Decay patterns of low-lying $N_{s\bar{s}}$ states via strangeness channels

Chun-Sheng An,[1] Ju-Jun Xie,[2] and Gang Li[3,*]
[1]School of Physical Science and Technology, Southwest University, Chongqing 400715, China
[2]Institute of Modern Physics, Chinese Academy of Sciences, Lanzhou 730000, China
[3]School of Physics and Engineering, Qufu Normal University, Qufu 273165, China

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Here we investigate the decay patterns of the low-lying hidden strangeness nucleon resonances ($\equiv N_{s\bar{s}}$) via the strangeness channels by employing the chiral Lagrangian approach, where the $N_{s\bar{s}}$ states are treated as compact pentaquark states. The S-wave decays of these states to the PB (pseudoscalar meson and baryon) and VB (vector meson and baryon) channels are studied. According to the obtained masses and decay properties, we find four states, namely, $N_{s\bar{s}}(1874)$ with quantum numbers $I(J^P) = 1/2(1/2^-)$, $N_{s\bar{s}}(1885)$ with $I(J^P) = 1/2(3/2^-)$, $N_{s\bar{s}}(2327)$ with $I(J^P) = 1/2(1/2^-)$, and $N_{s\bar{s}}(2252)$ with $I(J^P) = 1/2(3/2^-)$, may be associated with the well-established nucleon resonances $N^*(1895)$ and $N^*(1875)$ and with the resonances $N^*(2355)$ and $N^*(2250)$ newly predicted by the CLAS Collaboration, respectively. In addition, several other obtained hidden strangeness nucleon resonances may be expected to be dominant components of the predicted missing resonances in the literature.

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I. INTRODUCTION

Investigation of the hadron exotic states is always one of the most interesting subjects in hadronic physics. On both theoretical and experimental sides, great effort has been made to search for the exotic states, and lots of candidates for exotic mesons have been observed in the last decade, while most of the previous evidence on the existence of pentaquark states has been controversial, until the observation of two P_c^+ states was announced by the LHCb Collaboration in 2015 [1]. For recent reviews on the exotic states, see Refs. [2–5]. Note that kinds of the hidden charm pentaquark states like P_c^+ were first predicted in Ref. [6] and systematically studied using the constituent quark model in Ref. [7], and it is argued that the hidden charm states couple strongly to the charmness channels such as $J/\psi N$ because of the existence of the $c\bar{c}$ pair. On the other hand, a recent investigation of the charmness-nucleon σ term indicates that there should be $\approx 0.6\%$ charmness components in the ground state of the nucleon [8], and based on this result, one may also expect the existence of the nucleon resonances with $>99\%$ hidden charm pentaquark components above 4 GeV.

Analogously, in the strangeness sector, there should be sizable hidden strange pentaquark components in the nucleon and its excitations. In Refs. [9–12], the strangeness contributions to the magnetic moment, the spin, and the magnetic form factor of the nucleon are investigated, and the results show that the experimental data for the strangeness observable of the nucleon could be well reproduced by considering the compact strangeness components in the nucleon wave function. Meanwhile, it is claimed that the sea quark-antiquark pairs contribution to properties of the nucleon should be significant [13–15]. Besides, it is shown that the strangeness components in $N^*(1535)$ should account for the mass ordering of $N^*(1440)$, $N^*(1535)$, and $\Lambda^*(1405)$ [16] and the strong coupling of $N^*(1535)$ to the strangeness channels [17–19], which are consistent with the predictions of chiral unitary theory [20–22]. In addition, it is shown that data for the electromagnetic and strong decays of $N^*(1535)$ can also be well fitted by taking the strangeness contributions into account [23–25].

Recently, triggered by the observation of the P_c^+ pentaquark states, the hidden strange pentaquark states such as $\eta' N$ and ϕN bound states [26], $K\Sigma^*$ and $K^*\Sigma$ bound states [27,28], and compact five-quark states [29] were investigated. In Ref. [30], a possible ϕp resonance was investigated in the $\Lambda_c^+ \to \pi^0 \phi p$ decay by considering a triangle singularity mechanism, where the obtained ϕp invariant mass distribution agrees with the existing data. It is very interesting that the obtained hidden strange pentaquark states in those works lie in the energy range of several nucleon resonances with negative parity located, as listed in the Particle Data Group (PDG) reviews [31].

In Ref. [32], the strong decay behavior of the hidden strange meson-baryon molecular state was studied. Measurements on the decays of J/ψ and $\psi(2S)$ to $nK_s^0\bar{\Lambda}$ indicated that the nucleon excitations $N^*(1535)$, $N^*(1875)$, and $N^*(2120)$ may couple strongly to the $K\Lambda$ channel [33]. Moreover, it has been claimed that the nucleon resonances lying at ≈ 2 GeV contribute significantly to the ϕN production [34,35]. All the above evidence shows us that the hidden strange pentaquark configuration may be the dominant or notable component in some nucleon resonances.

*gli@mail.qfnu.edu.cn

Consequently, we investigate the decay patterns of the low-lying compact hidden strange pentaquark states ($N_{s\bar{s}}$) using the chiral Lagrangian approach. Wave functions for the $J^P = 1/2^-$ states obtained in Ref. [29] are employed, and those for the $J^P = 3/2^-$ states are derived using the same method. Limited by the employed approach, only the S-wave decay patterns for $N_{s\bar{s}}$ to strangeness channels are roughly estimated in the present work.

The paper is organized as follows. In Sec. II, we briefly present the employed formalism. The numerical results and discussions are given in Sec. III. Finally, Sec. IV presents a summary and conclusions.

II. FORMALISM

If the final meson is assumed to be emitted by a quark, to investigate the transitions $N_{s\bar{s}} \to PB(VB)$ (PB, pseudoscalar meson and baryon; VB, vector meson and baryon), one needs the explicit Lagrangian for quark-meson-quark interaction and the wave functions of the $N_{s\bar{s}}$ states. The wave functions for the hidden strange nucleon resonances have been explicitly studied in Refs. [29,36], and for the quark-quark-meson interaction, we employ the chiral Lagrangian approach, which is widely used in Refs. [37–41]. Accordingly, in this section, we briefly introduce the wave functions of the considered hidden strange nucleon resonances in Sec. II A, and we present a short review of the chiral Lagrangian approach and apply it to the five-quark system in Sec. II B.

A. Wave functions of the $N_{s\bar{s}}$ states

Following Refs. [29,36], a general wave function for the low-lying hidden strange nucleon resonances with $J^P = S^-$ can be written as

$$\psi^{(i)}_{t,s} = \sum_{a,b,c}\sum_{Y,y,T_z,t_z}\sum_{S_{4z},s_z} C^{[1^4]}_{[31]_a[211]_a} C^{[31]_a}_{[F^{(i)}]_b[S^{(i)}]_c}[F^{(i)}]_{b,Y,T_z} \times [S^{(i)}]_{c,S_{4z}}[211;C]_a(Y,T,T_z,y,\bar{t},t_z|1,1/2,t) \times (S_4,S_{4z},1/2,s_z|S,S_z)\bar{\chi}_{y,t_z}\bar{\xi}_{s_z}\varphi_{[5]}\,, \quad (1)$$

where $\bar{\chi}_{y,t_z}$ and $\bar{\xi}_{s_z}$ represent the isospinor and spinor of the antiquark, respectively, and $\varphi_{[5]}$ represents the completely symmetrical orbital wave function. The first summation involves symbols $C^{[.]}_{[..][...]}$, which are S_4 Clebsch-Gordan (CG) coefficients for the indicated color ([211]), flavor-spin ([31]), and flavor ([F]) and spin ([S]) wave functions of the $qqqq$ system. The second summation runs over the flavor indices in the SU(3) CG coefficient (with nine symbols) and the third one runs over the spin indices in the standard SU(2) CG coefficient.

Explicitly, according to Eq. (1), there are five possible pentaquark configurations that have appropriate symmetry structure and spin 1/2:

$$|1(1/2^-)\rangle = |qqqs([4]_X[211]_C[31]_{FS}[211]_F[22]_S)\otimes\bar{s},$$
$$|2(1/2^-)\rangle = |qqqs([4]_X[211]_C[31]_{FS}[211]_F[31]_S)\otimes\bar{s},$$
$$|3(1/2^-)\rangle = |qqqs([4]_X[211]_C[31]_{FS}[22]_F[31]_S)\otimes\bar{s},$$

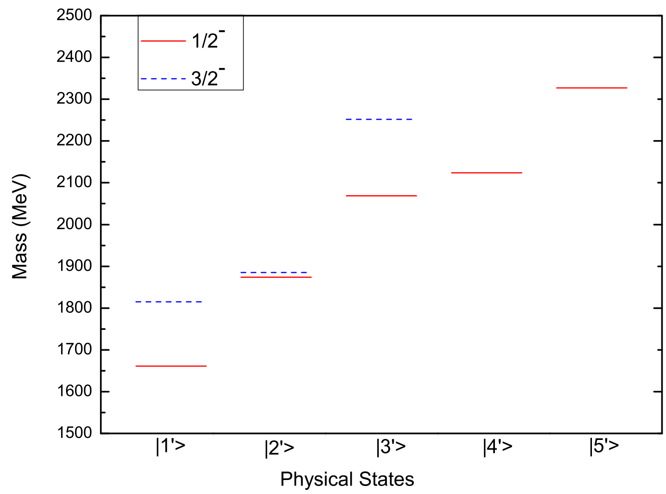

FIG. 1. Spectrum of the low-lying hidden strange nucleon resonances with $J^P = 1/2^-$ shown by the solid red lines and $J^P = 3/2^-$ shown by the dashed blue lines.

$$|4(1/2^-)\rangle = |qqqs([4]_X[211]_C[31]_{FS}[31]_F[22]_S)\otimes\bar{s},$$
$$|5(1/2^-)\rangle = |qqqs([4]_X[211]_C[31]_{FS}[31]_F[31]_S)\otimes\bar{s}, \quad (2)$$

and three configurations with spin 3/2:

$$|1(3/2^-)\rangle = |qqqs([4]_X[211]_C[31]_{FS}[211]_F[31]_S)\otimes\bar{s},$$
$$|2(3/2^-)\rangle = |qqqs([4]_X[211]_C[31]_{FS}[22]_F[31]_S)\otimes\bar{s},$$
$$|3(3/2^-)\rangle = |qqqs([4]_X[211]_C[31]_{FS}[31]_F[31]_S)\otimes\bar{s}. \quad (3)$$

In Ref. [29], the spectrum of the above hidden strange nucleon resonances with spin 1/2 was studied using the one-gluon-exchange model (OGE) and the Goldstone boson exchange model (GBE), respectively. One may note that the numerical results obtained using the OGE model (depicted in Fig. 1 by the solid red line) should be more reasonable. So hereafter we employ only the wave functions obtained in the OGE model, and the explicit probability amplitudes for the mixing between $|i(1/2^-)\rangle$ configurations can be found in Ref. [29]. One should note that the numerical results for the energies of the $I(J^P) = 1/2(1/2)^-$ states were obtained by using the empirical values for model parameters reported in the literature; if one changes values of the coupling strength in the OGE model by $\mp 10\%$, then the obtained masses for the five $I(J^P) = 1/2(1/2)^-$ $N_{s\bar{s}}$ states will be 1661 ± 51, 1874 ± 30, 2068 ± 11, 2124 ± 6, and 2327 ± 15 MeV, respectively. On the other hand, we apply the OGE model in Ref. [29] to the $3/2^-$ $N_{s\bar{s}}$ sector. Using the wave functions as in Eq. (3) and all the same parameters as in Ref. [29], we can get the numerical results shown in Fig. 1 by the dashed blue lines; the explicit values for the energies of the three obtained physical states are 1815 ± 36, 1885 ± 29, and 2252 ± 7 MeV, respectively.

Finally, using the empirical values for the coupling strength in the OGE model, the wave functions for the physical states

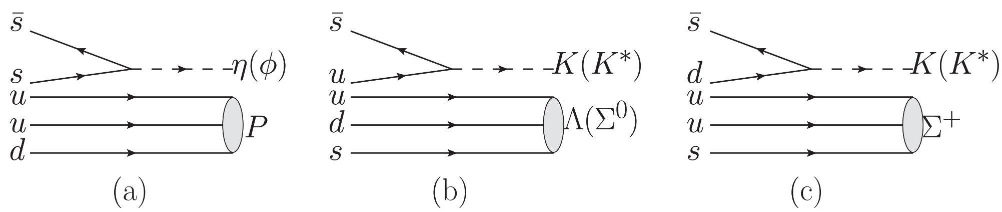

FIG. 2. The effective coupling of the $N_{s\bar{s}}$ states to $\eta(\phi)p$ (a), $K(K^*)\Lambda$ (b), and $K(K^*)\Sigma$ (c).

with $J^P = 3/2^-$ are as follows:

$$\begin{aligned}|1'(3/2^-)\rangle &= 0.989|1(3/2^-)\rangle - 0.151|2(3/2^-)\rangle\\ &\quad + 0.008|3(3/2^-)\rangle,\\ |2'(3/2^-)\rangle &= 0.150|1(3/2^-)\rangle + 0.985|2(3/2^-)\rangle\\ &\quad + 0.082|3(3/2^-)\rangle,\\ |3'(3/2^-)\rangle &= -0.020|1(3/2^-)\rangle - 0.080|2(3/2^-)\rangle\\ &\quad + 0.997|3(3/2^-)\rangle\,. \end{aligned} \tag{4}$$

B. The chiral Lagrangian approach

In this work, we assume that the final meson couples directly to a quark and the $\bar{s}$ quark, namely, a strange quark and the $\bar{s}$ quark are annihilated to emit the η or ϕ meson, while annihilation of a u or d quark with the $\bar{s}$ quark will emit a K or K^* meson, such kinds of quark-meson effective couplings are shown in Fig. 2.

In the chiral Lagrangian approach, the Hamiltonian for the quark and the pseudoscalar meson is

$$H_{\rm eff}^{Pqq} = \sum_j \bar{\psi}_j \gamma_\mu^j \gamma_5^j \psi_j \partial^\mu \phi_m, \tag{5}$$

where the summation on j runs over the quark in the initial hadron, ψ_j represents the quark field, and ϕ_m represents the pseudoscalar meson field. In the nonrelativistic approximation, Eq. (5) leads to the transition operator

$$\begin{aligned}T_d^{Pqq} &= \sum_j \left(\frac{\omega_M}{E_f + M_f}\sigma\cdot\vec{P}_f + \frac{\omega_M}{E_i + M_i}\sigma\cdot\vec{P}_i - \sigma\cdot\vec{k}_M \right.\\ &\quad \left. + \frac{\omega_M}{2\mu_q}\sigma\cdot\vec{p}_j \right) X_M^j \exp\{-i\vec{k}_M\cdot\vec{r}_j\},\end{aligned} \tag{6}$$

for the transition caused by the processes $q \to q'M$, and

$$\begin{aligned}T_a^{Pqq} &= \sum_j (m_j + m_{\bar{q}}) \mathcal{C}_{XFSC}^j \bar{\chi}_z^\dagger \mathcal{I}_2 \chi_z^j X_M^j\\ &\quad \times \exp\{-i\vec{k}_M\cdot(\vec{r}_j + \vec{r}_{\bar{q}})/2\},\end{aligned} \tag{7}$$

for the transitions caused by the processes $q\bar{q} \to M$. Accordingly, Eq. (7) is used in the calculations on the transition matrix elements of $N_{s\bar{s}} \to PB$ processes. And in Eq. (7), m_j and $m_{\bar{q}}$ are the constituent masses of the jth quark and the antiquark, respectively, $\mathcal{C}_{XFSC}^j$ denotes the operator to calculate the orbital, flavor, spin, and color overlap factor between the residual wave function of the pentaquark configuration after the quark-antiquark annihilation and the wave function of the final baryon, $\bar{\chi}_z^\dagger \mathcal{I}_2 \chi_z^j$ is the spin operator for the quark-antiquark annihilation, and X_M^j is the operator for a pseudoscalar meson emission, which can be defined as

$$\begin{aligned}&X_{K^\pm}^j = \mp\frac{1}{\sqrt{2}}(\lambda_4^j \mp i\lambda_5^j),\ X_{K^0,\bar{K}^0}^j = \mp\frac{1}{\sqrt{2}}(\lambda_6^j \mp i\lambda_7^j),\\ &X_\eta^j = \cos\theta\lambda_8^j - \sin\theta\sqrt{\frac{2}{3}}\mathcal{I},\quad X_{\eta'}^j = \sin\theta\lambda_8^j + \cos\theta\sqrt{\frac{2}{3}}\mathcal{I},\end{aligned} \tag{8}$$

for K, η, and η' emissions, where λ_i^j is the flavor SU(3) Gell-Mann matrices acting on the jth quark, $\mathcal{I}$ is the unit matrix in three-dimensional space, and θ is the mixing angle between η_1 and η_8, leading to the physical η and η'. Here we take the value $\theta = -23^\circ$ [42].

The Hamiltonian for the quark and the vector meson reads

$$H_{\rm eff}^{Vqq} = -\sum_j \bar{\psi}_j \left(a\gamma_\mu^j + \frac{ib\sigma_{\mu\nu}k_M^\nu}{2m_j} \right) \phi_m^\mu \psi_j, \tag{9}$$

where m_j is the constituent mass of the jth quark, k_M^ν represents the four-momentum of the vector meson, ϕ_m^μ is the vector meson field, and a and b are the vector and tensor coupling constants, respectively.

The quark–vector meson coupling in Eq. (9) results in the following transition operators,

$$\begin{aligned}T_{d,T}^{Vqq} &= \sum_j \left\{ i\frac{b'}{2m_j}\vec{\sigma}_j\cdot(\vec{k}_M\times\vec{\epsilon}) + \frac{a}{2\mu_q}\vec{p}_j\cdot\vec{\epsilon} \right\} X_M^j\\ &\quad \times \exp\{-i\vec{k}_M\cdot\vec{r}_j\},\end{aligned} \tag{10}$$

$$T_{d,L}^{Vqq} = \sum_j \frac{aM_V}{|\vec{q}|} X_M^j \exp\{-i\vec{k}_M\cdot\vec{r}_j\}, \tag{11}$$

for the transitions caused by the processes $q \to q'M$ with the emitted vector meson transversely polarized denoted by T and longitudinally polarized denoted by L, respectively, and in the following operators,

$$\begin{aligned}T_{a,T}^{Vqq} &= \sum_j \left\{ a - \frac{m_j + m_{\bar{q}}}{2m_j} b \right\} \vec{\sigma}\cdot\vec{\epsilon} X_V^j\\ &\quad \times \exp\{-i\vec{k}_M\cdot(\vec{r}_j + \vec{r}_{\bar{q}})/2\},\end{aligned} \tag{12}$$

$$\begin{aligned}T_{a,L}^{Vqq} &= \sum_j \left\{ a - \frac{m_j + m_{\bar{q}}}{2m_j} b \right\} \frac{E_V\vec{\sigma}\cdot\vec{q}}{M_V|\vec{q}|} X_V^j\\ &\quad \times \exp\{-i\vec{k}_M\cdot(\vec{r}_j + \vec{r}_{\bar{q}})/2\},\end{aligned} \tag{13}$$

for the transitions caused by the processes $q\bar{q} \to M$, where X_V^j is the vector meson emission operator that is defined very

TABLE I. The transition matrix elements for the hidden strange pentaquark configurations with $J^P = 1/2^-$ to PB and VB strangeness channels. Note that the following common factors are omitted: $(m_q + m_s)\langle\hat{\mathcal{O}}_X\rangle$ for $|i(1/2^-)\rangle \to PB$ transitions, in addition, $(2\cos\theta + \sqrt{2}\sin\theta)$ for ηN and $(2\sin\theta - \sqrt{2}\cos\theta)$ for $\eta' N$ decays; while for $|i(1/2^-)\rangle \to VB$ transitions, a common factor $(a - \frac{m_q+m_s}{2m_q}b)\langle\hat{\mathcal{O}}_X\rangle$ for the transitions with the final meson transversely polarized and $(a - \frac{m_q+m_s}{2m_q}b)\frac{E_V}{m_V}\langle\hat{\mathcal{O}}_X\rangle$ for the transitions with the final meson longitudinally polarized.

	$\|1\rangle$	$\|2\rangle$	$\|3\rangle$	$\|4\rangle$	$\|5\rangle$
ηN	$\sqrt{3}/3$	-1	$-\sqrt{2/3}$	$\sqrt{3}/3$	$\sqrt{3}/3$
$K\Lambda$	$-1/\sqrt{3}$	1	$\sqrt{6}$	$\sqrt{3}$	$\sqrt{3}$
$K\Sigma$	$-\sqrt{3}$	3	$-\sqrt{6}$	$\sqrt{3}/3$	$\sqrt{3}/3$
$\eta' N$	$\sqrt{3}/3$	-1	$-\sqrt{2/3}$	$\sqrt{3}/3$	$\sqrt{3}/3$
$K^*\Lambda^{(T)}$	$\sqrt{2/3}$	$\sqrt{2}/3$	$2/\sqrt{3}$	$-\sqrt{6}$	$\sqrt{2/3}$
$K^*\Lambda^{(L)}$	$1/\sqrt{3}$	$1/3$	$\sqrt{2/3}$	$-\sqrt{3}$	$1/\sqrt{3}$
$K^*\Sigma^{(T)}$	$\sqrt{6}$	$\sqrt{2}$	$-2/\sqrt{3}$	$-\sqrt{2/3}$	$\sqrt{6}/9$
$K^*\Sigma^{(L)}$	$\sqrt{3}$	1	$-\sqrt{2/3}$	$-1/\sqrt{3}$	$\sqrt{3}/9$
$\phi N^{(T)}$	2	$2/\sqrt{3}$	$2\sqrt{2}/3$	2	$-2/3$
$\phi N^{(L)}$	$\sqrt{2}$	$\sqrt{2/3}$	$2/3$	$\sqrt{2}$	$-\sqrt{2}/3$

similarly to X^j_M in Eq. (8) as

$$X^j_{K^{*\pm}} = \mp\frac{1}{\sqrt{2}}(\lambda^j_4 \mp i\lambda^j_5),$$
$$X^j_{K^{*0},\bar{K}^{*0}} = \mp\frac{1}{\sqrt{2}}(\lambda^j_6 \mp i\lambda^j_7), \tag{14}$$
$$X^j_\phi = \frac{\sqrt{2}}{3}\mathcal{I}^j - \sqrt{\frac{2}{3}}\lambda^j_8,$$

for ρ, K^* and $\bar{K}^*$, and ϕ emission. E_V and M_V are the energy and mass of the final meson, and here we have taken the polarization vector of the final meson to be

$$\epsilon^L_\mu = \frac{1}{M_V}\begin{pmatrix}|\vec{k}_M|\\ E_V\frac{\vec{k}_M}{|\vec{k}_M|}\end{pmatrix}, \quad \epsilon^T_\mu = \begin{pmatrix}0\\ \vec{\epsilon}\end{pmatrix}, \tag{15}$$

with $\vec{\epsilon}(\pm) = 1/\sqrt{2}(\mp 1, -i, 0)^T$. The three-momentum $\vec{k}_M$ is written as

$$|\vec{k}_M| = \frac{\sqrt{[M_i^2 - (M_f + m_M)^2][M_i^2 - (M_f - m_M)^2]}}{2M_i}, \tag{16}$$

where M_i, M_f, and m_M denote masses of the initial $N_{s\bar{s}}$ state, the final baryon, and the meson, respectively.

III. RESULTS AND DISCUSSIONS

Taking the wave functions of the pentaquark configurations $|i(1/2^-)\rangle$ and $|i(3/2^-)\rangle$ and the transition operators given in Sec. II, we obtain the S-wave transition elements for the configurations with $J^P = 1/2^-$ in Table I and those for configurations with $J^P = 3/2^-$ in Table II, respectively. Note that in the common factors given in the captions of these tables $m_q = m_s$ applies for the transitions with final states ηN, $\eta' N$, and ϕN; $m_q = m$ applies for the transitions with all the other final states; and $\langle\hat{\mathcal{O}}_X\rangle$ is the orbital matrix element that depends on the three-momentum $\vec{k}_M$, which reads

$$\langle\hat{\mathcal{O}}_X\rangle \propto \exp\{-3k_M^2/20\omega^2\}, \tag{17}$$

where ω is the harmonic oscillator parameter.

TABLE II. The transition matrix elements for $|i(3/2^-)\rangle \to VB$. Common factors: for $N_{s\bar{s}} \to VB$ decays, $(a - \frac{m_q+m_s}{2m_q}b)\langle\hat{\mathcal{O}}_X\rangle$ for the transitions with the final meson transversely polarized and $(a - \frac{m_q+m_s}{2m_q}b)\frac{E_V}{m_V}\langle\hat{\mathcal{O}}_X\rangle$ for the transitions with the final meson longitudinally polarized.

	$\|1\rangle$	$\|2\rangle$	$\|3\rangle$
$K^*\Lambda^{(T)}$	$-2/\sqrt{3}$	$-2/3$	$-2\sqrt{2/3}$
$K^*\Lambda^{(L)}$	$2\sqrt{2/3}$	$2\sqrt{2}/3$	$4/\sqrt{3}$
$K^*\Sigma^{(T)}$	$-2\sqrt{3}$	-2	$2\sqrt{2/3}$
$K^*\Sigma^{(L)}$	$2\sqrt{6}$	$2\sqrt{2}$	$-4/\sqrt{3}$
$\phi N^{(T)}$	$-2\sqrt{2}$	$-2\sqrt{2/3}$	$-4/3$
$\phi N^{(L)}$	4	$4/\sqrt{3}$	$4\sqrt{2}/3$

To get the numerical results, here we take the explicit empirical values for the model parameters: $m = 340$ MeV and $m_s = 460$ MeV for the constituent masses of the quarks [43], $f_K = f_\eta = 160$ MeV and $f_{\eta'} = 280$ MeV for the decay constants of the mesons [41,44], $\omega = 225$ MeV for the harmonic oscillator parameter [29], $a = -3$ and $b = 2$ for the vector and tensor coupling constants [41], and finally, masses of the final hadrons are taken from the PDG [31]. With these values for the parameters, we obtain the numerical results listed in Table III, where we name the obtained $N_{s\bar{s}}$ pentaquark states according to their masses obtained in this work. Because the partial decay widths of the $N_{s\bar{s}}$ states depend on a quark-meson coupling constant that falls in a large range, in Table III, we only show estimations on the ratios of the transition matrix element squares,

$$\mathcal{F} \equiv \frac{1}{2J_i+1}\sum_{J_{iz},J_{fz}} |\langle B, J_{fz}|T|N_{s\bar{s}}, J_{iz}\rangle|^2, \tag{18}$$

where J_i and J_f are the total angular momenta of the initial and final baryon states, respectively, and T represents the operators given in Eqs. (7), (12), and (13) for corresponding transitions. The ratios for the $J^P = 1/2^-$ states listed in Table III are defined as the obtained $\mathcal{F}$ for $N_{s\bar{s}} \to PB(VB)$ transitions over that for the $N_{s\bar{s}}(1661) \to \eta N$ transition, while the numerical results for $N_{s\bar{s}}(2252)$ with $J^P = 3/2^-$ are the ratios of the obtained $\mathcal{F}$ for the corresponding channels over that for the $N_{s\bar{s}}(2252) \to K^*\Lambda$ channel.

In Table III, the numerical results for the obtained $J^P = 1/2^-$ states to both PB and VB channels are presented, but for the $J^P = 3/2^-$ sector, only the results for $N_{s\bar{s}}(2252) \to VB$ are shown. Because the $N_{s\bar{s}}$ states with $J^P = 3/2^-$ decay into PB channels via the D wave, the partial decay widths of these channels should be much smaller than those of the VB channels. On the other hands, the two lower $N_{s\bar{s}}$ states with $J^P = 3/2^-$ are below the thresholds of all the studied VB strangeness channels. As we can see in a very recent work [32], the partial decay widths of the $K^*\Sigma$ and $K\Sigma^*$

TABLE III. Numerical results for the ratios of the transition matrix element squares $\mathcal{F}$.

	$J^P = 1/2^-$					$J^P = 3/2^-$
	$N_{s\bar{s}}(1661)$	$N_{s\bar{s}}(1874)$	$N_{s\bar{s}}(2069)$	$N_{s\bar{s}}(2124)$	$N_{s\bar{s}}(2327)$	$N_{s\bar{s}}(2252)$
ηN	1	0.19	0.10	0.03	≈ 0	–
$K\Lambda$	0.38	1.48	0.08	0.05	≈ 0	–
$K\Sigma$	–	0.65	0.11	≈ 0	≈ 0	–
$\eta' N$	–	–	0.19	0.05	≈ 0	–
$K^*\Lambda$	–	–	0.98	0.67	0.47	1
$K^*\Sigma$	–	–	–	4.05	0.22	2.22
ϕN	–	–	1.47	0.58	0.11	0.35

molecular states have been investigated, and the numerical results show that the branching ratios for $J^P = 3/2^-$ $N_{s\bar{s}}$ states are very small.

The lowest state $N_{s\bar{s}}(1661)$ lies at the energy very close to $N^*(1650)$, and $\approx$130 MeV higher than $N^*(1535)$; one may expect this state to be a sizable component in the two lowest S_{11} states. Note that the spin symmetry of the three-quark component of $S_{11}(1650)$ in the traditional quark model is expected to be the completely symmetric $[3]_S$ [43], which should weaken the transition between the three-quark component of $S_{11}(1650)$ and the present obtained $N_{s\bar{s}}(1661)$. In addition, as we can see in Table III, $N_{s\bar{s}}(1661)$ should couple strongly to the ηN and $K\Lambda$ channels; this is consistent with the large branching ratio of the $N^*(1535)$ resonance to the ηN channel [31] and the strong coupling of $N^*(1535)$ to the $K\Lambda$ channel predicted by the isobar model [16] and chiral perturbation theory [22]. One may notice that the nonvanishing coupling constant of $N^*(1650)$ to the $K\Sigma$ channel was reported in a very recent work [45], while as we can see in Ref. [25], strong coupling between $N^*(1650)$ and $K\Sigma$ can be obtained by taking probabilities of strangeness components in $N^*(1650)$ smaller than those in $N^*(1535)$.

Another nontrivial result is for the decay pattern of the obtained $N_{s\bar{s}}(1874)$, which falls in the energy range of the nucleon resonance $N^*(1895)$ [31]; considering the uncertainties of the present model, one may expect $N_{s\bar{s}}(1874)$ to be a dominant or sizable component in $N^*(1895)$. On the other hand, a very recent measurement on the $K^*\Lambda$ photoproduction showed that $N^*(1895)$ should contribute significantly to the $\gamma p \to K^*\Lambda$ reaction [46], and a partial-wave analysis on the $\gamma p \to K^+\Lambda$ and $\pi^- p \to K^0\Lambda$ reactions indicated that $N^*(1895)$ was unquestionably required in these processes [47]. Moreover, the measurement on the $\gamma p \to \eta p$ and $\gamma p \to \eta' p$ reactions showed strong couplings between $N^*(1895)$ and both ηp and $\eta' p$ [48]. All of this evidence is consistent with large strangeness components in the $N^*(1895)$ resonance. As we can see in Table III, the obtained $N_{s\bar{s}}(1874)$ in the present model may couple strongly to the $K\Lambda$ and $K\Sigma$ channels; the ratio of the coupling constants for the $N_{s\bar{s}}(1874)$ resonance to the $K\Lambda$, $K\Sigma$, and ηN channels is found to be $|g_{N_{s\bar{s}}K\Lambda} : g_{N_{s\bar{s}}K\Sigma} : g_{N_{s\bar{s}}\eta N}| \approx 2.81 : 1.86 : 1$, Accordingly, we can conclude that $N_{s\bar{s}}(1874)$ may be the dominant component of the nucleon resonance $N^*(1895)$.

The obtained $N_{s\bar{s}}(2069)$ and $N_{s\bar{s}}(2124)$ resonances in the present model fall in the energy range of the two-star nucleon resonance $N(2120)$ listed in Ref. [31], whose spin is identified to be 3/2, although the spin 1/2 may not be completely excluded. One can also expect $N_{s\bar{s}}(2069)$ and $N_{s\bar{s}}(2124)$ could be related to the missing nucleon resonances $N^*(2030)/N^*(2070)$ and $N^*(2145)/N^*(2195)$ predicted in Ref. [49], respectively, because all these resonances are predicted to couple with strangeness channels. Moreover, two S_{11} resonances located at 1846 ± 47 and 2113 ± 70 MeV were predicted by using a dynamical coupled channel approach in Ref. [50]; obviously, one can also associate the presently obtained $N_{s\bar{s}}(2124)$ with the latter one. About the decay properties of $N_{s\bar{s}}(2069)$ and $N_{s\bar{s}}(2124)$, as one can see in Table III, these two obtained states seem to couple strongly to the strangeness VB channels. For instance, if one compares the coupling constants for $N_{s\bar{s}}(2124)$ to the $K^*\Sigma$ and $K\Lambda$ channels, the obtained ratio is $|g_{N_{s\bar{s}}K^*\Lambda} : g_{N_{s\bar{s}}K\Lambda}| \approx 8.97 : 1$. Consequently, if one assumes that the presently obtained $N_{s\bar{s}}(2124)$ with $J^P = 1/2^-$ corresponds to a dominant component in $N(2120)$, significant evidence of this resonance must be shown in the reactions with the $K^*\Sigma$ final state, otherwise this assumption should not be appropriate. Meanwhile, it has been claimed that the nucleon resonances lying at $\approx$2 GeV may contribute to $N\phi$ production significantly [34,35]; this seems to coincide with our findings on the $N_{s\bar{s}}(2069)$ pentaquark state. As we can see in Table III, the ratio of the coupling constants (absolute value) for $N_{s\bar{s}}(2069)$ to ηN, $K^*\Lambda$, and ϕN is $\approx 1 : 3.17 : 3.86$. Finally, the highest $N_{s\bar{s}}(2327)$ state also shows strong coupling to the strangeness VB channels, but there is no solid experimental evidence on the nucleon resonance with negative parity around or above 2300 MeV [31], while one may expect that the $N_{s\bar{s}}(2327)$ state can be associated with the $N^*(2355)$ state with $J^P = 1/2^-$ predicted in Ref. [46], which should also couple to the $K^*\Lambda$ channel.

Next we come to the $3/2^-$ $N_{s\bar{s}}$ resonance sector. There are two states lying at $\approx 1850 \pm 50$ MeV, named $N_{s\bar{s}}(1815)$ and $N_{s\bar{s}}(1885)$, that are obtained in this work. Because these two obtained states are very close to the nucleon resonance $N^*(1875)$ with $J^P = 3/2^-$, especially the obtained $N_{s\bar{s}}(1885)$ state, one may expect them to take sizable components in $N^*(1875)$. While, as we have discussed above, the presently studied $N_{s\bar{s}}$ states with $J^P = 3/2^-$ decay into the PB channels via the D wave; hence the branching ratios for this kind of decay should be small [32]. In addition, all the strangeness pseudoscalar meson and decuplet baryon channels are above the thresholds of these two obtained hidden strange pentaquark states, so the decay channel with the large branching ratio may be the $\pi\Delta$ channel, although the coupling of $N_{s\bar{s}}$ states to nonstrange channels may be not so strong, the phase space for $N_{s\bar{s}}(1885)$ to $\pi\Delta$ channel is very large. However, the presently employed model is not well applicable to the transitions of $N_{s\bar{s}}$ states to nonstrange channels, so we can only give the above qualitative discussions. The highest obtained $N_{s\bar{s}}$ state, $N_{s\bar{s}}(2252)$, lies above 2 GeV, which couples strongly to the $K^*\Sigma$ and $K^*\Lambda$ channels, but relatively weakly to the ϕN channel; one may refer this obtained state to the predicted $N^*(2250)$ in Ref. [46].

Finally, as we have discussed above, one has to note the following points before going to the final conclusions. First, the obtained results in Table III are only ratios for the square of transition matrix elements of each $N_{s\bar{s}}$ state to the strangeness channels, but not the partial decay widths, which depend on the quark-meson coupling constant. Second, the present model is not well applicable to the transitions of $N_{s\bar{s}}$ states to nonstrangeness channels, e.g., $N_{s\bar{s}}(1885) \to \pi\Delta$, whose channel may be the one with the largest branching ratio of $N_{s\bar{s}}(1885)$.

IV. SUMMARY AND CONCLUSIONS

In the present work, we investigate the strong decay patterns of the low-lying hidden strange pentaquark states with $I(J^P) = 1/2(1/2^-)$ and $I(J^P) = 1/2(3/2^-)$ to the strangeness channels. Limited by the presently employed formalism, the decay patterns for nonstrange channels are not included; here only the S-wave decays of the $N_{s\bar{s}}$ states to the PB and VB strangeness channels are roughly estimated. In addition, as a direct extension of one of our previous works, we present the masses and wave functions of three $N_{s\bar{s}}$ pentaquark states with $I(J^P) = 1/2(3/2^-)$. One can see that the two obtained lower $N_{s\bar{s}}$ states are very close to the nucleon excitation $N^*(1875)$, and the highest one is located close to the $N^*(2250)$ resonance predicted by the CLAS Collaboration.

In the $1/2^-$ sector, the presently obtained $N_{s\bar{s}}(1661)$ resonance, which is ≈100 MeV higher than the nucleon resonance $N^*(1535)$, couples strongly to the ηN channel, which is in agreement with the large partial decay width of $N^*(1535) \to \eta N$ if we take $N_{s\bar{s}}(1661)$ to be the higher Fock component in the wave function of $N(1535)$. Note that the probability for $N_{s\bar{s}}(1661)$ in the $N^*(1650)$ resonance may be small, although $N_{s\bar{s}}(1661)$ lies so close to $N^*(1650)$, because the spin structure should weaken the coupling between the three-quark and the $N_{s\bar{s}}$ components. The $N_{s\bar{s}}(1874)$ state with $J^P = 1/2^-$ obtained in the present model couples very strongly to the $K\Lambda$ channel. This finding seems to coincide with the property of the nucleon resonance $N^*(1895)$; thus, one may expect $N_{s\bar{s}}(1874)$ to be the dominant component of $N^*(1895)$. For the $N_{s\bar{s}}(2069)$ state, the strangeness channel with the largest coupling is ϕN; this seems to be consistent with previous predictions that N^* located at ≈2 GeV should contribute significantly to the ϕN production. One can also expect the obtained $N_{s\bar{s}}(2069)$ and $N_{s\bar{s}}(2124)$ states could be related to the missing resonances predicted by Capstick and Roberts in Ref. [49]. The highest obtained hidden strange state, $N_{s\bar{s}}(2327)$, may be associated with the predicted $N^*(2355)$ resonance [46].

Two of the obtained $N_{s\bar{s}}$ states with $J^P = 3/2^-$, i.e., $N_{s\bar{s}}(1815)$ and $N_{s\bar{s}}(1885)$, which are very close to the nucleon resonance $N^*(1875)$, are below the thresholds of all strangeness VB channels. Meanwhile, these states decay into PB channels via the D wave, so these two states may mainly decay into the $\pi\Delta$ channel via the S wave, because the phase space for this channel is very large. We only present the numerical results for decay patterns of the $N_{s\bar{s}}(2252)$ resonance, which may mainly decays to $K^*\Sigma$ and $K^*\Lambda$ channels; this is consistent with the predicted $N^*(2250)$ resonance by a very recent measurement on K^* photoproduction [46]. We will wait for more precise experimental measurements on the K^* productions.

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