Published by Institute of Physics Publishing for SISSA

Received: *February 11, 2005*
Revised: *April 15, 2005*
Accepted: *May 11, 2005*
Published: *June 16, 2005*

Softness of supersymmetry breaking on the orbifold T^2/Z_2

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Abstract: We consider supersymmetry breaking due to a Scherk-Schwarz twist or localized mass terms in 6d $\mathcal{N} = 1$ supersymmetric gauge theory compactified on the orbifold T^2/Z_2. We show that the Scherk-Schwarz breaking in 6d is equivalent to the localized breaking with mass terms along the lines in extra dimensions. In the presence of the considered supersymmetry breaking, we find that there arises a finite one-loop mass correction to a brane scalar due to the KK modes of bulk gauge fields.

Keywords: Field Theories in Higher Dimensions, Large Extra Dimensions, Supersymmetry Breaking, Beyond Standard Model.

Contents

1. Introduction

Weak-scale supersymmetry(SUSY) [1] has been a promising candidate for physics beyond the Standard Model due to the natural solution to the hierarchy problem and the gauge coupling unification and etc. It is well known that supergravity mediation of SUSY breaking at the hidden sector generates all required soft SUSY breaking terms of order the weak scale [1]. However, it does not explain how soft masses approximately conserve flavor as required by bounds on flavor-changing neutral currents.

Recently, there has been a lot of attention to models with extra dimensions which give a new ground for understanding the SUSY breaking in a geometric way. Identifying extra dimensions by discrete actions leads to orbifolds [2], which lead to chiral fermions and the reduction of higher dimensional supersymmetry. Moreover, all or some of SM particles can be regarded to live on the appearing orbifold fixed points or branes.

Particularly, one can impose on bulk fields twisted boundary conditions in extra dimensions, *á la* Scherk-Schwarz(SS) [3]. Then, one can break further the remaining SUSY after orbifolding. In 5d $\mathcal{N}=1$ SUSY gauge theory compactified on S^1/Z_2, it was shown that in the presence of the SS breaking of SUSY, there arises a finite one-loop mass correction of the zero mode of a bulk scalar or a brane scalar due to the sum of Kaluza-Klein(KK) modes of bulk fields [4]. It turns out that the SS breaking is equivalent to the case with a nonzero auxiliary field (F term) of the radion multiplet in the off-shell 5d supergravity [5 − 7]. A nonzero twist parameter or F term can be determined dynamically after the radion stabilization [8, 9].

On the other hand, one can consider the localized breaking of SUSY at the orbifold fixed points [10 − 13, 7, 14 − 20]. For instance, when the remaining SUSY is broken at the hidden brane, only bulk fields such as gaugino or gravitino get nonzero masses at tree level and the broken SUSY is transmitted to the visible brane by bulk fields. Then, one can find that the mass spectrum of bulk fields and their coupling at the visible brane are equivalent to those in the SS breaking without brane mass terms [7, 14]. Therefore, there also appears the one-loop finiteness for a scalar mass of the visible brane [13, 19, 20], which is due to the geometric separation of SUSY breaking from the visible brane. If the broken SUSY is mediated dominantly by gaugino, the scalar mass becomes flavor-blind which sheds light on the supersymmetric flavor problem [12].

In this paper, we will consider the SUSY breaking in 6d $\mathcal{N}=1$ supersymmetric gauge theory compactified on the orbifold T^2/Z_2 [21]. Even if we consider only U(1) gauge group in the bulk, it is straightforward to extend to the non-abelian gauge group. The orbifold fixed points on T^2/Z_2 correspond to codimension-two branes. First we consider a generalization of the SS breaking of SUSY to the 6d case. Then, we show that the SS breaking is equivalent to the localized breaking with mass terms along the lines rather than points. This localized SUSY breaking can be realized by positioning the hidden sector at the fixed boundaries under additional Z_2 actions.

For the localized breaking with mass terms at the codimension-two brane, however, the classical solution of a bulk field is singular for an infinitely thin brane [22]. So, one must regularize the zero thickness of brane. Then, the regulator dependence in the classical solution is absorbed into the renormalized brane mass, which has a *classical* logarithmic RG running [22]. In that sense, the localized breaking at the codimension-two brane is sensitive to the ultraviolet physics of regularization even in a mild way with the log divergence. Actually, it has been shown [23] that in the presence of mass terms localized at the fixed points, the one-loop mass for a brane scalar due to bulk gauge fields has a log divergence due to the infinitely thin brane.

On the other hand, the localized mass terms at codimension-one branes are insensitive to the regularization of the brane thickness, as seen from the equivalence to the SS breaking. By using the off-shell action for 6d SUSY gauge theory with the bulk-brane coupling [21], we make a computation of one-loop mass correction to a brane scalar due to the SS breaking or the localized breaking along the distant lines. Thus, we find that the resulting one-loop correction is finite. In the limit of taking one extra dimension without a SS twist to be much smaller than the other one with a SS twist, we reproduce the 5d result with a SS twist. On the other hand, a small extra dimension with a nontrivial SS twist is not decoupled but rather its effect is dominant in the one-loop mass correction.

The paper is organized as follows. First we describe the SS twisted boundary conditions on the bulk gaugino and find the mass spectrum and mode expansion of gaugino. Next in the localized SUSY breaking with general Z_2-even mass terms along the lines, we obtain the similar result as in the SS twist. Then, we compute the one-loop mass correction to a brane scalar due to the KK modes of bulk gauge fields. Finally the conclusion is drawn.

2. Scherk-Schwarz breaking of SUSY on T^2/Z_2

Let us consider a 6d $\mathcal{N}=1$ supersymmetric U(1) gauge theory compactified on the T^2/Z_2 orbifold.[1] Two extra dimensions on a torus are identified as $x_5 \equiv x_5 + 2\pi R_5$ and $x_6 \equiv$

[1]It is straightforward to include hypermultiplets coupled to the U(1) [21] and extend to bulk non-abelian

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$x_6 + 2\pi R_6$ where R_5 and R_6 are radii of extra dimensions. By orbifolding on the torus by Z_2, we identify (x_5, x_6) with $(-x_5, -x_6)$. Then, there appear four orbifold fixed points,

$$(0,0)\,, \qquad (\pi R_5, 0)\,, \qquad (0, \pi R_6)\,, \qquad (\pi R_5, \pi R_6)\,. \tag{2.1}$$

The fundamental region is the half of the torus.

The kinetic term for the U(1) gaugino[2] is given by

$$\mathcal{L} = i\bar{\Omega}_i \Gamma^M \partial_M \Omega^i\,. \tag{2.2}$$

The gaugino Ω^i is a right-handed simplectic Majorana-Weyl fermion satisfying the chirality condition

$$\Gamma^7 \Omega^i = \Omega^i\,. \tag{2.3}$$

On writing the gaugino in a four dimensional Weyl representation $\Omega^i_R \equiv \lambda^i$, eq. (2.2) becomes

$$\mathcal{L} = i\bar{\lambda}_i \gamma^M \partial_M \lambda^i\,. \tag{2.4}$$

From the symmetry of the action on an orbifold T^2/Z_2, let us consider the orbifold boundary conditions and the Scherk-Schwarz(SS) twists on T^2/Z_2 as follows,

$$Z_2: \;\; \lambda(x, -x_5, -x_6) = \tau_3(i\gamma^5)\lambda(x, x_5, x_6) \equiv P\lambda(x, x_5, x_6)\,, \tag{2.5}$$

$$T_1: \;\; \lambda(x, x_5 + 2\pi R_5, x_6) = U_1 \lambda(x, x_5, x_6)\,, \tag{2.6}$$

$$T_2: \;\; \lambda(x, x_5, x_6 + 2\pi R_6) = U_2 \lambda(x, x_5, x_6) \tag{2.7}$$

where $U_i (i = 1, 2)$ are 2×2 twist matrices corresponding to $\mathrm{SU}(2)_R$ rotations. The SS twists on the orbifold are subject to the consistency conditions $U_i P U_i = P (i = 1, 2)$ and $U_1 U_2 = U_2 U_1$. We note that there is another possible choice of the parity matrix $P = \pm \mathbf{1}_2 (i\gamma^5)$, instead of $P = \tau_3(i\gamma^5)$. In this case, the consistency conditions lead to $U_i = \pm \mathbf{1}_2$ or $\pm\tau_3$. However, in this paper, let us focus on the case with $P = \tau_3(i\gamma^5)$ for which a continuous twist is possible.

The first condition $U_i P U_i = P (i = 1, 2)$ gives rise to the following form for either U_1 or U_2: a continuous twist connected to the identity,

$$U_i = e^{-i[2\pi\omega_i(\tau_1 \sin\phi_i + \tau_2 \cos\phi_i)]} \tag{2.8}$$

with ω_i, ϕ_i being real parameters, or a discrete twist not connected to the identity,

$$U_i = -\mathbf{1}_2\,. \tag{2.9}$$

By using the residual global invariance, a continuous twist(U_i with $i = 1$ or 2) can be always set to the one with $\phi_i = 0$. Therefore, also considering the second condition $U_1 U_2 = U_2 U_1$, we find that there are four possible twisted boundary conditions:

$$U_1 = e^{-2\pi i \omega_5 \tau_2}\,, \qquad U_2 = e^{2\pi i \omega_6 \tau_2}\,, \tag{2.10}$$

gauge groups. In these cases, one needs to remember that the bulk matter content is severely restricted due to genuine 6d anomalies [21].

[2] For notations and conventions, refer to [21].

$$U_1 = -\mathbf{1}_2\,, \qquad U_2 = e^{2\pi i\omega_6\tau_2}\,, \tag{2.11}$$
$$U_1 = e^{-2\pi i\omega_5\tau_2}\,, \qquad U_2 = -\mathbf{1}_2\,, \tag{2.12}$$
$$U_1 = U_2 = -\mathbf{1}_2 \tag{2.13}$$

where ω_5, ω_6 are real constant parameters. We note that the discrete choice of twist matrices corresponds to using R-parity of $\mathcal{N} = 1$ 4d supersymmetry as the global symmetry.

First, for the case with continuous twists in both extra dimensions given by eq. (2.10), let us make a redefinition of the gaugino as

$$\lambda(x, x_5, x_6) = e^{-i(\omega_5 x_5/R_5 - \omega_6 x_6/R_6)\tau_2}\tilde{\lambda}(x, x_5, x_6)\,. \tag{2.14}$$

Then, regarding $\tilde{\lambda}$ to be untwisted fields, we take the redefined gaugino to be a solution to the twisted boundary conditions (2.6) and (2.7) with eq. (2.10). Moreover, one can show that $\tilde{\lambda}$ satisfies the same orbifold boundary condition as λ in eq. (2.5).

Let us write the untwisted fields $\tilde{\lambda}$ in terms of 4d Majorana spinors $\tilde{\psi}^i (i = 1, 2)$ as

$$\tilde{\lambda}^1 = +\frac{1}{2}(1 + i\gamma^5)\tilde{\psi}^1 + \frac{1}{2}(1 - i\gamma^5)\tilde{\psi}^2\,, \tag{2.15}$$
$$\tilde{\lambda}^2 = -\frac{1}{2}(1 - i\gamma^5)\tilde{\psi}^1 + \frac{1}{2}(1 + i\gamma^5)\tilde{\psi}^2\,, \tag{2.16}$$

and similarly for the twisted fields λ in terms of 4d Majorana spinors $\psi^i (i = 1, 2)$. We note that the untwisted Majorana spinors are related to the twisted ones by

$$\begin{pmatrix} \psi^1 \\ \psi^2 \end{pmatrix} = \begin{pmatrix} \cos\left(\frac{\omega_5 x_5}{R_5} + \frac{\omega_6 x_6}{R_6}\right) & -\sin\left(\frac{\omega_5 x_5}{R_5} + \frac{\omega_6 x_6}{R_6}\right) \\ \sin\left(\frac{\omega_5 x_5}{R_5} + \frac{\omega_6 x_6}{R_6}\right) & \cos\left(\frac{\omega_5 x_5}{R_5} + \frac{\omega_6 x_6}{R_6}\right) \end{pmatrix} \begin{pmatrix} \tilde{\psi}^1 \\ \tilde{\psi}^2 \end{pmatrix}. \tag{2.17}$$

Then, from eq. (2.5), one can show that $\tilde{\psi}^i$ satisfy the following Z_2 boundary conditions,

$$\tilde{\psi}^1(x, -x_5, -x_6) = +\tilde{\psi}^1(x, x_5, x_6)\,, \tag{2.18}$$
$$\tilde{\psi}^2(x, -x_5, -x_6) = -\tilde{\psi}^2(x, x_5, x_6)\,, \tag{2.19}$$

and similarly for ψ^i. With this redefinition of fields, let us write the gaugino kinetic term (2.4) in terms of untwisted fields $\tilde{\psi}$ as

$$\begin{aligned}\mathcal{L} = {} & i\overline{\tilde{\psi}^1}\gamma^\mu\partial_\mu\tilde{\psi}^1 + i\overline{\tilde{\psi}^2}\gamma^\mu\partial_\mu\tilde{\psi}^2 - \overline{\tilde{\psi}^1}(\partial_5 + \gamma^5\partial_6)\tilde{\psi}^2 + \overline{\tilde{\psi}^2}(\partial_5 + \gamma^5\partial_6)\tilde{\psi}^1 - \\ & -\frac{\omega_5}{R_5}(\overline{\tilde{\psi}^1}\tilde{\psi}^1 + \overline{\tilde{\psi}^2}\tilde{\psi}^2) + \frac{\omega_6}{R_6}(\overline{\tilde{\psi}^1}\gamma^5\tilde{\psi}^1 + \overline{\tilde{\psi}^2}\gamma^5\tilde{\psi}^2).\end{aligned} \tag{2.20}$$

Equivalently, by writing $\tilde{\psi}^i = (\chi^i, \bar{\chi}^i)^T (i = 1, 2)$ with 4d Weyl spinors χ^i, the action becomes

$$\begin{aligned}\mathcal{L} = {} & \sum_{i=1,2}(i\chi^i\sigma^\mu\partial_\mu\bar{\chi}^i + i\bar{\chi}^i\bar{\sigma}^\mu\partial_\mu\chi^i) + \\ & +[-\chi^1(\partial_5 - i\partial_6)\chi^2 + \chi^2(\partial_5 - i\partial_6)\chi^1 + \text{c.c.}] + \mathcal{L}_m\end{aligned} \tag{2.21}$$

where $\mathcal{L}_m$ corresponds to the bulk mass terms given by

$$\mathcal{L}_m = -\left[\left(\frac{\omega_5}{R_5} + i\frac{\omega_6}{R_6}\right)(\chi^1\chi^1 + \chi^2\chi^2) + \text{c.c.}\right]. \tag{2.22}$$

Therefore, we find that the SS twist on the torus induces nonzero Z_2-even mass terms in the basis of the untwisted gaugino that we have introduced for redefining the gaugino. We note that the Z_2-even and odd untwisted fields take equal bulk masses as in the 5d case.

From the action (2.21), we can derive the equations of motion for gaugino as follows,

$$i\sigma^\mu\partial_\mu\bar{\chi}^2 + (\partial_5 - i\partial_6)\chi^1 - \left(\frac{\omega_5}{R_5} + i\frac{\omega_6}{R_6}\right)\chi^2 = 0\,, \tag{2.23}$$

$$i\bar{\sigma}^\mu\partial_\mu\chi^1 - (\partial_5 + i\partial_6)\bar{\chi}^2 - \left(\frac{\omega_5}{R_5} - i\frac{\omega_6}{R_6}\right)\bar{\chi}^1 = 0\,. \tag{2.24}$$

Therefore, solving the above equations, we find the solution for the untwisted gaugino as

$$\begin{pmatrix}\chi^1\\ \chi^2\end{pmatrix}(x,x_5,x_6) = \frac{1}{2\pi\sqrt{R_5R_6}}\sum_{n_5,n_6\in\mathbf{Z}}\begin{pmatrix}\cos\left(\frac{n_5}{R_5}x_5 - \frac{n_6}{R_6}x_6\right)\\ \sin\left(\frac{n_5}{R_5}x_5 - \frac{n_6}{R_6}x_6\right)\end{pmatrix}\eta^{(n_5,n_6)}(x) \tag{2.25}$$

where n_5, n_6 are integer, $i\sigma^\mu\partial_\mu\bar{\eta}^{(n_5,n_6)}(x) = M_{n_5,n_6}\eta^{(n_5,n_6)}(x)$ and the mass spectrum is given by

$$M_{n_5,n_6} = \frac{n_5+\omega_5}{R_5} + i\left(\frac{n_6+\omega_6}{R_6}\right). \tag{2.26}$$

Consequently, due to the relation (2.17), the solution for the twisted gaugino $\psi^i = (\zeta^i, \bar{\zeta}^i)^T$ $(i = 1, 2)$ becomes

$$\begin{pmatrix}\zeta^1\\ \zeta^2\end{pmatrix}(x,x_5,x_6) = \frac{1}{2\pi\sqrt{R_5R_6}}\sum_{n_5,n_6\in\mathbf{Z}}\begin{pmatrix}\cos\left(\frac{n_5+\omega_5}{R_5}x_5 - \frac{n_6+\omega_6}{R_6}x_6\right)\\ \sin\left(\frac{n_5+\omega_5}{R_5}x_5 - \frac{n_6+\omega_6}{R_6}x_6\right)\end{pmatrix}\eta^{(n_5,n_6)}(x)\,. \tag{2.27}$$

Similarly, for the case with a continuous twist in one direction and a discrete twist in the other direction given by eq. (2.11), we can make a redefinition of the gaugino with $\tilde{\lambda}$ as

$$\lambda(x,x_5,x_6) = e^{i(\omega_6x_6/R_6)\tau_2}\tilde{\lambda}(x,x_5,x_6)\,. \tag{2.28}$$

Then, $\tilde{\lambda}$ satisfies the following orbifold and twisted boundary conditions:

$$Z_2:\ \ \tilde{\lambda}(x,-x_5,-x_6) = \tau_3(i\gamma^5)\tilde{\lambda}(x,x_5,x_6) \tag{2.29}$$

$$T_1:\ \ \tilde{\lambda}(x,x_5+2\pi R_5,x_6) = -\tilde{\lambda}(x,x_5,x_6)\,, \tag{2.30}$$

$$T_2:\ \ \tilde{\lambda}(x,x_5,x_6+2\pi R_6) = \tilde{\lambda}(x,x_5,x_6)\,. \tag{2.31}$$

Consequently, plugging the redefined gaugino into the action, deriving the equation for $\tilde{\lambda}$ and imposing the above boundary conditions to $\tilde{\lambda}$, we find the corresponding solution for $\tilde{\lambda}$ in 4d Weyl representation as

$$\begin{pmatrix} \chi^1 \\ \chi^2 \end{pmatrix}(x, x_5, x_6) = \frac{1}{2\pi\sqrt{R_5 R_6}} \sum_{n_5, n_6 \in \mathbf{Z}} \begin{pmatrix} \cos\left(\frac{n_5+\frac{1}{2}}{R_5}x_5 - \frac{n_6}{R_6}x_6\right) \\ \sin\left(\frac{n_5+\frac{1}{2}}{R_5}x_5 - \frac{n_6}{R_6}x_6\right) \end{pmatrix} \eta^{(n_5,n_6)}(x) \tag{2.32}$$

where n_5, n_6 are integer, $i\sigma^\mu \partial_\mu \bar{\eta}^{(n_5,n_6)}(x) = M_{n_5,n_6}\eta^{(n_5,n_6)}(x)$ and the mass spectrum is given by

$$M_{n_5,n_6} = \frac{n_5 + \frac{1}{2}}{R_5} + i\left(\frac{n_6 + \omega_6}{R_6}\right). \tag{2.33}$$

Therefore, the solution for the twisted gaugino λ is given by eq. (2.27) with $\omega_5 = 1/2$. Also for the case with twist matrices (2.12), we only have to interchange $(n_5, R_5) \leftrightarrow (n_6, R_6)$ with $\omega_6 \to \omega_5$ in eq. (2.32), and then obtain the solution for the twisted gaugino λ given by eq. (2.27) with $\omega_6 = 1/2$.

Lastly, for the case with discrete twists in both extra dimensions, the solution for the twisted gaugino λ is given by eq. (2.27) with $\omega_5 = \omega_6 = 1/2$.

3. SUSY breaking due to localized gaugino masses

In this section, instead of the SS boundary twists of gaugino, let us consider a local breaking of supersymmetry which is parametrized by gaugino mass terms, and show the equivalence between the SS breaking and the localized breaking.

Let us take the most general Z_2-even mass terms[3] for gaugino, which are localized along the two lines intersecting at a fixed point $(\pi R_5, \pi R_6)$ on the orbifold,

$$\begin{aligned} \mathcal{L}_m = & -[2m(\chi^1\chi^1 + \rho\chi^2\chi^2) + \text{c.c}]\delta(x_5 - \pi R_5) - \\ & -[2im'(\chi^1\chi^1 + \rho'\chi^2\chi^2) + \text{c.c.}]\delta(x_6 - \pi R_6) \end{aligned} \tag{3.1}$$

where (m, ρ) and (m', ρ') are gaugino mass parameters and they are assumed to be real. Then, the lines with localized mass terms should be regarded as the fixed boundaries under two additional independent Z_2 actions [24]: Z_2': $(x_5, x_6) \to (-x_5, x_6)$ and Z_2'': $(x_5, x_6) \to (x_5, -x_6)$. In this case, it is conceivable that the localized mass terms are due to the SUSY breaking in the hidden sector located on the lines, rather than points.

In this case, the gaugino equations of motion are

$$i\sigma^\mu\partial_\mu\bar{\chi}^2 + (\partial_5 - i\partial_6)\chi^1 - 2(m\rho\delta(x_5 - \pi R_5) + im'\rho'\delta(x_6 - \pi R_6))\chi^2 = 0\,, \tag{3.2}$$

$$i\bar{\sigma}^\mu\partial_\mu\chi^1 - (\partial_5 + i\partial_6)\bar{\chi}^2 - 2(m\delta(x_5 - \pi R_5) - im'\delta(x_6 - \pi R_6))\bar{\chi}^1 = 0\,. \tag{3.3}$$

[3]If one introduces gaugino mass terms proportional to $\delta(x_5)$ and $\delta(x_6)$, there appears a non-supersymmetric gauge coupling at the origin due to the suppression of gaugino wave function [20]. Since we assume the visible sector fields to be localized at the origin, let us consider the gaugino mass terms only at distant lines.

Now let us take the solution of gaugino to the above equations as

$$\begin{pmatrix} \chi^1 \\ \chi^2 \end{pmatrix} (x, x_5, x_6) = \sum_M N_M \begin{pmatrix} u^1(x_5, x_6) \\ u^2(x_5, x_6) \end{pmatrix} \eta_M(x) \tag{3.4}$$

where N_M is the normalization constant and $i\sigma^\mu \partial_\mu \bar{\eta}_M(x) = M\eta_M(x)$. Then, the gaugino equations are

$$M\bar{u}^2 + (\partial_5 - i\partial_6)u^1 - 2(m\rho\delta(x_5 - \pi R_5) + im'\rho'\delta(x_6 - \pi R_6))u^2 = 0\,, \tag{3.5}$$

$$\bar{M}u^1 - (\partial_5 + i\partial_6)\bar{u}^2 - 2(m\delta(x_5 - \pi R_5) - im'\delta(x_6 - \pi R_6))\bar{u}^1 = 0\,. \tag{3.6}$$

Let us take u^1, u^2 to be real functions. Then, taking $M = M_5 + iM_6$ with real M_5 and M_6 and using eqs. (3.5) and (3.6), we obtain the equation for $t \equiv u^2/u^1$ as

$$\partial_5 t = M_5(1 + t^2) - 2m(1 + \rho t^2)\delta(x_5 - \pi R_5)\,, \tag{3.7}$$

$$\partial_6 t = -M_6(1 + t^2) + 2m'(1 + \rho' t^2)\delta(x_6 - \pi R_6)\,. \tag{3.8}$$

It is convenient to consider the Z_2-odd solution of t separately around different fixed points and match them in the overlap regions [20]. That is, let us consider the solution of t which satisfies the equations of motion inside a torus centered at each fixed point. Thus, we find the solution for t:

- $-\pi R_5 < x_5 < \pi R_5$ and $-\pi R_6 < x_6 < \pi R_6$,

$$t = \tan(M_5 x_5 - M_6 x_6)\,. \tag{3.9}$$

- $0 < x_5 < 2\pi R_5$ and $-\pi R_6 < x_6 < \pi R_6$,

$$t = \tan[M_5(x_5 - \pi R_5) - M_6 x_6 - \mathrm{arctan}\alpha(\rho, m\epsilon(x_5 - \pi R_5))] \tag{3.10}$$

  where $\epsilon(x_5 - \pi R_5)$ is a step function with $2\pi R_5$ periodicity given by

$$\epsilon(x_5) = \begin{cases} +1\,, & 0 < x_5 < \pi R_5\,, \\ 0\,, & x_5 = 0\,, \\ -1\,, & -\pi R_5 < x_5 < 0\,, \end{cases} \tag{3.11}$$

  and

$$\alpha(\rho, m\epsilon(x_5 - \pi R_5)) \equiv \frac{1}{\sqrt{\rho}} \tan(\sqrt{\rho} m\epsilon(x_5 - \pi R_5))\,. \tag{3.12}$$

- $-\pi R_5 < x_5 < \pi R_5$ and $0 < x_6 < 2\pi R_6$,

$$t = \tan[M_5 x_5 - M_6(x_6 - \pi R_6) + \mathrm{arctan}\alpha(\rho', m'\tilde{\epsilon}(x_6 - \pi R_6))]\,, \tag{3.13}$$

  where $\tilde{\epsilon}(x_6 - \pi R_6)$ is a step function with $2\pi R_6$ periodicity.

- $0 < x_5 < 2\pi R_5$ and $0 < x_6 < 2\pi R_6$,

$$\begin{aligned} t = \tan[&M_5(x_5 - \pi R_5) - M_6(x_6 - \pi R_6) - \mathrm{arctan}\alpha(\rho, m\epsilon(x_5 - \pi R_5)) + \\ &+\mathrm{arctan}\alpha(\rho', m'\tilde{\epsilon}(x_6 - \pi R_6))]\,. \end{aligned} \tag{3.14}$$

Identify the first two solutions in the overlap region of $0 < x_5 < \pi R_5$ and $-\pi R_6 < x_6 < \pi R_6$, we find

$$M_5 = \frac{1}{R_5}\left(n_5 + \frac{1}{\pi}\mathrm{arctan}\alpha(\rho, m)\right), \qquad n_5 = \text{integer}\,. \tag{3.15}$$

Likewise, identifying the first and third solutions in the overlap region of $-\pi R_5 < x_5 < \pi R_5$ and $0 < x_6 < \pi R_6$, we also find

$$M_6 = \frac{1}{R_6}\left(n_6 + \frac{1}{\pi}\mathrm{arctan}\alpha(\rho', m')\right), \qquad n_6 = \text{integer}\,. \tag{3.16}$$

Then, comparing the other solutions in the overlap regions does not lead to a new condition. Therefore, the mass spectrum is equivalent to the one with SS breaking when ω_5 and ω_6 in eq. (2.26) are identified with $\mathrm{arctan}\alpha(\rho, m)/\pi$ and $\mathrm{arctan}\alpha(\rho', m')/\pi$, respectively.

Moreover, the solutions of u^1 and u^2 are also given in the separate regions:

- $-\pi R_5 < x_5 < \pi R_5$ and $-\pi R_6 < x_6 < \pi R_6$,

$$u^1 = \cos(M_5 x_5 - M_6 x_6)\,, \tag{3.17}$$
$$u^2 = \sin(M_5 x_5 - M_6 x_6)\,. \tag{3.18}$$

- $0 < x_5 < 2\pi R_5$ and $-\pi R_6 < x_6 < \pi R_6$,

$$\begin{aligned} u^1 &= (-1)^{n_5} A(\rho, m\epsilon(x_5 - \pi R_5)) \times \\ &\quad \times \cos[M_5(x_5 - \pi R_5) - M_6 x_6 - \mathrm{arctan}\alpha(\rho, m\epsilon(x_5 - \pi R_5))]\,, \end{aligned} \tag{3.19}$$
$$\begin{aligned} u^2 &= (-1)^{n_5} A(\rho, m\epsilon(x_5 - \pi R_5)) \times \\ &\quad \times \sin[M_5(x_5 - \pi R_5) - M_6 x_6 - \mathrm{arctan}\alpha(\rho, m\epsilon(x_5 - \pi R_5))] \end{aligned} \tag{3.20}$$

where

$$A(\rho, m\epsilon(x_5 - \pi R_5)) \equiv \left(\frac{1 + \alpha^2(\rho, m\epsilon(x_5 - \pi R_5))}{1 + \rho\alpha^2(\rho, m\epsilon(x_5 - \pi R_5))}\right)^{1/2}. \tag{3.21}$$

- $-\pi R_5 < x_5 < \pi R_5$ and $0 < x_6 < 2\pi R_6$,

$$\begin{aligned} u^1 &= (-1)^{n_6} A(\rho', m'\tilde{\epsilon}(x_6 - \pi R_6)) \times \\ &\quad \times \cos[M_5 x_5 - M_6(x_6 - \pi R_6) + \mathrm{arctan}\alpha(\rho', m'\tilde{\epsilon}(x_6 - \pi R_6))]\,, \end{aligned} \tag{3.22}$$
$$\begin{aligned} u^2 &= (-1)^{n_6} A(\rho', m'\tilde{\epsilon}(x_6 - \pi R_6)) \times \\ &\quad \times \sin[M_5 x_5 - M_6(x_6 - \pi R_6) + \mathrm{arctan}\alpha(\rho', m'\tilde{\epsilon}(x_6 - \pi R_6))]\,. \end{aligned} \tag{3.23}$$

- $0 < x_5 < 2\pi R_5$ and $0 < x_6 < 2\pi R_6$,

$$\begin{aligned} u^1 &= (-1)^{n_5+n_6} A(\rho, m\epsilon(x_5 - \pi R_5)) A(\rho', m'\tilde{\epsilon}(x_6 - \pi R_6)) \times \\ &\quad \times \cos[M_5(x_5 - \pi R_5) - M_6(x_6 - \pi R_6) - \\ &\qquad - \mathrm{arctan}\alpha(\rho, m\epsilon(x_5 - \pi R_5)) + \mathrm{arctan}\alpha(\rho', m'\tilde{\epsilon}(x_6 - \pi R_6))]\,, \end{aligned} \tag{3.24}$$
$$\begin{aligned} u^2 &= (-1)^{n_5+n_6} A(\rho, m\epsilon(x_5 - \pi R_5)) A(\rho', m'\tilde{\epsilon}(x_6 - \pi R_6)) \times \\ &\quad \times \sin[M_5(x_5 - \pi R_5) - M_6(x_6 - \pi R_6) - \\ &\qquad - \mathrm{arctan}\alpha(\rho, m\epsilon(x_5 - \pi R_5)) + \mathrm{arctan}\alpha(\rho', m'\tilde{\epsilon}(x_6 - \pi R_6))]\,. \end{aligned} \tag{3.25}$$

In order to make a normalization of KK modes, let us insert the solutions in the action and integrate it over extra dimensions. Then, we obtain the normalization constant in the separate regions:

- $-\pi R_5 < x_5 < \pi R_5$ and $-\pi R_6 < x_6 < \pi R_6$,

$$N_M = \left(\int_{-\pi R_5}^{\pi R_5} dx_5 \int_{-\pi R_6}^{\pi R_6} dx_6 [(u^1)^2 + (u^2)^2] \right)^{-1/2} = \frac{1}{2\pi\sqrt{R_5 R_6}} . \tag{3.26}$$

- $0 < x_5 < 2\pi R_5$ and $-\pi R_6 < x_6 < \pi R_6$,

$$N_M = \left(\int_0^{2\pi R_5} dx_5 \int_{-\pi R_6}^{\pi R_6} dx_6 [(u^1)^2 + (u^2)^2] \right)^{-1/2} = \frac{1}{2\pi\sqrt{R_5 R_6}} A^{-1}(\rho, m) . \tag{3.27}$$

- $-\pi R_5 < x_5 < \pi R_5$ and $0 < x_6 < 2\pi R_6$,

$$N_M = \left(\int_{-\pi R_5}^{\pi R_5} dx_5 \int_0^{2\pi R_6} dx_6 [(u^1)^2 + (u^2)^2] \right)^{-1/2} = \frac{1}{2\pi\sqrt{R_5 R_6}} A^{-1}(\rho', m') . \tag{3.28}$$

- $0 < x_5 < 2\pi R_5$ and $0 < x_6 < 2\pi R_6$,

$$\begin{aligned} N_M &= \left(\int_0^{2\pi R_5} dx_5 \int_0^{2\pi R_6} dx_6 [(u^1)^2 + (u^2)^2] \right)^{-1/2} \\ &= \frac{1}{2\pi\sqrt{R_5 R_6}} A^{-1}(\rho, m) A^{-1}(\rho', m') . \end{aligned} \tag{3.29}$$

4. One-loop mass correction to a brane scalar

In the general case with nonzero gaugino masses, let us put a chiral multiplet at the $(0, 0)$ fixed point. Then, the scalar partner of the chiral multiplet does not feel the supersymmetry breaking directly but there exists a loop contribution to its mass due to the distant supersymmetry breaking. Only Z_2-even gaugino couples to the brane scalar. From the solution (3.17) with normalization (3.26) in the region $-\pi R_5 < 0 < \pi R_5$ and $-\pi R_6 < x_6 < \pi R_6$, we find that all KK modes of Z_2-even gaugino have the same brane coupling as the one of bulk gauge boson,

$$g_4 = \frac{g_6}{2\pi\sqrt{R_5 R_6}} \tag{4.1}$$

where g_6 is the six-dimensional gauge coupling. One has the KK mass spectrums for gauge bosons and gaugino running in loops, respectively,

$$M^2_{(0)n_5,n_6} = \left(\frac{n_5}{R_5}\right)^2 + \left(\frac{n_6}{R_6}\right)^2 , \tag{4.2}$$

$$M^2_{n_5,n_6} = \left(\frac{n_5 + \omega_5}{R_5}\right)^2 + \left(\frac{n_6 + \omega_6}{R_6}\right)^2 . \tag{4.3}$$

On the other hand, from the even-mode ζ^1 from eq. (2.27) at the $(0,0)$ fixed point and the mass spectrum in eq. (2.26), one can find that the SS breaking leads to the same brane coupling and mass spectrum of gaugino as in the localized breaking of supersymmetry. So, the brane scalar fields do not feel the difference between the SS twist and the localized gaugino masses along the distant lines.

Now let us consider the KK mode contribution to the one-loop mass correction for a brane scalar ϕ with charge Q under the U(1). For this, we note that the coupling of the bulk auxiliary field to the brane scalar is given by the following action [21],

$$\int d^6x\left[\frac{1}{2}(D^3)^2+\delta(x^5)\delta(x^6)g_6Q\phi^\dagger(-D^3+F_{56})\phi\right] \tag{4.4}$$

where D^3 is the third component of auxiliary field in the bulk vector multiplet and F_{56} is the extra component of field strength. After eliminating the auxiliary field by its equation of motion, we find the resulting coupling as

$$\int d^4x\left[-g_6Q\phi^\dagger F_{56}(x,x_5=0,x_6=0)\phi-\frac{1}{2}g_6^2Q^2(\phi^\dagger\phi)^2\delta(0)\delta(0)\right] \tag{4.5}$$

with

$$\begin{aligned}\delta(0)\delta(0) &= \frac{1}{4\pi^2R_5R_6}\sum_{n_5,n_6\in\mathbf{Z}}1\\ &= \frac{1}{4\pi^2R_5R_6}\sum_{n_5,n_6\in\mathbf{Z}}\frac{p^2-M^2_{(0)n_5,n_6}}{p^2-M^2_{(0)n_5,n_6}}\,.\end{aligned} \tag{4.6}$$

Therefore, considering the similar Feynman diagrams as in 5d [10, 18], in the dimensional regularization with $d=4-\epsilon$, bosonic and fermionic loop contributions to the scalar self energy are, at nonzero external momentum q^2, respectively,

$$-im_B^2(q^2)=4g_4^2Q^2\mu^{4-d}\sum_{n_5,n_6\in\mathbf{Z}}\int\frac{d^dp}{(2\pi)^d}\frac{p(q+p)}{(p^2-M^2_{(0)n_5,n_6})(q+p)^2} \tag{4.7}$$

and

$$-im_F^2(q^2)=-4g_4^2Q^2\mu^{4-d}\sum_{n_5,n_6\in\mathbf{Z}}\int\frac{d^dp}{(2\pi)^d}\frac{p(q+p)}{(p^2-M^2_{n_5,n_6})(q+p)^2}\,. \tag{4.8}$$

By using the Schwinger representation

$$\frac{1}{A^n}=\frac{1}{\Gamma(n)}\int_0^\infty dt\,t^{n-1}e^{-At}\,, \tag{4.9}$$

and performing the momentum integrations via the identities

$$\int_0^\infty dy\,y^{2n+d-1}e^{-y^2t}=\frac{\Gamma(d/2+n)}{2t^{d/2+n}}\,, \tag{4.10}$$

we find the one-loop corrections as

$$m_B^2(q^2) = \frac{g_4^2 Q^2 (\mu\pi R_5)^\epsilon}{4\pi^3 R_5^2} \int_0^1 dx \left[\left(2 - \frac{\epsilon}{2}\right) \mathcal{J}_2[0,0,c] + \pi x(1-x) q^2 R_5^2 \mathcal{J}_1[0,0,c]\right] \tag{4.11}$$

and

$$\begin{aligned} m_F^2(q^2) &= -\frac{g_4^2 Q^2 (\mu\pi R_5)^\epsilon}{4\pi^3 R_5^2} \int_0^1 dx \, \times \\ &\times \left[\left(2 - \frac{\epsilon}{2}\right) \mathcal{J}_2[\omega_5, \omega_6, c] + \pi x(1-x) q^2 R_5^2 \mathcal{J}_1[\omega_5, \omega_6, c]\right] \end{aligned} \tag{4.12}$$

with

$$\begin{aligned} \mathcal{J}_j[\omega_5, \omega_6, c] &\equiv \sum_{n_5, n_6 \in \mathbf{Z}} \int_0^\infty \frac{dt}{t^{j-\epsilon/2}} e^{-\pi t [c + a_5 (n_5 + \omega_5)^2 + a_6 (n_6 + \omega_6)^2]}, \qquad j = 1, 2; \quad a_{5,6}, c > 0; \\ a_5 &\equiv x, \qquad a_6 \equiv x \left(\frac{R_5}{R_6}\right)^2, \qquad c \equiv -x(1-x) q^2 R_5^2. \end{aligned} \tag{4.13}$$

For small positive[4] c, we obtain the following approximate formulas [25] for $\mathcal{J}_j[\omega_5, \omega_6, c]$,

$$\begin{aligned} \mathcal{J}_1[\omega_5, \omega_6, c \ll 1] &\simeq \frac{\pi c}{\sqrt{a_5 a_6}} \left[\frac{-2}{\epsilon}\right] - \ln \left| \frac{\vartheta_1(\omega_6 - iu\omega_5 | iu)}{(\omega_6 - iu\omega_5) \eta(iu)} e^{-\pi u \omega_5^2} \right|^2 - \\ &- \ln \left[\frac{(c + a_5 \omega_5^2 + a_6 \omega_6^2)}{a_6}\right], \qquad u \equiv \sqrt{\frac{a_5}{a_6}}, \end{aligned} \tag{4.14}$$

$$\begin{aligned} \mathcal{J}_2[\omega_5, \omega_6, c \ll 1] &\simeq -\frac{\pi^2 c^2}{2\sqrt{a_5 a_6}} \left[\frac{-2}{\epsilon}\right] + \frac{\pi^2 a_5}{3} \left(\frac{a_5}{a_6}\right)^{1/2} \left[\frac{1}{15} - 2\Delta_{\omega_5}^2 (1 - \Delta_{\omega_5})^2\right] + \\ &+ \left[\sqrt{a_5 a_6} \sum_{n \in \mathbf{Z}} |n + \omega_5| \mathrm{Li}_2(e^{-2\pi i z}) + \text{c.c.}\right] + \\ &+ \left[\frac{a_6}{2\pi} \sum_{n \in \mathbf{Z}} \mathrm{Li}_3(e^{-2\pi i z}) + \text{c.c.}\right] \end{aligned} \tag{4.15}$$

where $\Delta_{\omega_5} \equiv \omega_5 - [\omega_5]$ with $0 \leq \Delta_{\omega_5} < 1$ and $[\omega_5] \in \mathbf{Z}$, and $z \equiv \omega_6 - i\sqrt{\frac{a_5}{a_6}} |n + \omega_5|$. Here, ϑ_1 is the Jacobi theta function and η is the Dedekind eta function. And $\mathrm{Li}_2, \mathrm{Li}_3$ are the polylogarithm functions as

$$\mathrm{Li}_n(x) = \sum_{k=1}^\infty \frac{x^k}{k^n}, \qquad n = 2, 3. \tag{4.16}$$

Therefore, the resulting one-loop correction for the brane scalar is given by

$$\begin{aligned} m_\phi^2(q^2) &= m_B^2(q^2) + m_F^2(q^2) \\ &= \frac{g_4^2 Q^2}{2\pi^3 R_5^2} \int_0^1 dx \left[\mathcal{J}_2[0,0,c] - \mathcal{J}_2[\omega_5, \omega_6, c]\right] + \end{aligned}$$

[4] c is positive after a Wick rotation $q^2 = -q_E^2$.

$$+\frac{g_4^2 Q^2}{4\pi^2 R_5^2}(q^2 R_5^2)\int_0^1 dx\, x(1-x)\Big[\mathcal{J}_1[0,0,c]-\mathcal{J}_1[\omega_5,\omega_6,c]\Big]\,. \tag{4.17}$$

Consequently, we observe that both divergences of $\mathcal{J}_1$ and $\mathcal{J}_2$ are cancelled and there appear only finite corrections. Thus, we can take $c=0$ safely at the zero external momentum without involving the UV and IR mixing found in [25]. So, the mass correction with $q^2=0$ is given by

$$\begin{aligned} m_\phi^2(0) \;=\; & \frac{g_4^2 Q^2}{4\pi^3 R_5^2}\Bigg[\frac{2\pi^2}{3} r\Delta_{\omega_5}^2(1-\Delta_{\omega_5})^2+\frac{1}{r}(I_1(0,0)-I_1(\omega_5,\omega_6)+\text{c.c.})+ \\ & +\frac{1}{2\pi r^2}(I_2(0,0)-I_2(\omega_5,\omega_6)+\text{c.c.})\Bigg]\,, \qquad r\equiv\frac{R_6}{R_5} \end{aligned} \tag{4.18}$$

with

$$I_1(\omega_5,\omega_6)\equiv\sum_{n\in\mathbf{Z}}|n+\omega_5|\mathrm{Li}_2(e^{-2\pi|n+\omega_5|r-2\pi i\omega_6})\,, \tag{4.19}$$

and

$$I_2(\omega_5,\omega_6)\equiv\sum_{n\in\mathbf{Z}}\mathrm{Li}_3(e^{-2\pi|n+\omega_5|r-2\pi i\omega_6})\,. \tag{4.20}$$

In order to see the mass correction explicitly, let us simplify the sums as follows,

$$I_1(\omega_5,\omega_6)=\frac{1}{2}\sum_{k=1}^{\infty}\frac{e^{-2\pi ik\Delta_{\omega_6}}}{k^2\sinh^2(\pi kr)}\Big[\Delta_{\omega_5}\cosh(2\pi k(1-\Delta_{\omega_5})r)+(1-\Delta_{\omega_5})\cosh(2\pi k\Delta_{\omega_5}r)\Big]\,, \tag{4.21}$$

and

$$I_2(\omega_5,\omega_6)=\sum_{k=1}^{\infty}\frac{e^{-2\pi ik\Delta_{\omega_6}}}{k^3}\frac{\cosh(\pi k(1-2\Delta_{\omega_5})r)}{\sinh(\pi kr)}\,. \tag{4.22}$$

Here $\Delta_{\omega_6}\equiv\omega_6-[\omega_6]$ with $0\le\Delta_{\omega_6}<1$ and $[\omega_6]\in\mathbf{Z}$. Therefore, inserting the above expressions into eq. (4.18), we find that the resulting mass correction is finite as

$$\begin{aligned} m_\phi^2(0) \;=\; & \frac{g_4^2 Q^2}{4\pi^3 R_5^2}\Bigg[\frac{2\pi^2}{3} r\Delta_{\omega_5}^2(1-\Delta_{\omega_5})^2+ \\ & +\frac{1}{r}\sum_{k=1}^{\infty}\frac{1}{k^2\sinh^2(\pi kr)}\Big(1-\cos(2\pi k\Delta_{\omega_6})\{\Delta_{\omega_5}\cosh(2\pi k(1-\Delta_{\omega_5})r)+ \\ & +(1-\Delta_{\omega_5})\cosh(2\pi k\Delta_{\omega_5}r)\}\Big)+ \\ & +\frac{1}{\pi r^2}\sum_{k=1}^{\infty}\frac{1}{k^3\tanh(\pi kr)}\Big(1-\cos(2\pi k\Delta_{\omega_6})\frac{\cosh(\pi k(1-2\Delta_{\omega_5})r)}{\cosh(\pi kr)}\Big)\Bigg]\,. \end{aligned} \tag{4.23}$$

First let us consider the case with $\Delta_{\omega_5}=0$. Then, eq. (4.23) becomes

$$m_\phi^2(0) \;=\; \frac{g_4^2 Q^2}{4\pi^3 R_5^2}\Bigg[\frac{1}{r}\sum_{k=1}^{\infty}\frac{1}{k^2\sinh^2(\pi kr)}(1-\cos(2\pi k\Delta_{\omega_6}))+$$

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$$+\frac{1}{\pi r^2}\sum_{k=1}^{\infty}\frac{1}{k^3\tanh(\pi kr)}(1-\cos(2\pi k\Delta_{\omega_6}))\Bigg]. \tag{4.24}$$

In this case, let us take the limit of $\pi r \gg 1$, i.e. one extra dimension with radius R_5 to be much smaller than the other. Thus, the resulting mass correction reproduces exactly the 5d case with a SS twist [20],

$$m_\phi^2(0) \simeq \frac{g_4^2 Q^2}{4\pi^4 R_6^2}\sum_{k=1}^{\infty}\frac{1}{k^3}(1-\cos(2\pi k\Delta_{\omega_6}))\,. \tag{4.25}$$

On the other hand, when one takes $\Delta_{\omega_5}=\frac{1}{2}$, which is the case with a discrete twist in the fifth direction, eq. (4.23) becomes

$$m_\phi^2(0) = \frac{g_4^2 Q^2}{4\pi^3 R_5^2}\Bigg[\frac{\pi^2}{24}r+\frac{1}{r}\sum_{k=1}^{\infty}\frac{1}{k^2\sinh^2(\pi kr)}+ + \frac{1}{\pi r^2}\sum_{k=1}^{\infty}\frac{1}{k^3\tanh(\pi kr)}\left(1-\cos(2\pi k\Delta_{\omega_6})\frac{\pi kr}{\sinh(\pi kr)}\right)\Bigg]. \tag{4.26}$$

Again in the limit of $\pi r \gg 1$, the resulting mass correction is

$$m_\phi^2(0) \simeq \frac{g_4^2 Q^2}{96\pi R_5^2}\left(\frac{R_6}{R_5}\right)\left[1+\mathcal{O}\left(\frac{R_5^2}{R_6^2}\right)\right]. \tag{4.27}$$

Therefore, in this case, one extra dimension with small radius R_5 is not decoupled, but rather the effect due to the nontrivial SS twist in that direction is a dominant contribution to the mass correction. For other nonzero values of Δ_{ω_5}, such a non-decoupling of small extra dimension remains true because the first term in eq. (4.23) is dominant for $\pi r \gg 1$.

5. Conclusion

We considered supersymmetry breaking on the orbifold T^2/Z_2 via the SS twisted boundary conditions or the localized mass terms. It turns out that the SS breaking is equivalent to the localized breaking at the lines which should be regarded to be fixed boundaries under additional Z_2 actions. In this case, we have shown that in the presence of the SS twist or localized mass terms for the bulk gauge sector, there arises a finite one-loop mass correction to the visible brane scalar. In particular, for the case with one extra dimension much smaller than the other, we observe that the effect from the small extra dimension to the one-loop mass correction is not decoupled due to a nontrivial SS twist in that direction.

In order to know whether the contribution due to the bulk gaugino dominates over other contributions such as anomaly mediation [8], one needs to determine the SS twist parameter dynamically. At the level of 4d effective supergravity, one could think of the SS breaking to be equivalent to a nonzero F term of the corresponding radion multiplet for two extra dimensions as in 5d case [8], and introduce a radius stabilization mechanism to determine the F term dynamically. Moreover, in order to estimate supergravity loop corrections as in 5d case [26], it seems indispensible to understand the 6d off-shell supergravity, which is not available yet. Let us leave these issues in a future publication.

Acknowledgments

The author would like to thank W. Buchmüller, A. Falkowski and D. Ghilencea for useful comments and discussion.

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