PHYSICAL REVIEW D **83,** 111504(R) (2011)

# Composite and elementary natures of $a_1(1260)$ meson

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(Received 19 January 2011; published 20 June 2011)

We develop a practical method to analyze the mixing structure of hadrons consisting of two components of quark composite and hadronic composite. As an example, we investigate the properties of the axial vector meson $a_1(1260)$ and discuss its mixing properties quantitatively. We also make reference to the large $N_c$ procedure and its limitation for the classification of such a mixed state.

DOI: 10.1103/PhysRevD.83.111504 PACS numbers: 14.40.−n, 13.75.Lb

Hadrons, interacting with the strong force, are composite objects of quarks and gluons. One of recent interests in the hadron structure is whether hadrons are made up of quarks and gluons confined in a single-particle potential as described in the conventional quark model, or rather develop subcomponents of quark-clusters inside hadrons. A typical example of the latter is the deuteron, which is composed of a proton and a neutron, not of six quarks in a single confining potential [1]. It has been also suggested that some hadronic resonances could have substantially large components of hadronic composites [2,3]. Such recursive (nesting-box) structures are also seen in nuclear physics. For instance, the first $0^+$ excited state of $^{12}$C is described as a three-$\alpha$ cluster state [4].

In general, it is not easy to clearly identify the quark-cluster components, because the strong interaction scales for quarks and hadrons are not well separated. To simplify the situation, it would be a good starting point to set up a model space of a state described by a single-particle potential and one of several quark-clusters. The former may be identified with the "elementary" component[1] while the latter with the (hadronic) composite, the situation which was studied by Weinberg [1]. If hadronic resonant states are unavoidably mixtures of hadronic and quark-composites, an important issue is to clarify how these components are mixed in a hadron.

One good example to study the two features and their mixing is provided by the low-lying axial vector meson $a_1(1260)$. The $a_1$ meson is a candidate of the chiral partner of the $\rho$ meson [5–7] described as a $q\bar{q}$-composite, for example, in the Nambu-Jona-Lasinio model [8,9] and in the Lattice calculation [10]. It can also appear as a gauge boson of the hidden local symmetry [11], which is recently reconciled with the five-dimensional gauge field of the holographic QCD [12,13]. On the other hand, in coupled-channel approaches based on the chiral effective theory [14,15], the $a_1$ meson has been described as a dynamically generated resonance in the $\pi\rho$ scattering without introducing its explicit pole term. The $a_1$ nature has been studied by calculating physical observables such as the radiative decay width [16] or the $\tau$ decay spectrum into three pions [17,18]. Yet, the internal structure of the $a_1$ meson is not well understood.

In this paper, we focus on hadron structure having two components of quark-composite (we refer to it as the elementary component) and hadronic composite. We propose a method to disentangle their mixture appearing in the physically observed state by taking the $a_1$ meson as an example. The method provided here can be generally applied to other mixed systems if the interaction is given between the elementary component and possible constituents making the composite state. Our ingredients for the study of $a_1$ are therefore $\pi$ and $\rho$ mesons[2], which have potential to generate the *composite* $a_1$ meson with a suitable $\pi\rho$ interaction, and the *elementary* component of the $a_1$ meson (dominated by $q\bar{q}$), which couples to the $\pi\rho$ pair by a three-point interaction. We first solve the $\pi\rho$ scattering amplitude to find the poles corresponding to the physical $a_1$, and then develop a method to clarify the mixing nature of the two components by introducing appropriate bases for pure composite and elementary components. As an extension, we apply our method to the study of the large $N_c$ behavior of the $a_1$ state. The large $N_c$ limit is usually considered to give a guiding principle for the classification of state, in which the mass of the $q\bar{q}$ bound state scales as $\mathcal{O}(1)$ in powers of $N_c$ while that of a meson-meson molecule scales with higher order of $N_c$ [19,20]. Reference [21], however, has brought caution for the use of large $N_c$ argument for the dynamically generated scalar resonances. We discuss the validity of the classification for the mixed states of hadronic composite and elementary components.

Let us start with the composite $a_1$ meson, which is dynamically generated in the $s$-wave $\pi\rho$ scattering through the nonperturbative dynamics. The scattering amplitude $t$ satisfies the Bethe-Salpeter equation,

[1]The elementary component was referred to as the CDD pole or genuine quark state in the literature [3].

[2]We regard the $\rho$ meson as stable particles in the present model setting.

 

$t = v + vGt$, where $v$ is a four-point $\pi\rho$ interaction and $G$ the $\pi\rho$ two-body propagator and its formal solution is given by

$$t = \frac{v}{1 - vG}. \tag{1}$$

If the potential $v$ is sufficiently attractive, the amplitude develops a pole corresponding to a composite bound or resonant state of the scattering system at the energy satisfying $1 - vG = 0$. In the present case of $a_1$, the potential $v$ can be obtained from the $s$-wave projection of the Weinberg-Tomozawa interaction [14], and the pole in the $\pi\rho$ scattering amplitude can appear above the $\pi\rho$ threshold as a resonance in the second Riemann sheet, which is a consequence of the energy-dependent interaction. This pole corresponds to the $\pi\rho$-composite $a_1$ meson [14] without $q\bar{q}$ quark core [3,22].

Because the elementary $a_1$ meson has a coupling to the $\pi$ and $\rho$ mesons, it also contributes to the $\pi\rho$ scattering amplitude in the form of an effective $\pi\rho$ interaction going through the elementary $a_1$ pole:

$$v_{a_1} = g \frac{1}{s - m_{a_1}^2 + i\epsilon} g, \tag{2}$$

where $g$ is the coupling to $\pi\rho$ that can depend on $s$, and $m_{a_1}$ the bare mass of the elementary $a_1$ meson. The full scattering amplitude $T$ having both the $\pi\rho$ four-point interaction $v$ and the $a_1$ pole term $v_{a_1}$ is then written by

$$T = \frac{v + v_{a_1}}{1 - (v + v_{a_1})G}. \tag{3}$$

This amplitude generates poles corresponding to physical resonant states of the problem. They are expressed as a superposition of the basis states associated with the two poles of composite $a_1$ in Eq. (1) and elementary $a_1$ in Eq. (2), respectively.

Now let us study the mixing nature of the physical states. To this end, we first express equivalently the amplitude $t$ in Eq. (1) as

$$t \equiv g_R(s) \frac{1}{s - s_p} g_R(s), \tag{4}$$

where $s_p$ is the pole position of the amplitude $t$ in Eq. (1). In this form, we can interpret $(s - s_p)^{-1}$ as the one-particle propagator of the composite $a_1$ meson as shown in Fig. 1 by taking an analogy with the conventional discussion of the bound-state problem [22]. The vertex function $g_R(s)$ that is defined so as to reproduce Eq. (1) exactly is interpreted as the effective coupling of the composite $a_1$ to $\pi\rho$. This interpretation works well in the neighborhood of the pole, $s \sim s_p$. As $s$ is further away from $s_p$, $g_R(s)$ receives more contributions from the nonresonant background [23].

Having the form of Eq. (4), we now rewrite the scattering amplitude $T$ in Eq. (3) as

$$T = (g_R, g) \frac{1}{\hat{D}_0^{-1} - \hat{\Sigma}} \begin{pmatrix} g_R \\ g \end{pmatrix}, \tag{5}$$

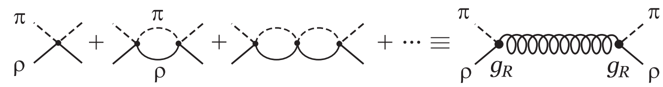


FIG. 1. Infinite set of diagrams that contributes to the meson-meson scattering amplitude in Eq. (1) and the definition of the propagator of the composite $a_1$ meson in Eq. (4).

where

$$\hat{D}_0^{-1} = \begin{pmatrix} s - s_p & \\ & s - m_{a_1}^2 \end{pmatrix}, \qquad \hat{\Sigma} = \begin{pmatrix} & g_R G g \\ g G g_R & g G g \end{pmatrix}. \tag{6}$$

The diagonal elements of the matrix $\hat{D}_0$ are the free propagators of the two $a_1$'s, one for the composite and the other for the elementary ones having the proper normalization, and the matrix $\hat{\Sigma}$ expresses the self-energy and interactions for these modes. One can prove that Eq. (5) is identical with Eq. (3) after some algebra.

We emphasize that the expression of Eq. (5) makes it possible to analyze the mixing nature of the physical $a_1$ in terms of the original two bases. Having the amplitude in the form of Eq. (5), the matrix $\hat{D} \equiv (\hat{D}_0^{-1} - \hat{\Sigma})^{-1}$ is identified with the propagators of the physical states represented by the bases of the elementary and composite $a_1$'s. The diagonal elements $D^{ii}$ indicate the dressed propagators of the composite and the elementary $a_1$'s as shown in Fig. 2, which express the $a_1$ mesons acquiring the quantum effects: e.g., the elementary $a_1$ introduced without width in Eq. (2) now has a width owing to its decay into $\pi\rho$ system. The important features of the propagators are that they have poles exactly at the same positions as the full amplitude $T$ in Eq. (3), and the residues of the diagonal elements $D^{ii} (i = 1, 2)$ defined by

$$D^{ii} = \frac{z_a^{ii}}{s - M_a^2} + \frac{z_b^{ii}}{s - M_b^2} + (\text{regular term}) \tag{7}$$

have the meaning of the wave function renormalization and then carry the information on the mixing rate of the physical resonant states. For instance, the residue $z_a^{11}$ means the probability of finding the original composite $a_1$ component in the resulting state having the mass $M_a$.

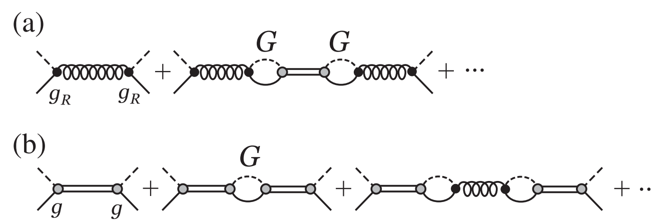


FIG. 2. Diagrams that contribute to the full amplitude in Eq. (5) going through (a) the component of the dressed propagator $D^{11}$ and (b) that of $D^{22}$. The dashed and solid internal lines indicate the $\pi$ and $\rho$ propagators, while the curved and double lines are those of the composite and elementary $a_1$'s.

The above discussions have close analogy to a two-level problem, which is well realized when the energy $s$ is not very far away from the original two poles $s_p$ and $m_{a_1}^2$. As $s$ is getting away from them, the orthogonality of the physical two-resonant states is lost and, furthermore, the residues can exceed unity. These problems arise owing to the energy dependence of $\hat{\hat{\Sigma}}$, especially those of $g_R$ and $g$ in $\hat{\hat{\Sigma}}$. We note that the energy dependence of $g_R$ is unavoidable for the resonant states generated dynamically.

So far, we have given a general framework to investigate the mixing properties of hadrons consisting of multiple components having different origins. To perform a concrete calculation, we shall employ a suitable model for $a_1$. Here we take only the single $\pi\rho$ channel into account for the composite $a_1$ because other channels have been found not important [14], but we can easily extend our framework to coupled-channel cases. Our ingredients for the study of the $a_1$ system are, therefore, the $\pi$ and $\rho$ and (elementary) $a_1$ mesons, for which a suitable model is given.

In this paper, as for interaction Lagrangians, we adopt the chiral Lagrangians induced by the holographic QCD approach as the Sakai-Sugimoto model [12] that is constructed in the strong-coupling limit of large-$N_c$ QCD. The relevant Lagrangians [12] are given by

$$\mathcal{L}_{\mathrm{WT}} = \frac{1}{f_\pi^2}\,\mathrm{tr}([\rho^\mu, \partial^\nu \rho_\mu][\pi, \partial_\nu \pi]), \tag{8}$$

$$\begin{aligned}\mathcal{L}_{a_1\pi\rho} = -ig_{a_1\pi\rho}\frac{4}{f_\pi}\{&\mathrm{tr}((\partial_\mu a_{1\nu} - \partial_\nu a_{1\mu})[\partial^\mu \pi, \rho^\nu]) \\ &+ \mathrm{tr}((\partial_\mu \rho_\nu - \partial_\nu \rho_\mu)[\partial^\mu \pi, a_1^\nu])\},\end{aligned} \tag{9}$$

where $\rho^\mu \equiv \vec{\rho}^\mu \cdot \frac{\vec{\tau}}{2}$, $\pi \equiv \vec{\pi} \cdot \frac{\vec{\tau}}{2}$, and $a_1^\mu \equiv \vec{a}_1^\mu \cdot \frac{\vec{\tau}}{2}$. The first Lagrangian gives the four-point Weinberg-Tomozawa (WT) interaction and the second gives the three-point vertex of $a_1\pi\rho$.

The advantage of using the concept of the holographic QCD approach is that the large-$N_c$ condition ensures that the $a_1$ field in the Lagrangian does not contain hadronic composite components, and hence we can avoid the double-counting in the analysis of the mixing nature. The model contains two inputs, $f_\pi = 92.4$ MeV and $m_\rho = 776$ MeV, giving the mass of the elementary $a_1$ meson $m_{a_1} = 1189$ MeV and the $a_1\pi\rho$ coupling constant $g_{a_1\pi\rho} = 0.26$. In the present study, we consider the $a_1$ field appearing in the Lagrangian with these physical constants to be the quark composite that survives in the large $N_c$ limit. For the mass of the pion, we employ the physical value $m_\pi = 138$ MeV that is isospin-averaged. These interactions in Eqs. (8) and (9) are essentially the same as those of the hidden-local symmetry [11], except for the actual values of $m_{a_1}$ and $g_{a_1\pi\rho}$.

Using these interactions, we take the WT potential $V_{\mathrm{WT}}$ [14] and the $a_1$ pole term $V_{a_1}$ as

$$V_{\mathrm{WT}} = \frac{\epsilon\cdot\epsilon'}{4f_\pi^2}\left\{3s - 2(m_\rho^2 + m_\pi^2) - \frac{1}{s}(m_\rho^2 - m_\pi^2)^2\right\}, \tag{10}$$

$$V_{a_1} = -\frac{8}{f_\pi^2}g_{a_1\pi\rho}^2\frac{\epsilon\cdot\epsilon'}{s - m_{a_1}^2 + i\epsilon}(s - m_\rho^2)^2, \tag{11}$$

after the $s$-wave projection with on-shell energies for external $\pi$ and $\rho$ [14]. In actual calculations, we should treat the polarization vectors $\epsilon$ and $\epsilon'$ for the $\rho$ meson appropriately. As reported in detail in Ref. [14], by substituting the coefficient $v_{\mathrm{WT}}(v_{a_1})$ defined by $V_{\mathrm{WT}}(V_{a_1}) \equiv -\epsilon\cdot\epsilon' v_{\mathrm{WT}}(v_{a_1})$ for the potentials in Eq. (3), we can obtain the scattering amplitude for the transverse polarization mode, in which we can find poles dynamically generated.

In solving the scattering equation, we need to calculate the propagator function $G$. Because the potentials in Eqs. (10) and (11) are separable, the formal solution of Eq. (1) becomes algebraic, and the function $G$ is then given by

$$G(\sqrt{s}) = i\int\frac{d^4q}{(2\pi)^4}\frac{1}{(P-q)^2 - m_\pi^2 + i\epsilon}\frac{1}{q^2 - m_\rho^2 + i\epsilon}, \tag{12}$$

where $P$ is the total four-momentum as $P^2 = s$. The integral (12) diverges, and hence, we regularize it by the dimensional regularization, and for the finite part we introduce a subtraction constant $a(\mu)$ as

$$\begin{aligned}G(\sqrt{s}) = \frac{1}{16\pi^2}\Big\{&a(\mu) + \ln\frac{m_\rho^2}{\mu^2} + \frac{m_\pi^2 - m_\rho^2 + s}{2s}\ln\frac{m_\pi^2}{m_\rho^2} \\ &+ \frac{q'}{\sqrt{s}}[\ln(s - (m_\rho^2 - m_\pi^2) + 2q'\sqrt{s}) \\ &+ \ln(s + (m_\rho^2 - m_\pi^2) + 2q'\sqrt{s}) \\ &- \ln(s - (m_\rho^2 - m_\pi^2) - 2q'\sqrt{s}) \\ &- \ln(s + (m_\rho^2 - m_\pi^2) - 2q'\sqrt{s}) - 2\pi i]\Big\},\end{aligned} \tag{13}$$

where $\mu$ is the scale parameter that is set to be 900 MeV in this paper and $q' = \lambda^{1/2}(s, m_\rho^2, m_\pi^2)/2\sqrt{s}$. A crucial point here is that the constant $a(\mu)$ is chosen to be a *natural* value [3],which ensures that the resulting resonant states, if they exist, are interpreted as a purely hadronic composite. Following the prescription in Ref. [3], we choose $a(\mu) = -0.2$ at the renormalization scale $\mu$, which satisfies the matching condition $\mathrm{Re}\,G(\sqrt{s} = m_\rho) = 0$. A choice of the subtraction constant $a(\mu)$ different from the natural value is equivalent to the introduction of the CDD (or elementary) pole that is not included in the model space of the scattering problem [3].

The polarization vector of the intermediate $\rho$ meson is treated in the same way as reported in Ref. [14].

Now we can evaluate the scattering amplitude (3) or (5) numerically. We find two poles at (a) $\sqrt{s} = 1033 - 107i$ MeV and at (b) $1728 - 313i$ MeV, corresponding to the physical states in the present model. These pole positions are significantly different from those of the two basis states, $\sqrt{s_p} = 1012 - 221i$ MeV and

$m_{a_1} = 1189$ MeV, because of the mixing effect. For further investigation, we vary the coupling strength of $a_1\pi\rho$ by introducing a parameter $x(0 \le x \le 1)$ and $g_{a_1\pi\rho} \to x g_{a_1\pi\rho}$, which controls the mixing strength.

In Fig. 3(a), we show the resulting pole-flow in the complex-energy plane by changing the mixing parameter $x$. At $x = 0$, the poles corresponding to the basis states of the composite and elementary $a_1$'s are found at the positions indicated by open square and open circle in Fig. 3(a), respectively. When the mixing is turned on, the pole starting from the composite $a_1$ (we refer to it as "pole-a") approaches the real axis, ending at $1033 - 107i$ MeV when $x = 1$ (solid square), while that from the elementary $a_1$ pole ("pole-b") goes far from the real axis and reaches $1728 - 313i$ MeV when $x = 1$ (solid circle).

In Fig. 3(b), we show the squared amplitude $|T|^2$ in Eq. (3) (or (5)) at $x = 1$. As shown in the figure, a peak structure is dominated by the pole-a, while a signal of the pole-b cannot be seen clearly because of its huge width ($\sim 600$ MeV). Therefore, the pole expected to be observed in experiments is the pole-a located at lower energy position [solid square in Fig. 3(a)] that comes from the composite $a_1$ pole. Indeed, the contribution of the pole-b is found to interfere destructively with the tail of the pole-a (around $\mathrm{Re}\sqrt{s} \sim 1.7$–$1.8$ MeV) as shown in Fig. 3(b). Our present model setting, however, is rather simple and may not be applicable to such a higher energy region where

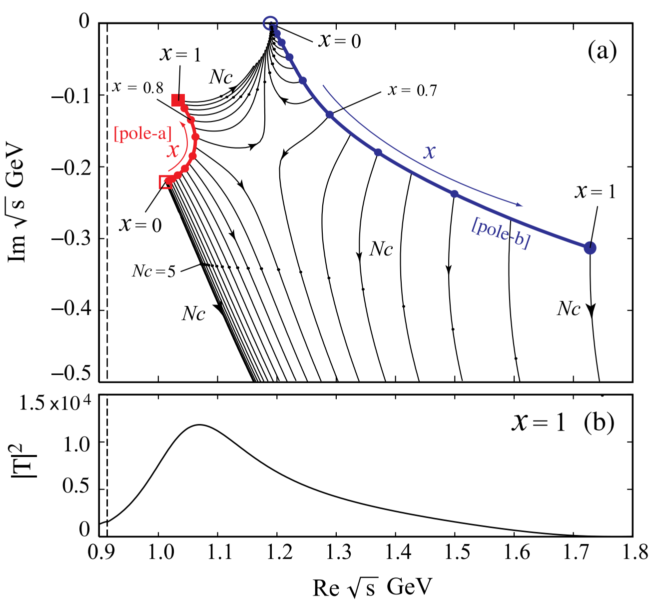


FIG. 3 (color online). (a) Trajectories of the poles in the full scattering amplitude in Eq. (3) by changing the mixing parameter $x$ (thick lines). The open square indicates the pole position of the composite $a_1$ and open circle indicates the elementary $a_1$ pole ($x = 0$). The other end points of solid circle and square correspond to the physical points ($x = 1$). Thin lines represent the pole-flows as $N_c$ is increased from $N_c = 3$ for fixed $x$ with small dots at $N_c = 5$. The vertical dashed line denotes the $\pi\rho$ threshold energy. (b) Squared amplitude $|T|^2$ of $\pi\rho \to \pi\rho$ process on the real energy axis for the mixing parameter $x = 1$.

contributions from higher coupled-channels neglected here should be important. A detailed investigation for an evidence of the second state of the $a_1$ meson, namely, the pole-b, is an interesting future work.

To study the mixing properties more quantitatively, we show the absolute values of the residues as functions of the mixing parameter $x$ in Fig. 4 (left panel). We can verify that, before turning on the mixing, the pole-a is purely the composite $a_1(z_a^{11} = 1$ and $z_a^{22} = 0$ at $x = 0)$ and the pole-b is the elementary $a_1$ as we intended. One of the most important messages can be read from the magnitude of $z_a^{11}$ and $z_a^{22}$ at $x = 1$, which are the residues of the possibly observed $a_1$ state. While we should carefully discuss the meaning of the residues of complex-value for resonant states, we can say that the pole-a at $x = 1$, although its location is close to the composite $a_1$ pole, has a component of the elementary $a_1$ meson *comparable* to that of the composite $a_1$. This conclusion can be drawn only after we look into the wave function by the residues $z$.

As for the pole-b, we find that the magnitude of the residues $z_b^{22}$ and $z_b^{11}$ interchange at around $x \sim 0.7$, meaning that the nature of the resonant state of pole-b changes from the elementary particle-like structure to the composite one. For larger $x(\gtrsim 0.8)$, $z_b^{11}$ and $z_b^{22}$ become larger than unity because the energy dependence of the potentials become stronger at the higher energy region where the pole-b is located. Once again, in such a higher energy region ($\mathrm{Re}\sqrt{s} \gtrsim 1.5$ GeV), our present model may not be applicable.

Next, we test the large $N_c$ dependence of the pole positions according to the scaling law of the pion-decay constant $f_\pi$ as $f_\pi \to f_\pi\sqrt{N_c/3}$. In Fig. 3(a), we also show the trajectories of the pole positions by changing the $N_c$ value for fixed mixing strength. At $x = 0$, as reported in Ref. [24], we see the composite $a_1$ pole tends to have a heavier mass and a wider width as $N_c$ is increased. For small mixing parameter $x \lesssim 0.6$, we find a similar behavior for the pole-a while the pole-b goes back to the elementary $a_1$ position. At $x \sim 0.7$, we find that the two $N_c$-trajectories flip simultaneously. For $x \gtrsim 0.7$, the pole-b goes away from the real axis as the composite $a_1$ does, while the pole-a approaches the elementary pole as $N_c$ is increased. In this way, the result of the large $N_c$ classification depends

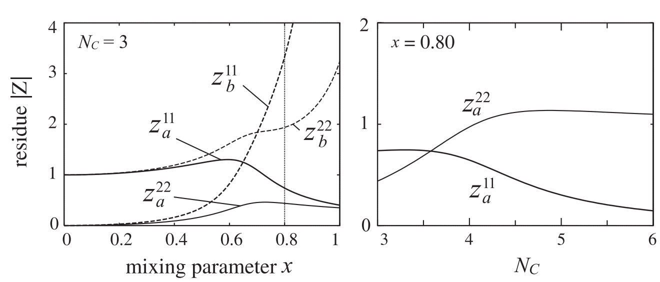


FIG. 4. Absolute value of the residues defined in Eq. (7). The left panel shows the mixing parameter $x$ dependence at $N_c = 3$ while the right panel is the $N_c$ dependence at $x = 0.8$. The meaning of each line is indicated in the figure.

strongly on the mixing parameter $x$, although the component of the composite $a_1$ is always larger than that of elementary ($z_a^{11} > z_a^{22}$) at $N_c = 3$ as shown in Fig. 4 (left panel).

In Fig. 4 (right panel), we show the $N_c$ dependence of the residues of pole-a at $x = 0.8$. There we find that the magnitudes of the residues of the pole-a, $z_a^{11}$ and $z_a^{22}$, interchange at $N_c \sim 3.5$. This indicates that the nature of the resonance *changes* as $N_c$ is varied. Thus, for the mixed system of elementary and composite components, the large $N_c$ limit does not always reflect the world at $N_c = 3$. This is a consequence of there being two sources of the $N_c$ dependence, one from the $a_1\pi\rho$ three-point vertex, which controls the mixing strength between the two basis states, and the other from the WT interaction, which determines the pole position ($s_p$) of the basis state for the composite $a_1$. Their competition determines the nature of the $a_1$ states. Actually, the $a_1\pi\rho$ interaction in Eq. (9) is of order $N_c^{-1/2}$ and vanishes in the large $N_c$ limit, where the two states of $q\bar{q}$ meson and hadronic composite states decouple. In the real world of $N_c = 3$, however, the interaction remains finite and gives an important contribution to the mixing dynamics as we showed. Therefore, we conclude that the large $N_c$ classification, which is often used to identify the character of resonances, does not necessarily work for such a mixed systems. In Ref. [25], we provide a general discussion of a two-level effective model of large $N_c$ behavior.

We have developed a general method to analyze the mixing structure of hadrons consisting of two components of quark and hadronic composites. As an example, the nature of the $a_1(1260)$ axial-vector meson has been explored. The present analysis points out theoretically that the $a_1$ meson has comparable amounts of the elementary $a_1$ component to the $\pi\rho$ composite $a_1$. Quest for evidences of the mixing nature in physical observables is an interesting future work.

The method proposed in the present paper makes it possible to discuss the validity of the classification in the large $N_c$ limit, which is considered to give a guiding principle to identify the nature of hadronic resonances. We have shown explicitly that the mixing nature of hadrons in the large $N_c$ limit could differ from that at finite $N_c = 3$. We conclude that the simple classification does not always work when admixture of components having different origins is important.

We express our thanks to T. Hyodo for useful discussions. This work is in part supported by Grant-in-Aid for Scientific Research on Priority Areas ‘‘Elucidation of New Hadrons with a Variety of Flavors’’ (Nos. 22105510 (H. N.), 22105509 (K. N.), 22105507 (D. J.), and E01:21105006 (A. H.)) from the ministry of Education, Culture, Sports, Science, and Technology of Japan. This work was done in part under the Yukawa International Program for Quark-hadron Sciences (YIPQS).

[1] S. Weinberg, Phys. Rev. **130**, 776 (1963); **137**, B672 (1965).
[2] V. Baru, J. Haidenbauer, C. Hanhart, Y. Kalashnikova, and A. E. Kudryavtsev, Phys. Lett. B **586**, 53 (2004).
[3] T. Hyodo, D. Jido, and A. Hosaka, Phys. Rev. C **78**, 025203 (2008).
[4] Y. Funaki, A. Tohsaki, H. Horiuchi, P. Schuck, and G. Ropke, Phys. Rev. C **67**, 051306 (2003).
[5] S. Weinberg, Phys. Rev. Lett. **18**, 507 (1967).
[6] C. W. Bernard, A. Duncan, J. LoSecco, and S. Weinberg, Phys. Rev. D **12**, 792 (1975).
[7] G. Ecker, J. Gasser, A. Pich, and E. de Rafael, Nucl. Phys. **B321**, 311 (1989).
[8] A. Dhar and S.R. Wadia, Phys. Rev. Lett. **52**, 959 (1984).
[9] A. Hosaka, Phys. Lett. B **244**, 363 (1990).
[10] M. Wingate, T. A. DeGrand, S. Collins, and U. M. Heller, Phys. Rev. Lett. **74**, 4596 (1995).
[11] M. Bando, T. Kugo, and K. Yamawaki, Phys. Rep. **164**, 217 (1988); N. Kaiser and U. G. Meissner, Nucl. Phys. **A519**, 671 (1990).
[12] T. Sakai and S. Sugimoto, Prog. Theor. Phys. **113**, 843 (2005); **114**, 1083 (2005).
[13] K. Nawa, H. Suganuma, and T. Kojo, Phys. Rev. D **75**, 086003 (2007).
[14] L. Roca, E. Oset, and J. Singh, Phys. Rev. D **72**, 014002 (2005).
[15] M. F. M. Lutz and E. E. Kolomeitsev, Nucl. Phys. **A730**, 392 (2004).
[16] L. Roca, A. Hosaka, and E. Oset, Phys. Lett. B **658**, 17 (2007); H. Nagahiro, L. Roca, and E. Oset, Phys. Rev. D **77**, 034017 (2008); H. Nagahiro, L. Roca, A. Hosaka, and E. Oset, Phys. Rev. D **79**, 014015 (2009).
[17] M. Wagner and S. Leupold, Phys. Rev. D **78**, 053001 (2008).
[18] D. Gomez Dumm, A. Pich, and J. Portoles, Phys. Rev. D **69**, 073002 (2004); D. Dumm, P. Roig, A. Pich, and J. Portoles, Phys. Lett. B **685**, 158 (2010).
[19] E. Witten, Nucl. Phys. **B160**, 57 (1979).
[20] R. L. Jaffe, AIP Conf. Proc. **964**, 1 (2007).
[21] J. R. Peláez and G. Ríos, Acta Phys. Pol. B **2**, 215 (2009).
[22] D. Lurie and A. J. Macfarlane, Phys. Rev. **136**, B816 (1964); D. Lurie, *Particle and Fields* (Interscience Publishers, New York, 1968).
[23] M. Doring, C. Hanhart, F. Huang, S. Krewald, and U. G. Meissner, Phys. Lett. B **681**, 26 (2009).
[24] L. S. Geng, E. Oset, J. R. Pelaez, and L. Roca, Eur. Phys. J. A **39**, 81 (2009).
[25] K. Nawa, H. Nagahiro, S. Ozaki, D. Jido, and A. Hosaka (unpublished).