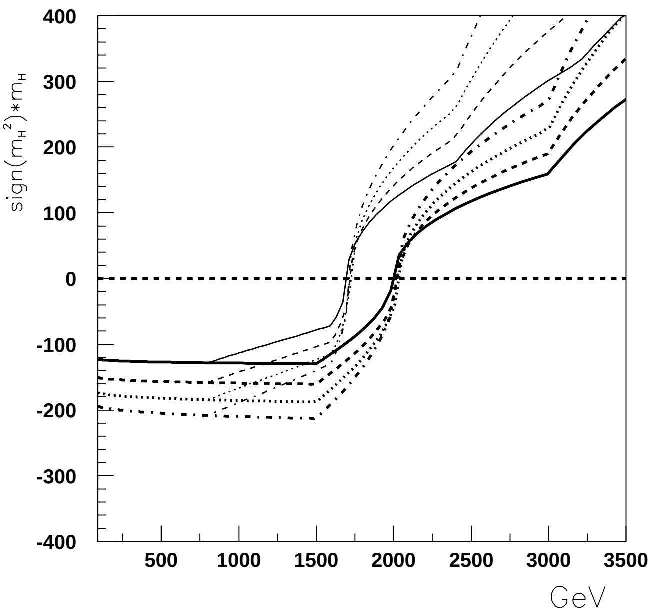

Fig. 1. The varying of the sign of m_H^2 and the value of m_H with the energy scale. The solid, dashed dot, and dash-dot lines represent $\lambda_{(M_Z)} = 1.0$, $\lambda_{(M_Z)} = 1.5$, $\lambda_{(M_Z)} = 2.0$ and $\lambda_{(M_Z)} = 2.5$, respectively. The group of wide (thin) lines corresponds to the case $M_c = 0.8\,\text{TeV}$ ($M_c = 1.5\,\text{TeV}$)

However, once the energy scale runs across the threshold of the first KK excitations and the degree of freedom of KK excitations are activated, m_H^2 can quickly change its sign from negative to positive, as shown in Fig. 1.

As we know, in the SM, the sign of the mass term m_H^2 completely determines whether the electroweak symmetry is broken or unbroken. The change of m_H^2 from the negative to the positive value hints to restoring of the broken electroweak symmetry. With an appropriate value of $\alpha_\lambda(M_Z)$, the process of the symmetry restoring can quickly occur once the running scale runs across the threshold of the first KK excitations, as we can see from (17). Therefore, in the up-to-down viewpoint, we can conclude that it is possible to have the radiative SEWSB mechanism in the ED scenarios of the SM, as in the case of the softly broken SUSY models. The underlying reason for this radiative SEWSB mechanism in ED to work is that the KK excitations of the bosons in the theory are quite massive.

We have also checked the case in which only vector bosons live in the bulk, and find that with properly chosen M_c and λ, the radiative breaking mechanism exists too. For the universal case, where all particles of the SM can propagate in the bulk, we find that due to the large contribution from KK excitations of the top quark and the large top Yukawa coupling, it is relatively hard to find an appropriate point which can both invoke SEWB and preserve the validity of perturbation theory. However, in the universal case of the two Higgs doublets extension of the SM, with the help of $\tan\beta$ (the ratio of the VEV of the two Higgs doublets), it is still possible to have the radiative SEWSB mechanism induced by massive KK bosons.

Although the Higgs self-energy is quadratically divergent in our case, and there are large unknown corrections given by physics at the cut-off scale, the positive value of the mass square of the Higgs mass at the UV cut-off scale might be attributed to the more fundamental theory, say supersymmetry breaking and decoupling.

In the low energy limit of the superstring theory, and in the case where the SUSY decouples before the KK excitations do (this case is possible for a TeV superstring [16]), KK excitations will be the new physics that interact with the SM sector. Then the REWSB mechanism given in this paper will naturally explain the symmetry breaking of the SM. This mechanism will also be helpful for the model construction of the NSUSY extra dimension model at a few TeV.

In summary, we investigate the radiative electroweak symmetry breaking in the NSUSY extra dimensions scenarios of the SM extension proposed by Antoniadis et al., Dienes et al. and Pomarol et al. Utilizing the decoupling theorem and the one-loop renormalization group equations, we find in the up-to-down viewpoint that those massive KK bosons can change the sign of the Higgs mass terms from positive to negative and therefore trigger the SEWSB. We conclude that the radiative mechanism can naturally exist in the NSUSY ED models, and neither new particle contents beyond the SM from the SUSY nor the TC are necessary for this mechanism to work.

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