Abstract

We investigate how the mass correction appears in 5-D with Scherk-Schwarz compactification and clarify whether the KK regularization is reliable method or not. In the extremely sharp cutoff limit of the 5-D regulator which preserves 5-D Lorentz invariance, we prove that the one loop correction to the mass does not depend on the ultraviolet physics for the Scherk-Schwarz breaking of supersymmetry. This is a unique property of Scherk-Schwarz breaking which is given by the boundary condition.
1 Introduction

Theories with extra dimension have been extensively studied by string theorists and phenomenologists. Extra dimension comes out naturally from string theory, and it looks desirable to see what we can get from extra dimensions. The best situation would be to derive all the physical predictions from an ultimate theory, but we are far away from the final dream. As a second best way, more practical and phenomenological bottom up approach has been done recently. Hierarchy problem, doublet-triplet splitting problem in GUT, and other problems are reconsidered within the framework of the models with extra dimensions. Among these, especially in some models [1, 2], Scherk-Schwarz (SS) breaking [3] plays an important role. Therefore it is very necessary to understand the physical properties of SS mechanism.

Calculations in the compact extra dimension models involve towers of Kaluza-Klein (KK) modes and the correct procedure of treating these fields has not been established. The classical decomposition of higher dimensional fields into their infinitely many discrete modes yields several interesting puzzles when we inquire quantum effects. The well known problem is the incapability of direct application of dimensional regularization (DR) to it. DR is the most popular regularization scheme that is widely used due to its powerful convenience in the calculation. However, the non-decoupling theorem does not hold for DR, and this makes it impossible to apply DR to KK towers directly. Matching and integrating out are essential ingredients in DR, and these concepts prevent us from direct application of DR to extra dimensional theory with KK modes.

The choice of the regularization is a matter of taste as long as the regularization keeps all the symmetries of the theory. In most cases, the regularization dependence appears only at the next loop order and is not important as long as the parameters are related to the measured quantities by the renormalization condition. The situation is different when we have a theory that can predict something. If one physical parameter does not have any tree level bare term for it and only can be induced by loop correction, we can predict it by measuring other parameters involved in the loop correction.

Recently one 5-D orbifold model [4] with SS supersymmetry breaking has been suggested, and in the model Higgs mass is predicted by measuring top quark mass and Z boson mass. One extra dimension is compactified to $S^1/Z_2 \times Z'_2$ and boundary condition is assigned to break supersymmetry such that only the Standard Model particles are remained as zero modes.
Yukawa couplings are allowed in the superpotential at two fixed points, and they generate one loop Higgs mass combined with mass splitting of top and stop induced by SS supersymmetry breaking. (There are other Yukawa corrections and gauge coupling corrections but top-stop gives the dominant contribution.) This triggered an interesting debate on the so-called “Kaluza-Klein regularization” scheme.

In the papers [4, 5, 6] they used KK regularization that sums up all KK modes first and then integrated over 4-D momentum. There is no scale above which supersymmetry is manifest, but the one loop correction to the Higgs mass depends only on the compactification radius. Position-momentum space mixed formalism has been used in [4] to show the exponentially decaying physics of higher 4-D momentum contributions for spatially separated propagation in the extra dimension. In [7], for generic momentum cutoff \( \Lambda \) and the truncation of KK modes at some cutoff \( N/R \), it was shown that quadratic divergences reappear in the calculation. One interesting observation was that quadratic divergences cancel separately in the bosonic and the fermionic sectors for KK regularization, i.e., for infinite sums of KK modes. In [10] it was stressed that the momentum cutoff and the truncation scale should be the same due to the 5-D Lorentz invariance, and KK regularization is an extremely anisotropic regularization that does not keep the 5-D Lorentz invariance. Furthermore, it was anticipated that the one loop mass correction to the Higgs would be highly sensitive to the UV physics if we use the physical isotropic regularization [13].

There appeared many papers [8, 9, 10, 11, 12] which support the finiteness (or UV insensitiveness) of one loop correction to the Higgs mass originally obtained by weird KK regularization. In [9], Gaussian distribution of fixed point brane has been used in order to soften the contributions of heavy KK modes, but the result is obtained only after infinite sums over KK modes. In [11], the calculation has been done in 4-D PV and 4-D DR, and also it needs infinite sums over KK modes which looks not sensible physically. All these calculations are based on their own regularizations which may contain hidden subtraction and can not rule out the possibility of the regularization which

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1In 4-D DR with minimal subtraction(MS) gives a finite answer for scalar mass loop correction and the divergence is not seen clearly. In odd dimensions, it is more puzzling since Gamma function \( \Gamma(\frac{2-D}{2}) \) appearing in the expression does not have a pole for odd \( D \) and the correction looks finite. In this case, the subtraction of \( (cubic/linear(D = 5/D = 3)) \) divergence is hidden in the process of replacing the divergent integral to finite Gamma function. All these divergences are seen clearly in the momentum cutoff regularization.
is UV sensitive. In the papers [11, 12] it was claimed that the quadratically divergent result obtained in [7, 13] is because the sharp momentum cut-off is not preserving the underlying symmetry, ‘local supersymmetry’ [14]. However, it is hard to accept the argument because we know that mass dependent regularization shows the power law divergent properties of physical quantities more clearly. Momentum cutoff itself also preserves supersymmetry and shows an excellent cancellations of bosonic and fermionic contribution in the supersymmetric and softly broken limit. Locality of the extra dimension was also stressed, but it is clear that for sufficient high momentum cutoff the violation of the locality would be extremely tiny and the smearing effect would not affect the result. Momentum cutoff of the theory is the scale below which the locality is valid. Therefore, it is remained as a puzzle and is still on the debate [15, 17, 18, 19].

In this paper we extend the observation given in [13] and obtain more precise result based on the regularization scheme which keeps 5-D Lorentz invariance [12] manifestly. To capture the UV physics above the compactification scale correctly, it is necessary to maintain 5-D Lorentz invariance since we can not distinguish the compactified extra dimension from other noncompact 4 dimension at short distances. In other words short distance physics does not discriminate KK modes from 4-D momentum, and the correct regularization should also have this property. The observation made in [13] is not complete because the 4-D momentum cutoff differs for each KK mode if we start from 5-D momentum cutoff. Let $R$ be the compactification radius, $N$ be the level of the KK modes such that KK modes heavier than $N/R$ are truncated, and $\Lambda$ be the cutoff of the 4-D momentum. The precise relation in Euclidean space is $p_{4}^2 + p_{5}^2 \leq \Lambda^2$. Thus $N = [R\Lambda]$ is fixed with $[ ]$ Gaussian integer and the 4-D momentum of the $n$th KK mode has a cutoff $p_{4n}^2 \leq \Lambda^2 - (n/R)^2$ and $p_{4N}$ should be nearly zero. In this paper we calculate the effect of these mode dependent cutoff motivated by 5-D Lorentz invariance and show that the UV insensitive result is recovered. This is a distinct feature of SS supersymmetry breaking mechanism.

2 Scherk-Schwarz mechanism

Scherk-Schwarz breaking of symmetry is realized by modifying the periodic boundary condition along the compact extra dimension with compactification
radius $R$. For the circle compactification,

$$\phi(x^\mu, y + 2\pi R) = e^{2\pi i a} \phi(x^\mu, y),$$

where $0 \leq a < 1$ (or $-\frac{1}{2} < a \leq \frac{1}{2}$) is a parameter controlling the breaking of the relevant symmetry. For maximal breaking of supersymmetry we give $a = \frac{1}{2}$ for bosons and $a = 0$ for its superpartners and vice versa. The simplest realization of the effect is

$$\phi(x^\mu, y) = \sum_{n=-\infty}^{\infty} e^{i(n+a)\frac{2\pi}{R}} \phi_n(x^\mu),$$

and the generalized realization for orbifolding setup is in [20, 21]. The spectrum is modified as

$$m_{nB,F}^2 = \left( \frac{1}{R} \right)^2 (n + a_{B,F})^2,$$

for bosons and fermions.

Scherk-Schwarz breaking of supersymmetry should be distinguished from orbifolding breaking of it. In order to get chiral theory from higher dimensional theory, orbifolding procedure is essential and $S^1/Z_2$ is assumed in which $Z_2$ is for $y \rightarrow -y$. It is the common setup that can have chiral zero modes. $Z_2$ breaks $N = 2$ to $N = 1$ in 4-D viewpoint, and there is $N = 1$ supersymmetry after the orbifolding and two fixed points. There are two ways of breaking this $N = 1$ further. One is the orbifolding with $Z'_2$ and the other is SS breaking of supersymmetry [1, 2]. Apparently it looks different, but as long as the spectrum is concerned, we can make the same spectrum of orbifolding by taking $a = \frac{1}{2}$ in SS breaking. For gauge theories, the equivalence of breaking by orbifold boundary condition and by background gauge fields is discussed in [19]. Thus the model in [4] is also considered as a kind of SS breaking of supersymmetry.

Though our main concern is the one loop mass correction to the Higgs from top and stop loop involving Yukawa couplings, we consider more general expressions that hold for corrections involving Yukawa couplings and gauge couplings. The expression is given already for its analytic continuation to the Euclidean space. We are interested in the generic properties of SS supersymmetry breaking, and set $a_B = a$ and $a_F = 0$ for simplicity. One loop mass correction to the scalar from bosons and fermion loops is given by

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2The author thanks H. P. Nilles for pointing out this.
\[ m^2 = C g^2 A \]  
\[ A = \sum_{n=-\infty}^{\infty} \int_0^{\infty} dp p^3 \left[ \frac{1}{p^2 + (n+a)^2} - \frac{1}{p^2 + n^2} \right], \]  
with \( C \) a constant which is irrelevant for later discussion and \( g \) is a coupling constant (Yukawa or gauge). We use rescaled dimensionless 4-D momentum \( p = RP \) of 4-D momentum \( P \) and \( \hat{\Lambda} = R\Lambda \) of the 5-D momentum cutoff \( \Lambda \) for notational simplicity. Since each term is divergent, we have to regulate the propagator. The first option is to use the Pauli-Villars regulator. For massless fields, the regulator corresponds to

\[ \frac{1}{p^2} \rightarrow \frac{1}{p^2 \left( \frac{\Lambda^2}{p^2 + \Lambda^2} \right)} . \]  

We can apply it to all of the KK modes.

\[ \frac{1}{p^2 + m^2} \rightarrow \frac{1}{p^2 + m^2} \left( \frac{\Lambda^2}{p^2 + m^2 + \Lambda^2} \right) = \left[ \frac{1}{p^2 + m^2} - \frac{1}{p^2 + m^2 + \Lambda^2} \right] . \]

Sharp momentum cutoff can be understood as a limit of higher PV regularization, \( L \rightarrow \infty \)

\[ \frac{1}{p^2 + m^2} \rightarrow \frac{1}{p^2 + m^2} \left( \frac{\Lambda^2}{p^2 + m^2 + \Lambda^2} \right)^L = \left[ \frac{1}{p^2 + m^2} - \sum_{l=1}^{L} \frac{1}{p^2 + m^2 + \Lambda^2} \right] . \]

Therefore, we get the correct sharp momentum cutoff in the limit \( L \rightarrow \infty \). This is slightly different from the previous one [4, 13]. In the previous calculation, four dimensional momentum cutoff\( (\Lambda_4) \) and KK momentum truncation \( (N) \) have been taken independently [4], and the cutoff of the four momentum and the KK truncation scale have been identified \( (\Lambda_4 R = N) \) [13] using the physical argument that they are constrained by the five dimensional cutoff. Here we have an improved regularization which avoids unphysical situation automatically. (anisotropic regularization [13])

Though \( L \rightarrow \infty \) corresponds to the sharp momentum cutoff limit which preserves 5-D Lorentz invariance, the above regulator is inadequate since it

\[^{3}\text{The regularization invented in this paper is indebted to V. A. Rubakov. The author thanks him for his keen suggestion.}\]
suppresses all the contributions except extremely IR contribution. For example, near the cutoff \( P^2 + m^2 \sim \Lambda \), the propagator is already \( (\frac{1}{2})^L \) suppressed and there remains rare correction.

Therefore we use different momentum dependent regularization that keeps the contribution near the cutoff finite even in the sharp cutoff limit. Let us start from the regulated propagator of 5-D theory

\[
\frac{1}{P^2} \rightarrow \frac{1}{P^2} e^{-\frac{\delta^2}{\Lambda^2}}. \tag{9}
\]

Though it is nonlocal and there is no Hamiltonian formulation for this modified propagator, it can be thought of as a perfectly good regulated propagator. Furthermore, we can generalize it for arbitrary integer \( L > 0 \) such that

\[
\frac{1}{P^2} \rightarrow \frac{1}{P^2} e^{-\frac{\delta^2}{\Lambda^2} L}. \tag{10}
\]

Since here \( \tilde{P}^2 \) corresponds to the 5-D momentum square and is decomposed into \( P^2 + (\frac{n+a}{R})^2 \). The KK mode dependent regulator keeps 5-D Lorentz invariance manifestly. Now near the cutoff, \( p^2 + (n+a)^2 = \bar{\Lambda}^2 \), the propagator is just \( 1/e \) times the original one and is independent of \( L \). Thus we can take \( L \rightarrow \infty \) limit while keeping the near cutoff contributions finite. The smooth function becomes a step function and the sharp momentum cutoff is realized. In the discussion, \( L \rightarrow \infty \) is not necessarily important but just simplifies the proof. Actually all the properties are maintained only if \( L \geq \bar{\Lambda}(= R \Lambda) \) since the step function can be an accurate approximation for \( L \geq \bar{\Lambda} \). The propagator is varied rapidly near \( \bar{\Lambda} \) just within the interval \( \bar{\Lambda} - 1 \) and \( \bar{\Lambda} + 1 \) for \( L \geq \bar{\Lambda} \). Having this keeping in mind, we use \( L \rightarrow \infty \) limit in the following discussion. The sum is truncated at \( N = [\bar{\Lambda}] \) and the 4-D momentum cutoff becomes

\[
C_n = \sqrt{(\bar{\Lambda}^2 - (n+a)^2)}. \tag{11}
\]

The one loop contribution in eq. (4),(5) is

\[
A = A(a) - A(0), \tag{12}
\]
\[ A(a) = \sum_{n=-N}^{N} \int_{0}^{C_n} dp \frac{p^3}{p^2 + (n + a)^2} \] (13)

\[ = \sum_{n=-N}^{N} \left[ \frac{1}{2} C_n^2 - \frac{1}{2} (n + a)^2 \log \left( \frac{C_n^2 + (n + a)^2}{(n + a)^2} \right) \right] \] (14)

\[ = \sum_{n=-N}^{N} \left[ \frac{1}{2} (\bar{\Lambda}^2 - (n + a)^2) - \frac{1}{2} (n + a)^2 \log \left( \frac{\bar{\Lambda}^2}{(n + a)^2} \right) \right]. \] (15)

Direct summation over \( N \) with \( \bar{\Lambda} = N + \frac{1}{2} \) and Stirling’s formula give us a simple expression in which \( N^3, N^2, N \) terms are cancelled.

\[ A(a) = \sum_{n=-N}^{N} \left[ \frac{1}{2} (N + \frac{1}{2})^2 - \frac{1}{2} (n + a)^2 
- \frac{1}{2} (n + a)^2 \log \left( \frac{(N + \frac{1}{2})^2}{n^2(1 + \frac{a}{n})^2} \right) \right] \] (16)

\[ A = A(a) - A(0) = -\frac{5}{12} \zeta(2)a^2 + O\left(\frac{a^2}{N^2}, a^4\right) \]
\[ = -\frac{5 \pi^2}{72} a^2 + O\left(\frac{a^2}{N^2}, a^4\right). \] (17)

From the expression, cubic divergence appears both in the bosonic and fermionic loops and cancel with each other. The final result does not depend on any positive powers of \( N \) (or \( \bar{\Lambda} \), the UV cutoff) and \( \log N \). This is a special property of Scherk-Schwarz supersymmetry breaking. Here the cancellation of quadratic divergence in the bosonic and the fermionic sector respectively is easy to understand from the property of SS supersymmetry breaking. If we expand \( A(a) \) in terms of \( a \), we get

\[ A(a) = \sum_{i=0}^{\infty} A_i a^i = A_0(\mathcal{O}(N^3)) + A_2(\mathcal{O}(N))a^2 + \mathcal{O}(a^4). \]

Kaluza-Klein towers have a \( Z_2 \) symmetry which changes \( a \rightarrow -a \) and

\[ ^4 \text{Theories with } a \text{ and } -a \text{ have the same spectrum.} \]
this prevents terms with odd powers of \( a \) \((A_1, A_3, \cdots)\). Therefore net contribution of boson and fermion is

\[
A = A_2(\mathcal{O}(N))a^2 + \mathcal{O}(a^4).
\]

Apparently there are contributions which are linearly divergent \( \mathcal{O}(N) \) \((\mathcal{O}(\Lambda))\) and logarithmically divergent \( \mathcal{O}(\log N) \) \((\mathcal{O}(\log \Lambda))\). It is puzzling that these two contributions are absent and the final result does not have UV sensitivity.

There are more convenient approach for general proof. When \( N \) is large enough \((\bar{\Lambda}(= R \Lambda) \gg 1)\), we can approximate the sum over finite modes by a definite integral over finite range.

\[
\sum_{n=-N}^{N} f(n+a) = \int_{-N}^{N} dx \ f(x+a) + \frac{1}{2} \left[ f(N+a) + f(-N+a) \right]
\]

\[
+ \sum_{m=1}^{L-1} \frac{1}{(2m)!} B_{2m} \left\{ f^{(2m-1)}(N+a) - f^{(2m-1)}(-N+a) \right\}
\]

\[
+ \frac{1}{(2L)!} B_{2L} \sum_{n=-(N-1)}^{N-1} f^{(2L)}(a+n+\theta),
\]

for some \( 0 < \theta < 1 \) and \( f^{(m)} = \frac{d^m f}{dx^m} \). This is the Euler-Maclaurin formula for \( f(x) \) whose first \( 2L \) derivatives are continuous on the interval \((-N+a, N+a)\).

The general proof for the softness of SS supersymmetry breaking is following. If there are quantities with the following structure

\[
A = \sum_{n=-N}^{N} \left[ f(n+a) - f(n) \right],
\]

then \( A \) does not depend on \( N \) if \( f(n) \) satisfies a few conditions. To be more specific in the 5-D case,

\[
f(x) = \frac{1}{2} \left( \bar{\Lambda}^2 - x^2 \right) - \frac{1}{2} x^2 \log \left( \frac{\bar{\Lambda}^2}{x^2} \right).
\]

\( \text{Note:} \) We assume \(-\frac{1}{2} < a < \frac{1}{2}\) in order to avoid unnecessary subtlety related to the mode located just at the boundary. For \( a = \frac{1}{2}\), we can think it as the limit starting from \( b = a - \epsilon \) with a tiny \( \epsilon > 0 \). This avoids the potential danger of sharp momentum cutoff and makes the sharp momentum cutoff good regulator preserving supersymmetry. Here it is easily maintained just by keeping the numbers of bosons and fermions equal. However this caution does not affect the conclusion because the contribution of these edge KK modes is extremely suppressed due to the constraint on 4-D momentum \( p^2 \leq \Lambda^2 - N^2 \sim 0 \).
Since \( f(x) \) is even function, we have

\[
\begin{align*}
    f(-x) &= f(x), \\
    f^{(2n-1)}(0) &= 0.
\end{align*}
\]

Using the Euler-Maclaurin formula and the special properties of \( f(x) \), we obtain after the Taylor series expansion around \( N \)

\[
A = \sum_{m=1}^{\infty} \left[ \frac{f^{(2m-1)}(N)}{(2m)!} a^{2m} + \frac{2 f^{(2m)}(N)}{(2m)!} a^{2m} + \frac{2 f^{(2m)}(0)}{(2m)!} a^{2m} \right] + \sum_{k=1}^{L-1} \frac{2 B_{2k} f^{(2k-1+2m)}(N)}{(2k)!(2m)!} a^{2m}
\]

where \( B_n \) is the Bernoulli number.

Let us look at the potentially UV dependent terms. Since \( f(x) \) is quadratic in \( x \), to investigate the behavior of \( f^{(1)}(x) \) and \( f^{(2)}(x) \) is enough for our purpose of checking UV dependence. \( f^{(1)}(x) \) can be \( \mathcal{O}(N) \) or \( \mathcal{O}(N \log N) \) and \( f^{(2)}(x) \) can be \( \mathcal{O}(\log N) \). All \( f^{(n)}(x) \) with \( n > 2 \) are \( \mathcal{O}(1) \) or less.) The first term contains \( f^{(1)}(N)a^2 \) term but \( f^{(1)}(N) = \mathcal{O}(1) \) because there is little room for 4-D momentum space of higher KK modes near the cutoff. The second and the third term contain \( f^{(2)}(N)a^2 \) and \( f^{(2)}(0)a^2 \). Therefore, to check \( f^{(2)}(N) \) and \( f^{(2)}(0) \) is enough to insure that there are no UV dependence in the mass correction. It turns out that \( f^{(2)}(N) = \mathcal{O}(1) \) and \( f^{(2)}(N) = \mathcal{O}(1) \) and there is no \( \log N \) (or \( \log \Lambda \)) dependence.

As is shown above, the contribution of KK modes from SS supersymmetry breaking shows extremely soft nature. 5-D Lorentz invariance naturally suppresses potentially dangerous contributions of heavy KK modes (near the cutoff) by reducing the available 4-D momentum space for those modes. It is natural to ask whether this is special in 5-D or generic in all dimensions. In 7-D model, \( f(x) = \mathcal{O}(x^4) \) and \( f^{(1)}(N), f^{(3)}(N), f^{(2)}(N), f^{(2)}(0), f^{(4)}(N) \) and \( f^{(4)}(0) \) are shown to be \( \mathcal{O}(1) \). These are enough to show that there are no UV dependence for the mass correction of the Higgs. Though we have not addressed the general proof, it is very plausible that in all dimensions, the SS supersymmetry breaking is extremely soft.

The properties of SS breaking \([22, 23]\) looks plausible from Wilson line interpretation of the gauge symmetry breaking \([15, 16]\). In that case nontrivial boundary condition can be identified to the configuration with constant background gauge field with periodic boundary condition. The constant gauge
field can be gauged away locally, and it becomes a physical one only for the compact space. UV sensitivity is determined from the local physics, UV insensitive result for SS symmetry breaking looks very natural since we can locally gauge it away.

3 Conclusion

Starting from the regulator preserving 5-D Lorentz invariance, we calculated one loop correction to the scalar mass with SS supersymmetry breaking spectrum in the extra dimensional model. On the contrary to the anticipation from the naive physical isotropic regularization [13] or truncation with momentum cutoff [7], the result is UV insensitive due to the special property of SS supersymmetry breaking. Soft breaking of supersymmetry in the bulk reduces the degree of divergence by two powers with the aid of boson-fermion cancellation, but Scherk-Schwarz breaking gets rid of all the divergences, i.e., UV sensitivities. The result depends only on the compactification radius and the SS supersymmetry breaking parameter. Sharp momentum cutoff is a good regularization consistent with supersymmetry as long as 5-D Lorentz invariance is preserved. The UV insensitive answer obtained by (apparently) unphysical KK regularization is confirmed to be right in the momentum cutoff regularization which is more physical and does not contain any disputable procedure.

While this paper was almost finished, three papers [24, 25, 26] appeared. The shift symmetry in those papers is naturally realized in the regulator used in this paper before sending $L \to \infty$. The shift symmetry is the discrete version of the boosting symmetry for the compact extra dimension, and 5-D Lorentz symmetry contains it.

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