Radius-dependent gauge unification in AdS5

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(November 10, 2018)

Abstract

We examine the relation of the 4-dimensional low energy coupling of bulk gauge boson in a slice of AdS5 to the 5-dimensional fundamental couplings as a function of the orbifold radius $R$. This allows us to address the gauge coupling unification in AdS5 by means of the radius running as well as the conventional momentum running. We then compute the radius dependence of 1-loop low energy couplings in generic AdS5 theory with 4-dimensional supersymmetry, and discuss the low energy predictions when the 5-dimensional couplings are assumed to be unified.
It has been noted that the large scale hierarchy between the weak and Planck scales can be naturally obtained in 5-dimensional (5D) theory on a slice of AdS5 \([1]\) with an appropriate radion stabilization mechanism \([2]\). In the original model of Randall and Sundrum, all the standard model fields are assumed to be confined in the TeV brane. An apparent drawback of this scenario is that one has to abandon the perturbative unification of gauge couplings at the fundamental scale of the model. An alternative scenario which may achieve gauge unification while solving the hierarchy problem is that the gauge fields propagate in 5D bulk spacetime \([3]\). In such case, the size of gauge coupling renormalization is generically of order 

$$[\ln(M_{Pl}/M_W)^2]/8\pi^2 \quad [4,5,6,7,8]$$

so the gauge unification can be achieved at a scale near the 4D Planck scale \(M_{Pl}\). In this paper, we first point out that it is convenient to consider the orbifold radius \(R\)-dependence of 4D couplings in addition to the momentum-dependence in order to address the unification of bulk gauge couplings in AdS5. We then compute the \(R\)-dependence of 1-loop 4D couplings in generic AdS5 theory with \(N = 1\) supersymmetry (SUSY) which is orbifolded by \(Z_2 \times Z_2'\) \([9]\), and examine the low energy consequences of unified 5D couplings.

Let us consider 5D gauge theory on a slice of AdS5 with orbifold radius \(R\). The action includes

$$\int d^4x dy \sqrt{G} \left(-\frac{1}{4g_5^2} F_{a}^{MN} F_{a}^{MN} - \sum_i \frac{\delta(y - n_i\pi)}{\sqrt{G_{55}}} \frac{1}{4g_i^2} F_{a \mu
u} F_{a}^{\mu\nu}\right), \quad (1)$$

where \(y = n_i\pi\) \((n_i = 0, 1)\) denote the 5-th coordinates of orbifold fixed points and the 5D metric \(G_{MN}\) is given by

$$ds^2 = G_{MN} dx^M dx^N = e^{-2kR |y|} g_{\mu\nu} dx^\mu dx^\nu + R^2 dy^2, \quad (2)$$

where \(k\) is the AdS curvature. For generic bulk fields in AdS5, Kaluza-Klein (KK) scale is set by \(M_{KK} \simeq \pi k/(e^{\pi kR} - 1) \quad [3]\). The quantity of our interest is the 1-loop low energy coupling of zero mode gauge boson with external 4D momentum \(p \lesssim M_{KK}\) for generic value of \(kR\):

$$\frac{1}{g^2_a(p, R, k)} = \left[\frac{1}{g_{5a}^2(\Lambda)} + \frac{\gamma_a}{8\pi^3} \Lambda\right] \pi R + \left[\sum_i \frac{1}{g_{ia}^2(\Lambda)} + \frac{b'_a}{8\pi^2} \ln \Lambda\right] + \frac{1}{8\pi^2} \tilde{\Delta}_a(p, R, k), \quad (3)$$
where $\Lambda$ is the cutoff scale measured by $G_{MN}$ and $p^2 = -g^{\mu\nu}\partial_\mu\partial_\nu$ is measured by $g_{\mu\nu}$. Here the linear divergence is from the bulk counter term $[10]$, while the log divergence is from the fixed point counter terms $[11][12]$, and the conventional momentum-running and finite KK threshold corrections are encoded in $\tilde{\Delta}_a$. In this paper, we focus on

$$\Delta_a = b'_a \ln \Lambda + \tilde{\Delta}_a$$

which is unambiguously calculable within 5D effective field theory, not on the uncalculable bare parameters $1/g^2_{ia}$ and $1/\hat{g}^2_{ia} \equiv 1/g^2_{ia} + \gamma_a \Lambda/8\pi^3$. In fact, $1/g^2_{ia}$ can be simply ignored under the strong coupling assumption $g^2_{ia}(\Lambda) = O(8\pi^2)$ $[13]$.

Let us first summarize some generic features of the 1-loop correction (4). Since the UV divergence structure is independent of $k$, $b'_a$ can be computed in the flat limit $k = 0$, yielding $[11]$

$$b'_a = \frac{1}{12} \left[ \sum T_a(\phi^{++}) - \sum T_a(\phi^{--}) - 23 \sum T_a(A^{++}_M) + 23 \sum T_a(A^{--}_M) \right],$$

where $T_a(\Phi) = \text{Tr}(T_a^2)$. Here $\phi^{zz'}$ and $A^{zz'}_M (z, z' = \pm)$ are the 5D real scalar and vector fields with the boundary conditions $\Phi^{zz}(-y) = z \Phi^{zz}(y)$, $\Phi^{zz'}(-y + \pi) = z' \Phi^{zz'}(y + \pi)$. When $p \lesssim M_{KK}$, the $p$-dependence of $1/g^2_a$ is given by $\partial \Delta_a/\partial \ln p = -b_a + O(p^2/M^2_{KK})$ where $b_a$ is the conventional 1-loop beta function coefficient in 4D effective theory. Also for $p \ll 1/R \ll \Lambda$ with $k = 0$, one finds

$$(\Delta_a)_{k=0} = b'_a \ln(\Lambda R) - b_a \ln(pR) + O(1).$$

In large radius limit $\pi k R \gg 1$, the $p$-dependence of $1/g^2_a$ can be determined within 5D field theory only for $p \lesssim e^{-k\pi R}\Lambda$. This can be easily seen for instance by considering the effects of higher-derivative terms in 5D lagrangian density, e.g. $F^{aMN}G^{PQ}\partial_P\partial_QF^{a}_{MN}/\Lambda$. Such term gives a contribution of order $p^2/e^{-2k\pi R}k\Lambda$ to $1/g^2_a$, which means that 5D field theory description breaks down for $p \gtrsim e^{-k\pi R}\sqrt{k\Lambda} [4]$. So one can not probe a possible gauge unification at $\Lambda$ by means of the momentum-running alone. Physically, this is to be expected since the gauge field zero mode is constant along $y$, so its amplitude receives an important
contribution from $y \approx \pi$ which has the cutoff $\sim e^{-\pi kR}\Lambda$. On the other hand, in small radius limit $\pi kR \ll 1$, the leading $p$-dependence is calculable within the 5D field theory of (II) as long as $p \lesssim \mathcal{O}(\Lambda)$. In particular, when $k \approx 1/\pi R \approx \Lambda$, we have $(\Delta a)_{k \approx 1/\pi R} \approx b_a \ln(\Lambda/p)$.

With the above observation, the gauge coupling renormalization in AdS5 can be described by Fig.1. First of all, the range of $[\ln(p/M_{Pl}), \pi kR]$ which allows a 5D field theory description is bounded to be below the line A representing $\ln(p/M_{Pl}) \approx -\pi kR$. The momentum-running of $g_a^2$ when $\pi kR \gg 1$ is allowed only for $p \lesssim e^{-\pi kR}\Lambda$ (the line 1), while for $\pi kR \ll 1$, the momentum-running is allowed up to $p = \mathcal{O}(\Lambda)$ (the line 3). At $p \approx \Lambda \approx 1/\pi R \approx k$, we have $1/g_a^2 \approx \pi R/g_{\text{Rad}}^2$. From this point, one can move along the dotted lines to arrive at the phenomenologically relevant point with $p \approx M_W$ and $\pi kR \gg 1$. This procedure involves always a radius-running along the line 2, so it is crucial to compute the $R$-dependence of $g_a^2$ over the range from $\pi kR \ll 1$ to $\pi kR \gg 1$ in order to determine $g_a^2$ at $p \approx M_W$ and $\pi kR \gg 1$. This suggests also that the gauge unification in AdS5 can be addressed by means of the double running along $\ln p$ and $R$. Suppose that the 5D bare couplings $g_{\text{Rad}}^2$ at $\Lambda$ have a unified value. Then the resulting prediction on $g_a^2$ at $p \approx M_W$ and $\pi kR \gg 1$ can be unambiguously computed by means of the $\ln p$ and $R$ runnings.

It is possible to directly compute the $R$-dependence of $g_a^2$ in generic AdS5 theory [14]. However in case with unbroken $N = 1$ SUSY, there is much simpler way to compute the $R$-dependence. In SUSY case, the radion $R$ forms a $N = 1$ superfield together with the 5-th component of the graviphoton $B_M$. In 4D effective supergravity (SUGRA), the field-dependence of gauge couplings is determined by the field-dependence of holomorphic gauge kinetic function $f_a$ and Kähler potential $K$ which can be expanded in powers of generic charged superfield $\Phi$: $K = K_0(T,T^*) + Z_\Phi(T,T^*)\Phi^*e^{-V}\Phi$, where $T$ denotes generic gauge singlet moduli superfield. Then the moduli-dependence of 1-loop low energy couplings are unambiguously determined to be [15]

$$
\frac{1}{g_a^2(p)} = \text{Re}(f_a) + \frac{b_a}{16\pi^2} \ln \left( \frac{M_{Pl}^2}{e^{-K_0/3}p^2} \right) - \sum_{\Phi} T_a(\Phi) \ln \left( e^{-K_0/3}Z_\Phi \right) + \frac{T_a(\text{Adj})}{8\pi^2} \ln (\text{Re}(f_a)),
$$

(7)
where $b_a = \sum_\Phi T_a(\Phi) - 3T_a(\text{Adj})$ and $M_{Pl}$ is the Planck scale of $g_{\mu\nu}$ which defines $p^2 = -g^{\mu\nu}\partial_\mu\partial_\nu$. With (7), one can determine the $R$-dependence of 1-loop couplings in AdS5 by computing the $R$-dependence of $f_a$ and $K$ in the corresponding 4D effective SUGRA. Obviously, (4) indicates that the 1-loop threshold corrections from massive KK modes are encoded in $f_a$, while the 4D field-theoretic loop effects of massless modes can be determined by the tree-level forms of $K$ and $f_a$. The $R$-dependent 1-loop $f_a$ appears to be the most nontrivial part to compute. However, $f_a$ is a holomorphic function of the radion superfield $T$ whose scalar component is given by $T = R + iB_5$, so its $R$-dependence can be determined by the $B_5$-dependence which is much easier to compute.

Let us consider a generic supersymmetric 5D theory on $S^1/Z_2 \times Z'_2$ with action

$$S = \int d^5x \sqrt{-G} \left[ \frac{M_5^3}{2} \left( R - \frac{3}{2} C_{MN} C^{MN} \right) + \frac{1}{g_{5a}^2} \left( \frac{1}{2} D_M \phi^a D^M \phi^a - \frac{1}{4} F^{aMN} F_{aMN}^a \right) + \frac{i}{2} \lambda^i a_M D_M \lambda^a_i + |D_M h_1^i|^2 + i \bar{\Psi}_I \gamma^M D_M \Psi_I + ic_I k\epsilon(y) \bar{\Psi}_I \Psi_I + ... \right]$$

(8)

where $R$ is the 5D Ricci scalar, $C_{MN} = \partial_M B_N - \partial_N B_M$ is the graviphoton field strength, $\phi^a, A_M^a$ and $\lambda^i a (i = 1, 2)$ are 5D scalar, vector and symplectic Majorana spinors constituting a 5D vector multiplet, $h_1^i$ and $\Psi_I$ are 5D scalar and Dirac spinor constituting the $I$-th hypermultiplet with kink mass $c_I k\epsilon(y)$. For nonzero $k$, $U(1)_R$ is gauged as

$$D_M h_1^i = \partial_M h_1^i - i \left( \frac{3}{2} (\sigma_3)^i_j - c_I \delta^i_j \right) k\epsilon(y) B_M h_1^j + ...$$

$$D_M \Psi_I = \partial_M \Psi_I + ic_I k\epsilon(y) B_M \Psi_I + ...$$

$$D_M \lambda^{ai} = \partial_M \lambda^{ai} - i \frac{3}{2} (\sigma_3)^i_j k\epsilon(y) B_M \lambda^{oj} + ...,$$

(9)

where the ellipsis stands for other gauge interactions. The 5D SUGRA multiplet is assumed to have the standard boundary conditions under $Z_2 : y \rightarrow -y$ and $Z'_2 : y + \pi \rightarrow -y + \pi$, leaving the 4D $N = 1$ SUSY unbroken. On the other hand, the vector and hypermultiplets can have arbitrary boundary conditions:

$$A_\mu^a(-y) = z_a A_\mu^a(y), \quad A_\mu^a(-y + \pi) = z'_a A_\mu^a(y + \pi),$$

$$h_1^i(-y) = z_I (\sigma_3)^i_j h_1^j(y), \quad h_1^i(-y + \pi) = z'_I (\sigma_3)^i_j h_1^j(y + \pi),$$

(10)
where \( z_{a,t}, z'_{a,t} = \pm 1 \) and the boundary conditions of other 5D fields are fixed by (10) in the standard manner. To derive the 4D effective SUGRA action, it is convenient to write the above 5D action in \( N = 1 \) superspace [10]. Among 5D gravity multiplet, we keep only \( T \) and replace other fields by their vacuum expectation values. Then following [16,17], we find (for \( M_5 = 1 \))

\[
S_1 = \int d^5x \left[ \int d^2\theta \left\{ \frac{1}{2} (T + T^*) e^{-(T + T^*)|y|} \left[ 1 + e^{(\frac{3}{2} + \epsilon)l(T + T^*)k|y|} H_I e^{-V} H_I \\
+ e^{(\frac{3}{2} + \epsilon)l(T + T^*)k|y|} H_I^* e^{V} H_I^* \right] + \frac{2}{g_5^2 a} e^{-(T + T^*)|y|} (\partial_y V^a - \frac{1}{\sqrt{2}}(\chi^a + \chi^{a*})^2) \right\} \\
+ \int d^2\theta \left\{ \frac{1}{4g_5^2} T W^{a\alpha} W_\alpha^a + H^c (\partial_y - \frac{1}{\sqrt{2}}\chi) H + h.c. \right\} \right],
\]

(11)

where \( W^a_\alpha \) is the chiral spinor superfield for the bulk gauge multiplet \( V^a = (A^a_\mu, \lambda^a) \) with \( \lambda^a = (1 - \gamma_5)\lambda^{a1}/2 \), and \( H_I = (h_I^1, \psi_I), H_I^c = (h_I^{2*}, \bar{\psi}_I^*), \chi^a = (\phi^a_5 + iA^a_5, \eta^a) \) are chiral superfields containing two-component fermions \( \psi_I = (1 - \gamma_5)\Psi_I/2, \bar{\psi}_I = (1 + \gamma_5)\Psi_I/2, \)
\( \bar{\eta}^a = (1 + \gamma_5)\lambda^{a2}/2 \). Note that \( H_I \) with \( z_I = z'_I = 1 \), \( H_I^c \) with \( z_I = z'_I = -1 \), \( V^a \) with \( z_a = z'_a = 1 \), \( \chi^a \) with \( z_a = z'_a = -1 \) can give massless 4D modes. The model can be easily generalized to include \( N = 1 \) superfields \( Q_{UV} \) \((Q_{IR})\) living on the UV (IR) brane at \( y = 0 \) (\( \pi \)). In fact, to rewrite (8) as (11), one needs to perform \( R \) and \( B_5 \)-dependent field redefinition, yielding an additional action through the chiral anomaly [18]:

\[
S_2 = \int d^5x d^2\theta \left\{ \frac{3}{4} T_a(\lambda^a) (z_b \delta(y) + z'_b \delta(y - \pi)) - \frac{1}{2} \epsilon c I T_a(\Psi_I) (z_I \delta(y) + z'_I \delta(y - \pi)) \\
- \frac{3}{2} T_a(\psi_{1IR}) \delta(y - \pi) \right\} (16\pi^2)^{-1} k|y| T W^{a\alpha} W_\alpha^a + h.c.,
\]

(12)

where \( \psi_{1IR} \) is the fermion component of \( Q_{1IR} \). Using the holomorphy property, this anomaly term can be easily determined by the following \( B_5 \)-dependent transformation of fermions:

\[
\lambda^{ai} \rightarrow e^{3ik|y|B_5/2} \lambda^{ai} , \quad \Psi_I \rightarrow e^{-ic_I k|y|B_5} \Psi_I , \quad \psi_{1IR} \rightarrow e^{-3ik\pi B_5/2} \psi_{1IR} .
\]

S_1 has been derived in [17] using different superfield basis for \( H_I \) and \( H_I^c \). A nice feature of our field basis is that \( B_5 \) does not have any non-derivative interaction in \( S_1 \) other than the Chern-Simons coupling. As a result, integrating out massive KK modes does not generate
1-loop $B_5$-coupling to $F\tilde{F}$ other than those in $S_2$, so no $R$-dependent 1-loop $f_a$ other than those from $S_2$. It is then straightforward to derive the Kähler metrics of the massless 4D fields from $H_I, H^c_I$ and of the brane fields $Q_{UV}, Q_{IR}$, and also the gauge kinetic function of the massless 4D gauge fields \[17,19,20\]:

\[\begin{align*}
Y_{H_I, H^c_I} &= 2M_5(e^{\left(\frac{1}{2} - z_{lc}I\right)\pi k(T + T^*)} - 1)/(1 - 2z_{lc}I)k, \\
Y_{Q_{IR}} &= e^{-\pi k(T + T^*)}, \quad Y_{Q_{UV}} = 1, \\
f_a &= \frac{\pi T}{g_5^2} - \frac{3}{8\pi^2} \left( \frac{1}{2} \sum_b z_b' T_a(\lambda^b) + \frac{3}{2} \sum_I z_I' c_I T_a(\Psi_I) \right) k\pi T, \\
\end{align*}\]

where $Y_\Phi = e^{-K_0/3}Z_\Phi$ for the Kähler metric $Z_\Phi$, and $M_5^2 = e^{-K_0/3}M_5^2 = (1 - e^{-\pi k(T + T^*)})M_5^2/k$. The calculation of $Y_{\chi^b}$ for the massless 4D fields from $\chi^b$ involves the heavy tadpole with one 4D derivative, i.e. the tadpole of $A^b_\mu$ with $z_b = z'_b = -1$ \[13\]:

\[A^b_\mu = \partial_\mu A^b_5 \left[ y - \pi \frac{e^{(T + T^*)ky - 1}}{e^{(T + T^*)k\pi} - 1} \right],\]

yielding

\[Y_{\chi^b} = k/(e^{\pi k(T + T^*)} - 1)M_5.\]

Applying \[13\] and \[14\] to \[7\], we find

\[\begin{align*}
\Delta_a &= T_a(Q_{UV}) \ln(M_5/p) + T_a(Q_{IR}) \ln(M_5 e^{-\pi kR}/p) \\
&\quad - T_a(\text{Adj}) \left( 3 \ln(M_5/p) - 3\pi kR/2 - \ln(M_5R) \right) \\
&\quad + \sum_{z_I = z'_I} T_a(H_I) \left( \ln(k/p) - z_I c_I \pi kR - \ln \left[ (e^{(1-2z_Ic_I)\pi kR} - 1)/(\pi(1 - 2z_Ic_I)) \right] \right) \\
&\quad - \sum_{z_I = -z'_I} z'_I T_a(H_I) c_I \pi kR + \sum_{z_b = -z'_b} z'_b T_a(V^b) 3\pi kR/2 \\
&\quad + \sum_{z_b = z'_b = -1} T_a(\chi^b) \left( \ln(M_5^2/pk) + \pi kR/2 + \ln(1 - e^{-2\pi kR}) \right) \\
\end{align*}\]

which is valid for $p \lesssim M_{KK} \approx \pi k/(e^{\pi kR} - 1)$. The above result obtained by 4D effective SUGRA calculation can be confirmed by an explicit loop calculation summing all the loops of KK modes \[14\], which assures the validity of our 4D SUGRA calculation.
As a simple example to show the effects of radius-running, let us consider a supersymmetric model with the MSSM gauge and Higgs superfields living in 5D bulk spacetime with \( z_a = z'_a = z_I = z'_I = 1 \), the MSSM lepton superfields living on the UV brane and the MSSM quark superfields on the IR brane. We choose \( k = 5 \times 10^{17} \, \text{GeV}, M_5 = 1.5 \times 10^{18} \, \text{GeV} \), and assume the \( SU(5) \)-like boundary conditions of 5D couplings at \( M_5 \): \( \hat{g}_{SU(3)}^2 = \hat{g}_{SU(2)}^2 = \frac{5\hat{g}_{U(1)}^2}{3} \). We also choose \( c_I = 1/2 \) for the Higgs hypermultiplets, so that the Higgs zero modes are constant along \( y \). As we have noted, the unification of \( \hat{g}_{SU(a)}^2 \) implies that \( g_a^2 \) are unified at \( p \approx 1/R \approx M_5 \). If our universe has \( 1/R \approx M_5 \), the 4D effective theory for \( p \lesssim M_5 \) is the MSSM, and then this model can not be compatible with the observed low energy couplings which indicate that gauge couplings are unified at \( 2 \times 10^{16} \, \text{GeV} \), not at \( M_5 \approx 10^{18} \, \text{GeV} \). However if our universe has \( \pi k R \approx 10 \), the momentum running along the line 1 is allowed only for \( p \lesssim 7 \times 10^{13} \, \text{GeV} \). One then has to include the effects of radius running along the line 2, making the observed low energy couplings to be consistent with the 5D unification at \( M_5 \) as depicted in Fig.2.

To conclude, we have pointed out that the gauge coupling renormalization in AdS5 can be studied by means of the double running along \( \ln p \) and \( R \), as illustrated in Fig. 1. It is then crucial to compute the \( R \)-dependence of \( g_a^2 \) to address the issue of gauge unification in AdS5. Using the gauged \( U(1)_R \) and chiral anomaly in 5D SUGRA and also the holomorphic property of 4D effective SUGRA, we could compute the \( R \)-dependence of 1-loop 4D couplings in generic AdS5 theory with \( N = 1 \) SUSY which is orbifolded by \( Z_2 \times Z'_2 \). Our result (15) can be used to study gauge unification in supersymmetric AdS5 model.
FIGURES

FIG. 1. The domain of $[\ln p, \pi kR]$ which can be described by 5D effective field theory. The line A represents $\ln(p/M_5) \approx -\pi kR$.

FIG. 2. The running of 1-loop $\Delta_a$. The region 1 represents the momentum running up to $p \approx 10^{13}$ GeV along the line 1 of Fig. 1, the region 2 is the radius running from $\pi kR = 10$ to $\pi kR = 1$ along the line 2, and the region 3 is the momentum running to $p \approx M_5$ along the line 3.
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