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Holographic model for the baryon octet

Zhen Fang*

School of Physical Sciences, University of Chinese Academy of Sciences,
No. 19A Yuquan Road, Beijing 100049, China

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By adopting the nonlinear realization of chiral symmetry, a holographic model for the baryon octet is proposed. The mass spectra of the baryon octet and their low-lying excited states, which show good consistency with experiments, are calculated. The couplings of pion to nucleons are derived in two gauges and are shown to be equivalent with each other. It also shows that only derivative couplings of pion to nucleons appear in this holographic model. The coupling constant is then calculated.

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I. INTRODUCTION

In the recent decades, there have been many studies on anti-de Sitter (AdS)/QCD since the discovery of AdS/CFT [1–3]. Various holographic QCD models have been constructed, either in the top-down approach [4–6] or in the bottom-up approach [7–9]. Much research has been focused on the phenomenology of low-energy QCD [9–19], such as the hadron spectrum [10–13], the deconfining or chiral phase transitions [14–18], the hadron form factors [9], etc. The AdS/QCD method has shown powerful ability in the description of the low-energy QCD phenomenon.

In this paper, we focus on the baryon properties and try to construct a sensible holographic model for the baryon octet, based on the AdS/QCD model for mesons [7,8], which have shown superiority in the low energy QCD phenomenology. Indeed, considerable effort has been made in the construction of the holographic baryon models, such as in Refs. [20,21]. In Ref. [22], the authors built a holographic model for spin-$\frac{1}{2}$ baryons, which naturally incorporates the important phenomenon of the parity doublet [23]. However, one large limitation of this work is that they only considered the two-flavor case. To generalize the model, one cannot bypass the baryons with heavy quarks, which makes the issue about chiral symmetry troublesome. The meson physics in the three-flavor case has been considered in the holographic framework in Ref. [24]. For baryons with three flavors, chiral symmetry is believed to be broken from the beginning, and the picture of chiral symmetry breaking in Ref. [22] seems to be invalid.

Note that the holographic construction in Ref. [22] shares similarities with the linear sigma model and the massless modes of baryons before chiral symmetry breaking form the representation of the chiral group. However, for heavy baryons, it is more suitable to accept the nonlinear realization of chiral symmetry in the construction of baryon models [25]. Following this way, a holographic model for the baryon octet will be given. The whole mass spectra of the baryon octet and their low-lying excited states, which show good agreement with experiments, will be calculated. We also show that this holographic model only contains gradient couplings of the pion with baryon fields, at least in the chiral limit, which is in parallel with the nonlinear construction of the effective baryon model in the four-dimensional (4D) field-theoretic method [26].

II. OUTLINE OF THE MODEL

As in Refs. [7,8], we take a slice of AdS_5 space-time with the metric ansatz

$$ds^2 = e^{2A(z)}(\eta_{\mu\nu}dx^\mu dx^\nu - dz^2), \qquad \epsilon \le z \le z_m, \quad (1)$$

where $A(z) = -\ln z$, $\eta_{\mu\nu} = \mathrm{diag}(+,-,-,-)$ and the IR brane $z_m = 1/\Lambda_{\mathrm{QCD}}$.

To produce the possible parity-doubling pattern of excited baryon states, we follow Ref. [22] (for a brief review of this work, see the Appendix) and introduce two bulk baryon fields $B_{1,2}$ which correspond to the chiral baryon operators $\mathcal{O}_{L,R}$ in the boundary, respectively, but with the isodoublet in the two-flavor case replaced by the baryon octet in the three-flavor case:

$$B = \sum_{a=1}^{8}\frac{\lambda^a B^a}{\sqrt{2}} = \begin{pmatrix} \frac{\Sigma^0}{\sqrt{2}} + \frac{\Lambda}{\sqrt{6}} & \Sigma^+ & p \\ \Sigma^- & -\frac{\Sigma^0}{\sqrt{2}} + \frac{\Lambda}{\sqrt{6}} & n \\ \Xi^- & \Xi^0 & -\frac{2\Lambda}{\sqrt{6}} \end{pmatrix}. \quad (2)$$

The holographic model is formulated on the basis of the chiral gauge group $G = SU(3)_L \times SU(3)_R$, which is broken spontaneously into a vectorial subgroup $H = SU(3)_V$. The bulk action comprises two sectors. The first one, which contains the covariant kinetic terms and the bulk mass terms, can be written as

*fangzhen@ucas.ac.cn

$$S_{b1} = \int d^5x\sqrt{g}\sum_{j=1}^{2}\mathrm{Tr}(\bar{B}_j i e_A^M \Gamma^A \nabla_M B_j - m_5 \bar{B}_j B_j), \quad (3)$$

where the metric $g_{MN} = e_M^A e_N^B \eta_{AB}$ with the vielbein $e_M^A = \frac{1}{z}\eta_M^A$, the Dirac matrices $\Gamma^A = (\gamma^\mu, -i\gamma^5)$ satisfy $\{\Gamma^A, \Gamma^B\} = 2\eta^{AB}$, and the mass of the bulk baryon fields m_5 can be determined by the AdS/CFT correspondence as $m_5 = \pm\frac{5}{2}$ with the sign dictated by the analysis of the chirality of baryon fields (conventionally, we set $m_5 = \frac{5}{2}$ for B_1 and $m_5 = -\frac{5}{2}$ for B_2) [22]. The Lorentz and gauge covariant derivatives of the bulk baryon fields have the form

$$\nabla_M B_1 = \partial_M B_1 + \frac{i}{4}\omega_M^{AB}\Gamma_{AB}B_1 - i[\mathcal{L}_M, B_1], \quad (4)$$

$$\nabla_M B_2 = \partial_M B_2 + \frac{i}{4}\omega_M^{AB}\Gamma_{AB}B_2 - i[\mathcal{R}_M, B_2], \quad (5)$$

where the spin connection $\omega_M^{AB} = \partial_z A(z)(\delta_M^A\delta_z^B - \delta_z^A\delta_M^B)$, $\Gamma^{AB} = \frac{1}{2i}[\Gamma^A, \Gamma^B]$ is the Lorentz generator for spinor fields and $\mathcal{L}_M(\mathcal{R}_M)$ can be formulated by the gauged Maurer-Cartan 1-forms as

$$\mathcal{L}_M = i\xi^\dagger D_M\xi, \qquad \mathcal{R}_M = i\xi D_M\xi^\dagger \quad (6)$$

with the covariant derivatives defined by

$$D_M\xi = \partial_M\xi - i(A_L)_M\xi, \quad (7)$$

$$D_M\xi^\dagger = \partial_M\xi^\dagger - i(A_R)_M\xi^\dagger, \quad (8)$$

where A_L^M and A_R^M are the chiral gauge fields and $\xi = e^{iT_a\pi_a}$ is the element of the coset space G/H [$T_a = \lambda_a/2$ is the generator of the $SU(3)$ group with λ_a the Gell-Mann matrices, and π_a denotes the bulk pseudoscalar field]. We can see that $\mathcal{L}_M$ and $\mathcal{R}_M$ contain both the gauge fields and the pseudoscalar fields and are indeed the only terms containing these fields in the model, which is different from that in Ref. [22].

Note that in the nonlinear representation the bulk baryon fields B_j only transform under subgroup H as $B_j' = HB_jH^\dagger$. The element $\xi(\pi)$ of the coset space G/H transforms under the chiral gauge group G as $\xi' = L\xi(\pi)H^\dagger = H\xi(\pi)R^\dagger$ with $L \in SU(3)_L$, $R \in SU(3)_R$, and the chiral gauge fields A_L^M and A_R^M transform in the following way:

$$(A_L)'_M = L(A_L)_M L^\dagger + iL\partial_M L^\dagger, \quad (9)$$

$$(A_R)'_M = R(A_R)_M R^\dagger + iR\partial_M R^\dagger. \quad (10)$$

Then, the group transformations of $\mathcal{L}_M$ and $\mathcal{R}_M$ can be obtained as

$$\mathcal{L}'_M = H\mathcal{L}_M H^\dagger + iH\partial_M H^\dagger, \quad (11)$$

$$\mathcal{R}'_M = H\mathcal{R}_M H^\dagger + iH\partial_M H^\dagger. \quad (12)$$

The other sector of the bulk action which generates chiral symmetry breaking can be formulated as

$$S_{b2} = -\int d^5x\sqrt{g}[c_1\mathrm{Tr}(\bar{B}_1\{\chi_+, B_2\}) + c_2\mathrm{Tr}(\bar{B}_1[\chi_+, B_2]) + (c_2 - c_1)\mathrm{Tr}(\bar{B}_1 B_2)\mathrm{Tr}\chi_+ + \mathrm{h.c.}], \quad (13)$$

where h.c. denotes the Hermitian conjugate, χ_+ is related to the three-flavor generalization of the bulk scalar field X (see also Sec. IV) by

$$\chi_+ = \frac{1}{2}(\xi^\dagger X\xi^\dagger + \xi X^\dagger\xi), \quad (14)$$

and χ_+ transforms in the nonlinear representation as $\chi'_+ = H\chi_+H^\dagger$. The bulk scalar field X can be defined as

$$X \equiv \xi(X_0 + T_0S_0 + T_aS_a)\xi, \quad (15)$$

where S_0 and S_a represent the scalar singlet and octet, X_0 is the vacuum expectation value (VEV) of the scalar field X and is presumed to be of the form

$$X_0 = \frac{1}{2}\begin{pmatrix} v_u & & \\ & v_u & \\ & & v_s \end{pmatrix} \quad (16)$$

with $v_u = m_u z + \sigma_u z^3$ and $v_s = m_s z + \sigma_s z^3$, where m_u, σ_u denote the mass and chiral condensate of u, d quarks and m_s, σ_s denote the ones of s quark [7,8]. In terms of Eq. (15), χ_+ can be reduced to

$$\chi_+ = X_0 + T_0S_0 + T_aS_a, \quad (17)$$

which is only associated with the scalar mesons.

As in Ref. [22], the bulk baryon fields B_j can be written in the chiral form: $B_j = B_{jL} + B_{jR}$ with $B_{jL,R}$ decomposed by the Kaluza-Klein (KK) and Fourier decomposition as

$$B_{jL,R}(x,z) = \sum_n \int \frac{d^4p}{(2\pi)^4} e^{-ipx} f_{jL,R}^{(n)}(z)\psi_{L,R}^{(n)}(p), \quad (18)$$

where $f_{jL,R}$ are the Kaluza-Klein profiles of the bulk baryon fields, $\psi_{L,R}(p)$ are the 4D spinors satisfying $\gamma^5\psi_L(p) = \psi_L(p)$, $\gamma^5\psi_R(p) = -\psi_R(p)$, and $\not{p}\psi_{L,R}(p) = |p|\psi_{R,L}(p)$. Now, in terms of the profiles $f_{jL,R}$ (note that the superscript n has been neglected), the equation of motion (EOM) for the baryon octet can be derived from the variation of the bulk action as

$$\begin{pmatrix} \partial_z - \frac{9}{2z} & -\frac{v(z)}{2z} \\ -\frac{v(z)}{2z} & \partial_z + \frac{1}{2z} \end{pmatrix} \begin{pmatrix} f_{1L} \\ f_{2L} \end{pmatrix} = -|p| \begin{pmatrix} f_{1R} \\ f_{2R} \end{pmatrix},$$
$$\begin{pmatrix} \partial_z + \frac{1}{2z} & \frac{v(z)}{2z} \\ \frac{v(z)}{2z} & \partial_z - \frac{9}{2z} \end{pmatrix} \begin{pmatrix} f_{1R} \\ f_{2R} \end{pmatrix} = |p| \begin{pmatrix} f_{1L} \\ f_{2L} \end{pmatrix}, \qquad (19)$$

where $v(z)$ has the following form for different particles in the baryon octet:

$$v(z) = \begin{cases} (3c_2 - c_1)v_u(z) & \text{for } (p,n), \\ \left(2c_2 - \frac{4c_1}{3}\right)v_u(z) + \left(\frac{c_1}{3} + c_2\right)v_s(z) & \text{for } \Lambda, \\ 2c_2 v_u(z) + (c_2 - c_1)v_s(z) & \text{for } \Sigma\text{s}, \\ (c_2 - c_1)v_u(z) + 2c_2 v_s(z) & \text{for } \Xi\text{s}. \end{cases} \qquad (20)$$

From the parity transformations of the baryon fields, we get $f_{1L} = f_{2R}$ and $f_{1R} = -f_{2L}$ for even-parity states, while $f_{1L} = -f_{2R}$ and $f_{1R} = f_{2L}$ for odd-parity states [22,27]. As in Ref. [22], by solving Eqs. (19) with the boundary conditions $f_{1L}(\epsilon \to 0) = 0$ and $f_{1R}(z_m) = 0$, we obtain the mass spectra of the baryon octet and their excited states.

III. MASS SPECTRA OF THE BARYON OCTET AND THEIR EXCITED STATES

To get the baryon spectrum, we first fix the same parameters as the ones in Ref. [22] by fitting the masses of the nucleons (p, n) and their first excited state N(1440): $m_u = 0$, $\sigma_u = (198\ \text{MeV})^3$, $z_m = (205\ \text{MeV})^{-1}$, and $3c_2 - c_1 = 14.8$. Then, we take $m_s = 100$ MeV and $\sigma_s = \sigma_u = (198\ \text{MeV})^3$ and the last parameter c_2 fixed by the mass of Λ to be $c_2 = 4.52$. The experimental and numerical results for the baryon-octet spectrum are shown in Table I, where we can see that the model predictions for the masses of Σs and Ξs match with experimental data very well. It can be easily shown that the mass hierarchy of different particles in the baryon octet is exclusively due to the large mass of the strange quark. Table II lists the masses of the first excited baryon-octet states with both even and odd parities, which are also consistent with experiments, except for the first odd-parity state of Λ. It should be noted that the masses of the excited states of Ξs are not given due to the lack of experimental data. For the masses of higher excited states, this model cannot give consistent results with experiments, the reason for which might be attributed to the sharp IR cutoff of the fifth dimension of AdS_5, which can be remedied through a soft dilaton term [27,28].

TABLE I. Experimental value and model predictions for the mass spectrum of the baryon octet.

Baryon octet	p, n	Λ	Σs	Ξs
Exp.(MeV)	939[a]	1115[a]	1190	1320
Model (MeV)	939	1115	1184	1321

[a]Input value, and the experimental data are taken from Ref. [29].

TABLE II. Experimental value and model predictions for the mass spectra of the first excited even-parity and odd-parity states.

1st even	p, n	Λ	Σs	Ξs
Exp.(MeV)	1440[a]	1600	1660	—
Model (MeV)	1440	1548	1598	—
1st odd	p, n	Λ	Σs	Ξs
Exp.(MeV)	1535	1405	1620	—
Model (MeV)	1505	1588	1627	—

It should be remarked that the baryon masses in this model are acquired by a different way than that in Ref. [22], where the VEV of the scalar field X in a Yukawa coupling term breaks the chiral symmetry and gives the nucleon mass. In our model, with a nonlinear realization of chiral symmetry, the chiral symmetry of the baryon action *per se* has been broken from the beginning, which can be seen from the group transformations of the baryon fields B_j that keep the action invariant. However, it is still the VEV of the scalar field χ_+ that contributes to the mass spectrum of the baryon octet.

A unique feature of this bottom-up holographic framework is that one cannot separate the explicit chiral symmetry breaking terms from the ones generating spontaneous symmetry breaking in the bulk action as the quark mass and the chiral condensate are sewed into a single term, i.e., the VEV of the bulk scalar field X_0. In the three-flavor case, X_0 does not vanish due to the large mass of the s quark, even if there is no spontaneous chiral symmetry breaking. It leads us to accept the nonlinear realization of chiral symmetry. However, as the bulk baryon action does not realize the full chiral symmetry in the nonlinear representation, the baryon states cannot form the Wigner-Weyl modes of the chiral group, which might cause some vagueness for the explanation of the parity-doublet pattern of excited baryons. We will not go into details about the reason of the parity-doubling phenomenon, which is still inconclusive [23].

IV. PION-NUCLEON COUPLING

In the above, we construct a holographic model for the baryon octet by using a nonlinear realization of chiral symmetry, which has been proven to be more suitable for describing baryons with larger masses [25]. The explanation of the parity-doubling pattern in terms of two baryon fields $B_{1,2}$ with inverse chirality (parity) has been incorporated in this framework, though the bulk baryon action does not realize the full chiral symmetry. Another characteristic of this holographic model is that there are only derivative interactions of pions with baryons, which is also a desired feature in the 4D effective baryon models [25,26]. Next, we will derive the expressions for the pion-nucleon couplings and show directly the nature of derivative interactions.

As the pion wave function is necessary for the derivation of the pion-nucleon coupling, we first present the bulk action of the meson sector which has been studied in Refs. [7,8],

$$S_m = \int d^5x\sqrt{g}\mathrm{Tr}\left[|DX|^2 + 3|X|^2 - \frac{1}{4g_5^2}(F_L^2 + F_R^2)\right], \tag{21}$$

where the covariant derivative of the bulk scalar field X is $D^M X = \partial^M X - iA_L^M X + iXA_R^M$ and the chiral gauge field strength is $F_{L,R}^{MN} = \partial^M A_{L,R}^N - \partial^N A_{L,R}^M - i[A_{L,R}^M, A_{L,R}^N]$. The vector and axial-vector fields are defined as $V^M = \frac{1}{2}(A_L^M + A_R^M)$ and $A^M = \frac{1}{2}(A_L^M - A_R^M)$, respectively.

Note that the scalar part of the meson action (21) can also be recast into a nonlinear form by the Eq .(15), and the results do not depend on which forms we used. Below, we will work in two gauges, as has been used in Refs. [7] and [8], respectively. The results in both gauges will be shown to be equivalent to each other, from which the gauge-independence manifests itself obviously.

Let us first work in the $A_5 = 0$ gauge accepted in Ref. [8]. The EOM of the pseudoscalar meson in this gauge can be derived as

$$\partial_z\left(\frac{\partial_z f_\varphi}{z}\right) + \frac{g_5^2 v_u^2}{z^3}(f_\pi - f_\varphi) = 0, \tag{22}$$

$$-q^2\partial_z f_\varphi + \frac{g_5^2 v_u^2}{z^2}\partial_z f_\pi = 0, \tag{23}$$

where $f_\pi(z)$ comes from the KK decomposition of the bulk pion field and $f_\varphi(z)$ comes from that of the radial component of axial-vector field ($A_\mu^a = A_{\mu\perp}^a + \partial_\mu \varphi^a$). As in the calculation of baryon spectrum, we only consider the $m_u = 0$ case in which the pion as Nambu-Goldstone boson has zero mass ($q^2 = 0$), so we have $f_\pi'(z) = 0$, which also indicates, with the normalization condition $f_\pi(z \to 0) = 0$, that $f_\pi(z) = 0$. Then Eq. (22) can be reduced to the following form:

$$z\partial_z\left(\frac{\partial_z f_\varphi}{z}\right) - \frac{g_5^2 v_u^2}{z^2} f_\varphi = 0. \tag{24}$$

The pion-nucleon couplings are supplied by the covariant kinetic term of the bulk baryon action (3), which is the unique term containing the pseudoscalar and gauge fields, as noted above. In the $A_5 = 0$ gauge, the interaction terms (Lagrangian density) of pion and nucleons can be extracted from the action (3) as

$$\begin{aligned}\mathcal{L}_{\pi NN}^{(1)} &= z[(f_\pi - f_\varphi)(|f_{1\mathrm{L}}|^2 - |f_{2\mathrm{L}}|^2)\bar\psi\gamma^5 \partial\!\!\!/ \pi^a t^a \psi \\ &\quad + \partial_z f_\pi (f_{1\mathrm{L}}^* f_{1\mathrm{R}} - f_{2\mathrm{L}}^* f_{2\mathrm{R}}) i\bar\psi\gamma^5 \pi^a t^a \psi] \\ &= -z f_\varphi (|f_{1\mathrm{L}}|^2 - |f_{2\mathrm{L}}|^2)\bar\psi\gamma^5 \partial\!\!\!/ \pi^a t^a \psi,\end{aligned} \tag{25}$$

where the superscript $*$ denotes complex conjugate, $t^a = \frac{\sigma^a}{2}$ with σ^a the Pauli matrices, and $\psi(x)$ denotes the 4D isodoublet of nucleons. The Lagrangian density (25) shows obviously the derivative coupling of pion to nucleons, which is a characteristic feature of the nonlinear representation. Note that it is different from the holographic baryon model proposed in Ref. [22], where a Yukawa coupling term was introduced for chiral symmetry breaking, which mimics the structure of the linear sigma model.

For comparison, we take another gauge-fixing scheme used in Ref. [7] to work out the coupling of pion to nucleons and show that the result is equal to the one in the $A_5 = 0$ gauge. The pion meson in this gauge is associated with the axial-vector and pseudoscalar fields which can be extracted (up to quadratic terms) from the action (21) as follows:

$$\begin{aligned}S_{\pi A} = \int d^5x \Big[&-\frac{1}{2g_5^2 z}(\partial_\mu A_\nu^a \partial^\mu A^{a\nu} - \partial_\mu A_\nu^a \partial^\nu A^{a\mu} \\ &- (\partial_\mu A_5^a - \partial_z A_\mu^a)^2) + \frac{v_u^2}{2z^3}((\partial_\mu \pi^a - A_\mu^a)^2 \\ &- (\partial_z \pi^a - A_5^a)^2)\Big].\end{aligned} \tag{26}$$

To eliminate the mixing terms between A_5^a (π^a) and A_μ^a, we add the following gauge-fixing term:

$$S_{\mathrm{g.f.}} = \int \frac{-d^5x}{2\xi_A g_5^2 z}\left[\partial^\mu A_\mu^a - \xi_A z\partial_z\left(\frac{A_5^a}{z}\right) + \frac{\xi_A g_5^2 v_u^2}{z^2}\pi^a\right]^2. \tag{27}$$

In the unitary gauge, we take $\xi_A \to \infty$ to decouple the term $z\partial_z(\frac{A_5^a}{z}) - \frac{g_5^2 v_u^2}{z^2}\pi^a$ from the A_μ^a terms. Note that the orthogonal combination of A_5^a and π^a remains massless if

$$z\partial_z\left(\frac{A_5^a}{z}\right) - \frac{g_5^2 v_u^2}{z^2}\pi^a = 0. \tag{28}$$

With the ansatz $A_5^a(x,z) = f_0(z)\pi^a(x)$, Eq. (28) becomes

$$z\partial_z\left(\frac{f_0(z)}{z}\right) - \frac{g_5^2 v_u^2}{z^2} f_\pi(z) = 0. \tag{29}$$

After fixing the gauge, the pseudoscalar part of the action (26) can be written as

$$S_\pi = \int d^5x \left[\frac{1}{2g_5^2 z} \partial_\mu A_5^a \partial^\mu A_5^a + \frac{v_u^2}{2z^3} \partial_\mu \pi^a \partial^\mu \pi^a - \frac{v_u^2}{2z^3} (\partial_z \pi^a - A_5^a)^2 \right], \quad (30)$$

from which one can see that the condition for the pion to be massless is $\partial_z \pi^a - A_5^a = 0$, which implies $f_0(z) = f'_\pi(z)$. In this gauge, we can also derive the coupling of pion to nucleons as

$$\mathcal{L}_{\pi NN}^{(2)} = z[f_\pi(|f_{1L}|^2 - |f_{2L}|^2)\bar{\psi}\gamma^5 \partial\!\!\!/ \pi^a t^a \psi + (\partial_z f_\pi - f_0)(f_{1L}^* f_{1R} - f_{2L}^* f_{2R}) i\bar{\psi}\gamma^5 \pi^a t^a \psi] = z f_\pi (|f_{1L}|^2 - |f_{2L}|^2)\bar{\psi}\gamma^5 \partial\!\!\!/ \pi^a t^a \psi. \quad (31)$$

To prove the equivalence of $\mathcal{L}_{\pi NN}$ in Eqs. (25) and (31), one just notes that Eq. (29) is the same as Eq. (24) in view of the relation $f_0(z) = f'_\pi(z)$ derived in the unitary gauge. It indicates that the couplings of the pion with nucleons are equivalent to each other in the two gauges. The gradient coupling constant can be defined as

$$\frac{g_A}{F_\pi} \equiv \int_0^{z_m} \frac{dz}{z^4} f_\pi(z)(|f_{2L}(z)|^2 - |f_{1L}(z)|^2), \quad (32)$$

where F_π denotes the pion decay constant, $f_{1L}(z)$ and $f_{2L}(z)$ are normalized by the condition

$$\int_0^{z_m} \frac{dz}{z^4} (|f_{1L}(z)|^2 + |f_{2L}(z)|^2) = 1, \quad (33)$$

and $f_\pi(z)$ is solved from Eq. (29) with the normalization of $f_0(z)$:

$$\int_0^{z_m} dz \left[\frac{1}{g_5^2 z} f_0^2 + \frac{z^3}{g_5^4 v^2} \left(\partial_z \left(\frac{f_0}{z} \right) \right)^2 \right] = 1. \quad (34)$$

The pion decay constant F_π can be computed from the meson action (21) as $F_\pi^2 = -(g_5^2 z)^{-1} \partial_z A(0,z)|_{z=\epsilon}$ with $A(q,z)$ the bulk-to-boundary propagator of the axial-vector field [8]. Using the parameters fixed by the baryon spectrum, we get $F_\pi \simeq 54.5$ MeV and $g_A \simeq 0.33$, which are too small compared with the experimental value: $F_\pi \simeq 92.4$ MeV and $g_A \simeq 1.27$ [29].

The discrepancy between the holographic model for spin-$\frac{1}{2}$ baryons and the one for mesons has been revealed in Ref. [22], where the masses of $N(1440)$ and $N(1535)$ cannot be obtained from the fitting of the meson sectors [7], which is ascribed to the nonvanishing anomalous dimension of baryons in Ref. [22]. Another important reason for this inconsistency, as has been noted above, is that the sharp IR cutoff of the AdS_5 metric causes no linear confinement which is manifest in the hadron spectrum [27,28].

The coupling of the pion to nucleons may also receive contributions from nonminimal gauge interactions, such as the magnetic gauge coupling term [30]

$$\mathcal{L} = i\kappa[\bar{B}_1 \Gamma^{MN} (\mathcal{F}_L)_{MN} B_1 - \bar{B}_2 \Gamma^{MN} (\mathcal{F}_R)_{MN} B_2], \quad (35)$$

where the transformed gauge field strengths are $\mathcal{F}_L^{MN} = \xi^\dagger F_L^{MN} \xi$ and $\mathcal{F}_R^{MN} = \xi F_R^{MN} \xi^\dagger$. Note that this term also only generates gradient coupling of the pion to nucleons.

V. SUMMARY AND CONCLUSION

In this paper, we propose a holographic model for the baryon octet with a nonlinear realization of chiral symmetry on the basis of the AdS/QCD model for mesons [7,8]. The mass spectra of the baryon octet and their low-lying excited states have been calculated and show good agreement with experiments, especially for the ground states. However, the model cannot give consistent results with experiments for highly excited baryon states, which might be attributed to the sharp IR cutoff of the AdS_5 metric [27,28]. We have also shown that the model only contains gradient coupling of pion to nucleons, as in the 4D effective baryon models with a nonlinear realization of chiral symmetry [25]. However, the derivative coupling constant is much smaller than that obtained from experimental measurements, which might be rooted in the ignorance of other gauge interaction terms contributing to the coupling of the pion to nucleons, as has been noted above.

Note that two bulk baryon fields are introduced in this holographic model, although they transform in the same way under the gauge group H. This is an unusual point of the model which is different from the previous studies [22]. As the bulk action does not realize the full chiral symmetry, the parity-doubling property of excited baryon states cannot be explained with a Wigner-Weyl realization of the chiral group, so other reasons must be called for [23]. On the other hand, this holographic baryon model with nonlinear realization can be easily transformed into a linear representation with the full chiral symmetry restored and the baryon spectrum unchanged. However, it will lead to lengthier action terms, and the derivative coupling of pion with baryons will be lost.

Anyway, a holographic model for the baryon octet has been constructed, which incorporates the parity-doublet pattern of excited baryon states in an unusual way and leads to unique derivative coupling of the pion to nucleons. Nevertheless, there is much room to improve the model, and much more low-energy hadron physics related to the baryon octet can be studied in this holographic framework. How to reconcile the physics of the baryon sector and the meson sector in a desired way still deserves further studies.

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APPENDIX REVIEW OF THE TWO-FLAVOR HOLOGRAPHIC BARYON MODEL

Here, we give a brief review of the holographic model for spin-$\frac{1}{2}$ baryons first proposed in Ref. [22], where the authors introduced two bulk baryon fields $N_{1,2}$ which correspond to the 4D chiral baryon operators $\mathcal{O}_{L,R}$, respectively. The bulk action for the baryons can be written as

$$S_N = \int d^5x\sqrt{g}\sum_{j=1}^{2}\left[\frac{i}{2}\bar{N}_j e_A^M \Gamma^A \nabla_M N_j - \frac{i}{2}(\nabla_M^\dagger \bar{N}_j) e_A^M \Gamma^A N_j - m_5 \bar{N}_j N_j\right] \quad \text{(A1)}$$

with a Yukawa coupling term,

$$S_Y = -\int d^5x\sqrt{g} g_Y(\bar{N}_1 X^\dagger N_2 + \bar{N}_2 X N_1), \quad \text{(A2)}$$

where the five-dimensional Dirac matrices $\Gamma^A = (\gamma^\mu, -i\gamma^5)$ satisfy $\{\Gamma^A, \Gamma^B\} = 2\eta^{AB}$, the vielbein e_A^M satisfies $g_{MN} = e_M^A e_N^B \eta_{AB}$, and the Lorentz (gauge) covariant derivative has the form

TABLE III. The results of baryon spectrum taken from Ref. [22].

z_m(GeV^{-1})	g_Y	(p,n)(GeV)	$N(1440)$	$N(1535)$	3rd	4th
$(0.33)^{-1\,a}$	8.67	0.94[a]	2.14	2.24	3.25	3.30
$(0.205)^{-1}$	14.4	0.94[a]	1.44[a]	1.50	2.08	2.12

[a]Input value, and $\sigma = \frac{\sqrt{2}\xi}{g_5 z_m^3}$ with $3.4 \le \xi \le 4$.

$$\nabla_M = \partial_M + \frac{i}{4}\omega_M^{AB}\Gamma_{AB} - i(A_{L,R}^a)_M t^a \quad \text{(A3)}$$

with $\Gamma^{AB} = \frac{1}{2i}[\Gamma^A, \Gamma^B]$ and $\omega_M^{AB} = -\frac{1}{z}(\delta_M^A \delta_z^B - \delta_z^A \delta_M^B)$. Note that we also have $m_5 = \frac{5}{2}$ for N_1 and $m_5 = -\frac{5}{2}$ for N_2.

As only the two-flavor case with vanishing current quark mass was considered in Ref. [22], the VEV of the bulk scalar field takes the form $X = \frac{1}{2}\sigma z^3$. The EOM for the baryon doublet can be derived from the above action and has the same form as Eq. (19) with $v(z) = g_Y \sigma z^3$. The baryon spectrum can be obtained by solving the EOM with the same boundary conditions as before. The results of Ref. [22] are listed in Table III.

The second-row results are obtained by inputting the parameters of previous studies of the meson sector [7,8], while the third row shows the best-fitting results. Obviously, this model can realize the parity doubling pattern naturally, although it cannot give a good baryon spectrum which is consistent with experiments, especially for the highly excited baryon resonances.

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