Integrating out Holographic QCD back to Hidden Local Symmetry

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We develop a previously proposed gauge-invariant method to integrate out infinite towers of vector and axialvector mesons arising as Kaluza-Klein (KK) modes in a class of holographic models of QCD (HQCD). We demonstrate that HQCD can be reduced to the chiral perturbation theory (ChPT) with the hidden local symmetry (HLS) (so-called HLS-ChPT) having only the lowest KK mode identified as the HLS gauge boson, and the Nambu-Goldstone bosons. The $O(p^4)$ terms in the HLS-ChPT are completely determined by integrating out infinite towers of vector/axialvector mesons in HQCD: Effects of higher KK modes are fully included in the coefficients. As an example, we apply our method to the Sakai-Sugimoto model.

1. Introduction

Holography, based on gauge/gravity duality [1], has been of late fashion to reveal a part of features in strongly coupled gauge theories involving the application to QCD (so-called holographic QCD (HQCD)). There are two types of holographic approaches: One is called “top-down” approach starting with a stringy setting; the other is called “bottom-up” approach beginning with a five-dimensional gauge theory defined on an AdS (anti-de Sitter space) background. It is a key point to notice that in whichever approach one eventually employs a five-dimensional gauge model with a characteristic induced-metric and some boundary conditions on a brane configuration.

In the low-energy region, any model of HQCD is reduced to a certain effective hadron model in four-dimensions. Such effective models include vector and axialvector mesons as infinite towers of Kaluza-Klein (KK) modes together with the Nambu-Goldstone bosons (NGBs) associated with the spontaneous chiral symmetry.
breaking. Green functions in QCD are evaluated straightforwardly from the effective model following the holographic dictionary. Full sets of infinite towers of exchanges of KK modes (vector and axialvector mesons) contribute to Green functions involving current correlators and form factors and could mimic ultraviolet behaviors in QCD, although such a hadronic description would not be reliable above a certain high-energy scale. This implies that appropriate/gauge-invariant holographic results require calculations including full set of KK towers, which, however, would not be practical because of forms written in terms of an infinite sum.

It was pointed out [2] that the infinite tower of KK modes is interpreted as a set of gauge bosons of the hidden local symmetries (HLSs) [3, 4]. This implies an interesting possibility that, in the low-energy region, any holographic models can be reduced to the simplest HLS model, provided that the infinite tower of KK modes is integrated out keeping only the lowest one identified with the $\rho$ meson and its flavor partners. Effects from the higher KK modes would then be fully incorporated into higher derivative terms ($O(p^4)$ terms) in the HLS effective field theory as an extension of the conventional chiral perturbation theory (ChPT) [5], so-called the HLS-ChPT [4, 6] which is manifestly gauge-invariant formulation and makes it possible to calculate any Green functions order by order in derivative expansion. Once holographic models are expressed in terms of the HLS-ChPT, one can even calculate meson-loop corrections of subleading order in $1/N_c$ expansion. This would give a new insight into the HQCD which as it stands is valid only in the large $N_c$ limit. Indeed, in the previous work [7], we proposed a consistent method of integrating out the infinite towers of vector and axialvector mesons in the Sakai-Sugimoto (SS) model [8, 9] into the HLS-ChPT, with the $O(p^4)$ terms explicitly given by the integrated higher modes effects. Then we were able to do the first calculations of $1/N_c$ corrections to the SS model.

In this talk, we report the results of our work [12] developing in detail the integrating-out method proposed in Ref. [7]. First of all, we work in a class of holographic models to introduce our integrating-out method. Next, as an example, we apply our procedure to the Sakai-Sugimoto (SS) model [8, 9] to give the HLS model with a full set of $O(p^4)$ terms determined. We then calculate the pion electromagnetic form factor to demonstrate how powerful our formulation is even before including the $1/N_c$ effects through loop. As we will see explicitly, the momentum-dependence of the form factor is evaluated including full set of contributions from KK modes without performing infinite sums. Our method can straightforwardly be applied to other types of holographic models such as those given in Refs. [10, 11]. More details are presented in Ref. [12].

2. A gauge-invariant way to integrate out HQCD

In this section, starting with a class of HQCD models including the SS model [8, 9], we introduce a way to obtain a low-energy effective model in four-dimension described only by the lightest vector meson identified as $\rho$ meson based on the
HLS together with the NGBs. Suppose that the fifth direction, spanned by the coordinate $z$, extends from minus infinity to plus infinity $(-\infty < z < \infty)$\(^a\). We employ a five-dimensional gauge theory which has a vectorial $U(N)$ gauge symmetry defined on a certain background associated with the gauge/gravity duality. As far as gauge-invariant sector such as the Dirac-Born-Infeld part of the SS model [8, 9] is concerned, the five-dimensional action in large $N_c$ limit can be written as\(^b\)

$$S_5 = N_c \int d^4xdz \left(-\frac{1}{2} K_1(z) \text{tr}[F_{\mu\nu}F^{\mu\nu}] + K_2(z) M_{KK}^2 \text{tr}[F_{\mu z}F^{\mu z}]\right), \quad (1)$$

where $K_{1,2}(z)$ denote a set of metric-functions of $z$ constrained by the gauge/gravity duality. $M_{KK}$ is a typical mass scale of KK modes of the gauge field $A_M$ with $M = (\mu, z)$. The boundary condition of $A_M$ is chosen as $A_M(x^\mu, z = \pm \infty) = 0$. A transformation which does not change this boundary condition satisfies $\partial_M g(x^\mu, z)|_{z=\pm \infty} = 0$, where $g(x^\mu, z)$ is the transformation matrix of the gauge symmetry. This implies an emergence of global chiral $U(N)_L \times U(N)_R$ symmetry in four-dimension characterized by the transformation matrices $g_{R,L} = g(z = \pm \infty)$.

Following Refs. [7–9], we work in $A_z = 0$ gauge. There still exists a four-dimensional gauge symmetry under which $A_\mu(x^\mu, z)$ transforms as $A_\mu \rightarrow h \cdot A_\mu \cdot h^\dagger - i\partial_\mu h \cdot h^\dagger$ with $h = h(x^\mu)$. This gauge symmetry is identified [7–9] with the HLS [3, 4]. In the $A_z = 0$ gauge, the NGB fields $\pi(x^\mu)$ disappear from the chiral field $U = e^{i2\pi/F_z}$ since $U \rightarrow 1$. They are instead included [7] in $A_\mu(x^\mu, z)$ at the boundary as $A_\mu|_{z=\pm \infty} = \alpha^{R,L}_\mu = i\xi_{R,L} \partial_\mu \xi^\dagger_{R,L}$, where $\xi_{L,R}$ form the chiral field $U$ as $U = \xi^\dagger_L \cdot \xi_R$. Since $\xi_{L,R} \rightarrow h \cdot \xi_{L,R} \cdot g_{L,R}$ [3, 4], $\alpha^{R,L}_\mu$ transform as $\alpha^{R,L}_\mu \rightarrow h \cdot \alpha^{R,L}_\mu \cdot h^\dagger - i\partial_\mu h \cdot h^\dagger$.

We introduce infinite towers of massive KK modes for vector ($V^{(n)}_\mu(x^\mu)$) and axial-vector ($A^{(n)}_\mu(x^\mu)$) meson fields, where we treat $V^{(n)}_\mu$ as the HLS gauge fields transforming under the HLS as $V^{(n)}_\mu \rightarrow h \cdot V^{(n)}_\mu \cdot h^\dagger - i\partial_\mu h \cdot h^\dagger$, while $A^{(n)}_\mu$ as matter fields transforming as $A^{(n)}_\mu \rightarrow h \cdot A^{(n)}_\mu \cdot h^\dagger$. The five-dimensional gauge field $A_\mu(x^\mu, z)$ is now expanded as

$$A_\mu(x^\mu, z) = \alpha^R_\mu(x^\mu) \phi^R(z) + \alpha^L_\mu(x^\mu) \phi^L(z) + \sum_{n=1}^{\infty} \left(A^{(n)}_\mu(x^\mu) \psi_{2n-1}(z) - V^{(n)}_\mu(x^\mu) \psi_{2n-1}(z)\right). \quad (2)$$

The functions $\{\psi_{2n-1}(z)\}$ are the eigenfunctions satisfying the eigenvalue equation obtained from the action $[11]: -K^{-1}_1(z) \partial_\zeta (K_2(z) \partial_\zeta \psi_n(z)) = \lambda_n \psi_n(z)$, where $\lambda_n$ denotes the $n$th eigenvalue. On the other hand, the gauge-invariance requires the functions $\phi^{R,L}(z)$ to be different from the eigenfunctions: From the trans-
formation properties for $A_\mu(x^\mu, z)$, $\alpha_{R,L}^{(n)}$, and $V_\mu^{(n)}$ we see that the functions $\phi^{(R,L)}(z)$, $\{\psi_{2n-1}(z)\}$, and $\{\psi_{2n}(z)\}$ are constrained as

$$\phi^R(z) + \phi^L(z) - \sum_{n=1}^{\infty} \psi_{2n-1}(z) = 1. \quad (3)$$

Using this, we may rewrite Eq. (2) to obtain $A_\mu(x^\mu, z) = \alpha_{\mu||}(x^\mu) + \alpha_{\mu\perp}(x^\mu)(\phi^R(z) - \phi^L(z)) + \sum_{n=1}^{\infty} A_{\mu}^{(n)}(x^\mu)\psi_{2n}(z) + \sum_{n=1}^{\infty} \left(\alpha_{\mu||}(x^\mu) - V_{\mu}^{(n)}(x^\mu)\right)\psi_{2n-1}(z)$, where $\alpha_{\mu||,\perp} = \frac{\alpha_{R,L}}{2}$ transform under the HLS as $\alpha_{\mu||} \rightarrow h \cdot \alpha_{\mu||} \cdot h^\dagger - i\partial_\mu h \cdot h^\dagger$ and $\alpha_{\mu\perp} \rightarrow h \cdot \alpha_{\mu\perp} \cdot h^\dagger$, respectively. Note that $\alpha_{\mu\perp}$ includes the NGB fields as $\alpha_{\mu\perp} = \frac{\alpha_{R,L}}{2}\partial_\mu \pi + \cdots$. The corresponding wave function $(\phi^R - \phi^L)$ should therefore be the eigenfunction for the zero mode, $\psi_0: \phi^R(z) - \phi^L(z) = \psi_0(z)$. From this and Eq. (3) we see that the wave functions $\phi^R$ and $\phi^L$ are not the eigenfunctions but are given as $\phi^{R,L}(z) = \frac{1}{2}[1 + \sum_{n=1}^{\infty} \psi_{2n-1}(z) \pm \psi_0(z)]$.

By substituting Eq. (2) with Eq. (3) into the action (1), the five-dimensional theory is now described by the NGB fields along with the infinite towers of the vector and axialvector meson fields in four dimensions.

We first naively try to truncate towers of the vector and axialvector meson fields simply eliminating $V_\mu^{(n)}$ and $A_{\mu}^{(n)}$ for $n > N$. Then we find

$$S_5^{\text{truncation}} \equiv \int dzd^4x K_2(z) \sum_{n=N+1}^{\infty} \lambda_{2n-1} \psi_{2n-1}(z) \text{tr}[\alpha_{\mu||}(x^\mu)]^2, \quad (4)$$

which explicitly breaks the chiral symmetry as well as the HLS, because $\alpha_{\mu||} \rightarrow h \cdot \alpha_{\mu||} \cdot h^\dagger - i\partial_\mu h \cdot h^\dagger$. Naive truncation of tower of vector meson fields thus forces us to encounter the explicit violation of the chiral symmetry.

Now we shall propose a method to truncate towers of vector and axialvector meson fields in a gauge-invariant manner. Consider a low-energy effective theory below the mass of $n = N + 1$ level. Such an effective theory can be obtained by integrating out mesons with $n \geq N + 1$ via the equations of motion. Neglecting terms including the derivatives acting on the heavy fields $V_\mu^{(k)}$ and $A_\mu^{(k)}$ with $k > N$, the equations of motion for them take the following forms: $V_\mu^{(k)} = \alpha_{\mu||}, A_\mu^{(k)} = 0$ ($k = N+1, N+2, \cdots, \infty$). Putting these solutions into the action, we have, instead of Eq. (4):

$$S_5^{\text{integrate out}} \equiv \int dzd^4x K_2(z) \sum_{n=1}^{N} \lambda_{2n-1} \psi_{2n-1}(z) \text{tr}[\alpha_{\mu||}(x^\mu) - V_{\mu}^{(n)}(x^\mu)]^2, \quad (5)$$

which is certainly gauge-invariant. Note also that the gauge-invariance now requires not the constraint in Eq. (3) but $\phi^R(z) + \phi^L(z) - \sum_{n=1}^{N} \psi_{2n-1}(z) = 1$.

Let us now consider a low-energy effective model obtained by integrating out all the higher vector and axialvector meson fields except the lowest vector meson field $V_{\mu}^{(1)} \equiv V_\mu$. Following the gauge-invariant way proposed above, the expansion of $A_\mu(x^\mu, z)$ is expressed as

$$A_\mu(x^\mu, z) = \alpha_{\mu\perp}(x^\mu)\psi_0(z) + (\tilde{\alpha}_{\mu||}(x^\mu) + V_\mu(x^\mu)) + \tilde{\alpha}_{\mu||}(x^\mu)\psi_1(z), \quad (6)$$
where $\bar{\alpha}_\mu = -V_\mu + \alpha_\mu$. One can further introduce the external gauge fields by gauging the global chiral $U(N)_L \times U(N)_R$ symmetry. (For details, see Ref. [12].) Then we obtain the low-energy effective model including only the lightest HLS field $V_\mu$ and the NGB fields $\pi$ described by the HLS-ChPT with $O(p^4)$ terms [4, 6]: The $O(p^4)$ terms include the effects from infinite towers of higher KK modes and are completely determined as explicitly shown in Ref. [12]; sum rules such as those introduced in Ref. [9] are also fully built in the HLS-ChPT Lagrangian. Our formulation is thus more practical and useful.

Finally, we once again emphasize that our methodology presented here is applicable to any models of HQCD.

3. Application to Sakai-Sugimoto Model

In this section, we apply our methodology to the Sakai-Sugimoto (SS) model [8, 9] based on $D8$/$D8$/$D4$ brane configuration. The five-dimensional gauge-invariant portion (so-called the Dirac-Born-Infeld (DBI) part) of the low-energy effective action in the SS model is given by [8, 9]

$$S_{SS}^{DBI} = N_c G \int d^4 x dz \left( -\frac{1}{2} K^{-1/3}(z) \text{tr}[F_{\mu\nu}F^{\mu\nu}] + K(z) M_{KK}^2 \text{tr}[F_{\mu\nu}F^{\mu\nu}] \right), \quad (7)$$

where $K(z) = 1 + z^2$ is the induced metric of the five-dimensional space-time; the overall coupling $G$ is the rescaled t Hooft coupling expressed as $G = N_c g_{YM}^2/(108\pi^3)$ with $g_{YM}$ being the gauge coupling of the $U(N_c)$ gauge symmetry on the $N_c$ D4-branes [8, 9]; the mass scale $M_{KK}$ is related to the scale of the compactification of the $N_c$ D4-branes onto the $S^1$. Comparing Eq. (7) with Eq. (1), we read off $K_1(z) = GK^{-1/3}(z)$, $K_2(z) = GK(z)$, so that we find the equation of motion,

$$-K^{-1/3}(z) \partial_z (K(z) \partial_z \psi_n) = \lambda_n \psi_n$$

with the eigenvalues $\lambda_n$ and the eigenfunctions $\psi_n$ of the KK modes of the five-dimensional gauge field $A_\mu(x^\mu, z)$.

Application to the Chern-Simons term is straightforward that is explicitly demonstrated in Ref. [12].

As emphasized in the end of the previous section, without introducing any sum rules, we are able to calculate amplitudes straightforwardly from the effective Lagrangian which includes contributions from full set of higher KK modes. To see it more explicitly, as an example, we shall study the pion electromagnetic (EM) form factor $F_\pi^\pm(Q^2)$ at tree-level of the present model. $F_\pi^\pm$ is readily constructed from the Lagrangian written in terms of the HLS-ChPT presented in Ref. [12]:

$$F_\pi^\pm(Q^2)|_{HLS} = g_{\gamma\pi\pi}(Q^2) + \frac{g_{\rho}(Q^2)m_{\pi}(Q^2)}{m_{\pi}^2 + Q^2},$$

where $Q^2 = -p^2$ denotes a momentum-squared in space-like region and $g_{\gamma\pi\pi}(Q^2)$ is given by $g_{\rho}(Q^2) = \frac{m_{\pi}^2}{g} \left( 1 + g_{\rho}(Q^2) \right)$. The applicable momentum range should be restricted to $0 \leq Q^2 \ll \{m_{\rho}^2, m_{\pi}^2, \cdots \}$ since we have integrated out higher KK modes keeping only the $\rho$ meson. Note that our form factor $F_\pi^\pm(0)|_{HLS}$ automatically ensures the EM gauge-invariance, $F_\pi^\pm(0)|_{HLS} = 1$. 


One can easily show \[12\] that if towers of vector and axialvector mesons had naïvely been truncated as in Eq. (4) one would have \(F_{\pi^\pm}(0)|_{\text{truncation}} \neq 1\), leading to a violation of the EM gauge-invariance. (It turns out \[12\] that higher KK modes actually play the crucial role to maintain the EM gauge-invariance.) We further rewrite \(F_{\pi^\pm}(Q^2)|_{\text{HLS}}\) as

\[
F_{\pi^\pm}(Q^2)|_{\text{HLS}} = \left(1 - \frac{1}{2} \hat{a}\right) + \bar{z} \frac{Q^2}{m_\rho^2} + \frac{g_\rho g_{\rho \pi \pi}}{m_\rho^2} + Q^2,
\]

where \(\hat{a} = a \left(1 - \frac{g^2 z_2}{2} - g^2 z_3 + \frac{(g^2 z_2)(g^2 z_3)}{2}\right)\) and \(\bar{z} = \frac{1}{4} a \left(g^2 z_6 + (g^2 z_3)(g^2 z_4)\right)\) \[12\] which are expressed in terms of the five-dimensional theory as \[3\] \[12\] \(\hat{a} = \frac{2g_\rho g_{\rho \pi \pi}}{m_\rho^2} = \frac{1}{8} \lambda_1 \left(\langle \psi_1 | (1 - \psi_2^0) \langle \psi_1 \rangle\right)\), \(\bar{z} = \frac{1}{8} \lambda_1 \left(\langle \psi_1 | (1 - \psi_2^0) \langle \psi_1 \rangle\right) - (1 - \psi_2^0)\). These \(\hat{a}\) and \(\bar{z}\) are calculated independently of any inputs to be \(\hat{a} \simeq 2.62\) and \(\bar{z} \simeq 0.08\), where we have used \(\lambda_1 \simeq 0.669\).

Using the expression \[8\] and the values of \(\hat{a}\) and \(\bar{z}\), we evaluate the momentum-dependence of \(F_{\pi^\pm}\) which was actually not possible in the original SS model \[9\] because of the form written in terms of the infinite summation. In Fig. 1 we show the predicted curve of \(F_{\pi^\pm}\) with respect to \(Q^2\) together with the experimental data from Ref. \[13\]. The \(\chi^2\)-fit results in good agreement with the data (\(\chi^2/\text{d.o.f} = 147/53 \simeq 2.77\)). Comparison with the result derived from the lowest vector meson dominance (LVMD) hypothesis with \(\hat{a} = 2\) and \(\bar{z} = 0\) is shown by a dashed curve in the left panel of Fig. 1. The right panel of Fig. 1 shows a comparison with the result obtained by fitting the parameters \((\hat{a}, \bar{z})\) to the experimental data. It is interesting to note that the best-fit values of \(\hat{a}\) and \(\bar{z}\) are quite close to those in the predicted curve. This fact reflects that the predicted curve fits well with the experimental data.

### 4. Summary

In this talk, we developed a methodology to integrate out arbitrary parts of infinite towers of vector and axialvector mesons arising as KK modes in a class of HQCD models. It was shown that our method is gauge-invariant in contrast to a naive truncation [See Eq. (4)]. It was demonstrated that any models of HQCD in the low-energy region can be described by the HLS-ChPT. We applied our method to the SS model and evaluated the momentum-dependence of the pion EM form factor as an example, which demonstrated power of our formulation. The predicted form factor was shown to be fitted well with the experimental data in the low-energy (space-like momentum) region. This was difficult in the original SS model due to the forms of the form factors written in terms of infinite sum of vector meson exchanges. More on phenomenological applications of our formulation to the SS model is presented in Ref. \[12\].

\[d\] Here \(g_\rho \equiv g_\rho(Q^2 = -m_\rho^2)\) and \(g_{\rho \pi \pi} \equiv g_{\rho \pi \pi}(Q^2 = -m_\rho^2)\).

\[e\] \(\langle A \rangle \equiv \int dz K^{-1/3}(z) A(z)\).
Fig. 1. The predicted curve of the pion EM form factor $F_{\pi}^{\mp}$ with respect to space-like momentum-squared $Q^2$ (denoted by solid line) fitted with the experimental data from Ref. [13] with $\chi^2$/d.o.f $= 147/53 \simeq 2.77$. In the left panel, the dashed curve corresponds to the form factor in the LVMD hypothesis of the HLS model with $\tilde{a} = 2$ and $\tilde{z} = 0$ taken. ($\chi^2$/d.o.f $= 226/53 \simeq 4.26$). The dashed curve in the right panel corresponds to the form factor fitted with the experimental data, yielding the best fit values of $\tilde{a}$ and $\tilde{z}$: $\tilde{a}_{\text{best}} = 2.44$, $\tilde{z}_{\text{best}} = 0.08$ ($\chi^2$/d.o.f $= 81/51 \simeq 1.56$).

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