Strange Tribaryons as $\bar{K}$-Mediated Dense Nuclear Systems

Yoshinori Akaishi$^{a,b,*}$, Akinobu Doté$^{c,†}$ and Toshimitsu Yamazaki$^{b,‡}$

$^a$College of Science and Technology, Nihon University, Funabashi, Chiba 274-8501, Japan,

$^b$DRI, RIKEN, Wako, Saitama 351-0198, Japan

$^c$Institute of Particle and Nuclear Studies, KEK, Tsukuba, Ibaraki 305-0801, Japan

We discuss the implications of recently discovered strange tribaryons in $^4\text{He}(\text{stopped}-K^-,p)S^0(3115)$ and $^4\text{He}(\text{stopped}-K^-,n)S^1(3140)$ within the framework of deeply bound $\bar{K}$ states formed on shrunken nuclear cores. $S^1(3140)$ corresponds to $T=0$ $ppnK^-$, whereas $S^0(3115)$ to $T=1$ $pnnK^-$, which is an isobaric analog state of $pppK^-$, predicted previously. The observed binding energies can be accounted for by including the relativistic effect and by invoking a medium-enhanced $\bar{K}N$ interaction by 15%. We propose various experimental methods to further study these and related bound systems.

1. Introduction

In a series of publications in recent years [1,2,3,4,5,6], we have predicted deeply bound narrow $\bar{K}$ nuclear states based on bare $\bar{K}N$ interactions, which were derived from empirical data ($\bar{K}N$ scattering and kaonic hydrogen) together with the ansatz that $\Lambda_{1405}$ is a bound state of $\bar{K}N$. The presence of such hitherto unknown kaonic nuclear states results from a very attractive $\bar{K}N$ interaction in the $I=0$ channel, which persists to be strong for discrete bound states of finite nuclei, and causes not only a strong binding of $K^-$ in proton-rich nuclei, but also an enormous shrinkage of $\bar{K}$-bound nuclei despite the hard nuclear incompressibility. Thus, a $\bar{K}$ produces a bound state with a $\bar{K}$-mediated “condensed nucleus”, which does not exist by itself. For example, $ppp\bar{K}^-$ with a total binding energy of $-E_K^p = 48$ MeV, $ppn\bar{K}^-$ with $-E_K^n = 118$ MeV, and $ppp\bar{K}^-$ with $-E_K^p = 97$ MeV. The calculated rms distances in the $ppn\bar{K}^-$ system are: $R_{NN} = 1.50$ fm and $R_{KN-N} = 1.19$ fm, whereas $R_{K-N} = 1.31$ fm in $\Lambda_{1405}$. The $NN$ distance in $ppn\bar{K}^-$ is substantially smaller than the normal inter-nucleon distance ($\sim 2.0$ fm), and the average nucleon density, $\rho_{NN} = 3.1 \times 10^{-3}$, is much larger than the normal density, $\rho_0 = 0.17$ fm$^{-3}$. An observation of such a deeply bound state would not only confirm the underlying physics framework, but would also provide profound information on an in-medium modification of the $\bar{K}N$ interaction, nuclear compressibility, chiral symmetry restoration, kaon condensation regime and possible transition to a quark-gluon phase. Here, the predicted values for the separation energies and widths will serve as benchmarks to examine the above problems.

An experimental search for $ppnK^-$ in the $^4\text{He}(\text{stopped}-K^-,n)$ reaction [1] was carried out by experimental group E471 of Iwasaki et al. at KEK. The first evidence was found in the proton-emission channel [7].

$^4\text{He}(\text{stopped}-K^-,p)[N'NN\bar{K}]_{T=1}^{Q=0} : S^0(3115), (1)$

where a distinct narrow peak appeared at a mass of $M = 3117 \pm 5$ MeV/$c^2$ with a total binding energy (separation energy of $K^-$) of $S_K = -E_K = -(M - M_p - 2M_n - M_{K^-})c^2 = 194 \pm 5$ MeV and a width of less than 25 MeV. This state, populated in reaction [1], has a unique isospin, $(T,T_3) = (1,-1)$. It was an unexpected discovery, since the $T=1$ bound state of $K^-$ on a triton ($(ppn)_{T=1/2}$) had been predicted to be shallow and broad [11]. On the other hand, an ex-
otic $T = 1$ $pppK^-$ state with $S_K = 97$ MeV had been previously predicted [25], and the observed $S^0(3115)$ can be identified as its isobaric analog state.

An indication of another species was observed in the neutron-emission channel [3].

$^4\text{He(stopped--$K^-$, $n$)}[N'NN\bar{K}]_{T=0}^{Q=1} : S^1(3140), (2)$ in which a peak corresponding to a total mass of $M = 3141 \pm 6$ MeV/$c^2$ with a total binding energy of $S_K + BE(^4\text{He}) = -E_K \equiv -(M - 2 M_p - M_n - M_{K^-})c^2 = 169 \pm 6$ MeV and a width less than 25 MeV was revealed. This can be identified as the originally predicted $T = 0$ $ppnK^-$. Surprisingly, the observed total binding energies of both $S^0(3115)$ and $S^1(3140)$ (194 and 169 MeV, respectively) are much larger than the predicted ones (97 and 118 MeV, respectively). Furthermore, the former $T = 1$ state lies below the latter $T = 0$ state, contrary to a naive expectation.

In the present paper we show that these surprising observations can be understood within the framework of deeply bound $\bar{K}$ nuclei.

2. Spin-isospin structure of strange tribaryons

Let us first discuss what kinds of states are expected to be low lying in the strange tribaryon system. The nomenclature we adopt, $[N'NN\bar{K}]_{(T,T_3)}^Q$, with $Q$ being a charge and $(T, T_3)$ a total isospin and its 3rd component, persists no matter whether the constituent $\bar{K}$ keeps its identity or not. We also use a conventional charge-state configuration, such as $ppnK^-$, for representing an isospin configuration without any loss of generality. The first two nucleons occupy the ground orbital (0s), whereas the third nucleon ($N'$) in the case of $T_{NN} = 3/2$ has to occupy an excited orbital (the lowest one is $0p_{3/2}$). An overall view of the tribaryon system together with the experimental information is presented in Fig. [1].

Intuitively speaking, the level ordering depends on the number of strongly attractive $I = 0 \bar{K}N$ pairs in each state. Thus, it is instructive to count the projected number of pairs in the $\bar{K}$-nucleus interaction in each state. The originally predicted $T = 0$ state ($ppnK^-$) has partial isospins of $T_{NN} = 1$ and $T_{NN'} = 1/2$ and spin and parity (including that of $\bar{K}$) of $J^*=1/2^-$, in which the attractive interaction is represented by

$$V_{T=0}^{(T=0)} = \frac{3}{2}v^{I=0}. \quad (3)$$

In this case, the bound “nucleus” is a shrunk $(nnn)_T=1/2$. For $T = 1$, on the other hand, there are four possible configurations, and the most favorite state is a linear combination of two configurations with $T_{NN'} = 1, T_{NN} = 1, T_{NN'N} = 3/2$ and with $T_{NN'} = 0, T_{NN} = 1, T_{NN'N} = 1/2$, in which the attractive interaction is represented by

$$V_{T=1}^{(T=1)} = \frac{2}{3}v^{I=0} + \frac{4}{3}v^{I=0}. \quad (4)$$

In the case that $v^{I=0} \sim v^{I=0}$, the attractive interaction for the $T = 1$ state amounts to $\sim 2v^{I=0}$, which is larger than that for the $T = 0$ state. This is a key to understanding the level ordering of the $T = 1$ and $T = 0$ states; in the former, the stronger $\bar{K}$-core interaction tends to compensate for the large internal energy of the $T_{NN'N} = 3/2$ core.

The predicted $pppK^-$ state corresponds to this lowest $T = 1$ state. The “core nucleus” that is bound by $K^-$ is not at all like a triton ([nnn])$_T=1/2$, but is close to a non-existing $ppp$. Figure [1] shows an overview of the lowest $T = 1$ (isospin triplet) and $T = 0$ states, together with the originally predicted energy levels and density distributions of $[pppK^-]_{T=0}$ and $[pppK^-]_{T=1}$ (upper part) and the observed energy levels in this framework (lower part). Now that $[NN\bar{K}]_{(T,T_3)=(1,-1)}^{Q=0}$ has been observed as $S^0(3115)$, another isospin partner $S^1_{(T,T_3)=(1,0)}^{Q=1}$ should also exist, and is expected to appear in a spectrum of $^4\text{He(stopped--$K^-$, $n$)}$ with a marginal strength [3].

It should be noted that the larger number of attractive $\bar{K}N^{I=0}$ in the $T = 1$ state may cause a lowering of the $T = 1$ state, even below the $T = 0$ state, although the third nucleon in the $T = 1$ state should be flipped up to the excited orbital ($0p_{3/2}$). In the following we discuss this possibility by addressing the following questions: 1) the nuclear compression, 2) the relativistic effect, 3) the spin-orbit interaction, and 4) the
Thus, to avoid confusion, it is convenient to divide the total $K$ potential ("separating" potential) into the core part and a $K$-core "binding" potential as:

$$U_K(r) = \Delta E_{\text{core}}(r) + U_{\bar{K} - \text{core}}(r). \quad (5)$$

We distinguish between the separation energy of $\bar{K}$ ($S_K$) and the $\bar{K}$ binding energy ($B_K \equiv -E_{\bar{K} - \text{core}}$):

$$-S_K = \langle \Delta E_{\text{core}} \rangle + E_{\bar{K} - \text{core}},$$

where $\langle \Delta E_{\text{core}} \rangle$ is an expectation value of the core compression energy with respect to $\bar{K}$ distribution. The calculated shrunk-core energy ($\Delta E_{\text{core}}$), $\bar{K}$-core potentials ($U_{\bar{K} - \text{core}}$) and $\bar{K}$-core binding energies ($-E_{\bar{K} - \text{core}}$) in the $T = 0$ state are shown in Fig. 2.

The relativistic effect can be taken into account by using a linearized Klein-Gordon (KG) equation for $\bar{K}$,

$$\left\{-\frac{\hbar^2}{2m_{\bar{K}}} \nabla^2 + U_{\bar{K} - \text{core}}\right\} |\Phi\rangle = \left(\varepsilon_{\text{KG}} + \frac{\varepsilon_{\text{KG}}^2}{2m_{\bar{K}} c^2}\right) |\Phi\rangle,$$  \quad (7)

where $\varepsilon_{\text{KG}}$ ($\equiv E_{\bar{K} - \text{core}}$) is the energy of $\bar{K}$ without its rest mass energy, and $m_{\bar{K}}$ the rest mass of $\bar{K}$. The optical potential, $U_{\bar{K} - \text{core}}$, is given on the assumption that $\bar{K}$ is in a scalar mean-field potential as:

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provided by the shrunk nuclear core. When we make a transformation of the KG energy as

\[
\left( \varepsilon_{KG} + \frac{\varepsilon_{KG}^2}{2m_Kc^2} \right) \rightarrow \varepsilon_S.
\]

Eq. (6) becomes equivalent to a Schrödinger-type equation with an energy solution of \( \varepsilon_S \). Thus, the KG energy can be estimated from a Schrödinger solution, which we obtain in the NR calculation, by using

\[
\varepsilon_{KG} = m_Kc^2 \left( \sqrt{1 + \frac{2\varepsilon_S}{m_Kc^2}} - 1 \right). \tag{9}
\]

This “exact” relation means that, when the Schrödinger energy \( \varepsilon_S \) drops down to \(-m_Kc^2/2\), and the relativistic energy becomes \(-m_Kc^2\), namely, the total mass becomes 0 (“kaon condensation” regime), as shown in Fig. 2.

In this “mapping” treatment, we also made consistent corrections on the threshold energies of \( \Lambda + \pi \) and \( \Sigma + \pi \) and the complex energy of \( \Lambda_{1405} \), and obtained re-fitted \( \bar{K}N \) interaction parameters. It is to be noted that, since the internal energy of the shrunk nucleus is so large, the \( U_{\bar{K}-\text{core}} \) for the KG equation is very deep. Thus, the relativistic treatment gives a substantial negative correction to the energy \( E_{\bar{K}-\text{core}} \).

Let us consider the case of \( [NNN\bar{K}]_{T=0} \). The original NR total binding energy, \( E_K = -118 \text{ MeV} \), was readjusted to \(-111 \text{ MeV} \) after taking into account the relativistic effect on the \( \bar{K}N \) binding in \( \Lambda_{1405} \). The nuclear core energy from the core shrinkage is \(< \Delta E_{\text{core}} > \approx 50 \text{ MeV} \). Thus, we obtain \( V_0 = -570 \text{ MeV} \), \( W_0 = 18 \text{ MeV} \) and \( a_K = 0.923 \text{ fm} \) in the expression for the \( \bar{K} \)-core potential as

\[
U_{\bar{K}-\text{core}}(r) = (V_0 + iW_0) \exp \left[ - \left( \frac{r}{a_K} \right)^2 \right]. \tag{10}
\]

In this case, the relativistic correction is \(-23 \text{ MeV} \), yielding \( E_K = -134 \text{ MeV} \). The total binding energy is still smaller by 35 MeV than the experimental value, 169 MeV.

Next, we consider the case of \( [NNN\bar{K}]_{T=1} \). This state has a larger \(< \Delta E_{\text{core}} > \) than the \( T = 0 \) state, because the energy to excite one nucleon from the \( 0s \) shell to the \( 0p \) is estimated
to be 50 MeV. Thus, starting from the NR result, we obtain a deeper $\bar{K}$-core potential, $V_0 = -652$ MeV and $W_0 = -12$ MeV, and the $\bar{K}$ binding energy is subject to a large relativistic correction, $\Delta \varepsilon_{RC} = -46$ MeV.

Roughly speaking, the relativistic effect accounts for about half of the discrepancies in $S_K$. It is to be noted that this large correction is a consequence of a shrinkage of the nuclear core. Namely, the large $<\Delta E_{\text{core}}>$(compression energy), which translates into a larger negative $\bar{K}$ potential ($U_{\bar{K}-\text{core}}$), causes a larger relativistic correction. However, the resulting total binding energy (143 MeV) is still smaller than the observed one (194 MeV).

4. Spin-orbit splitting

In the $T = 1$ state the third nucleon occupies a $0p_{3/2}$ orbital and behaves like a compact satellite halo [5], as shown in Fig. 1. In our previous prediction we neglected the spin-orbit splitting between $0p_{3/2}$ and $0p_{1/2}$. Here, we note that the one-body spin-orbit interactions may give a large contribution for a shrunk nucleus, because it depends on the gradient of the nuclear surface. To estimate the effect of the one-body spin-orbit interaction, we used the well known Thomas-type $l \cdot s$ potential,

$$V_{ls}(r) = -(i\hbar^2/2M)\frac{2}{dr}\left[1 + \frac{U_{\text{nucl}}(r)}{2Mc^2}\right], \quad (11)$$

and found that $\Delta E(J^\pi = 3/2^+) \simeq -5$ MeV and $\Delta E(J^\pi = 1/2^+) \simeq 10$ MeV in the case of a shrunk $0p$ orbital.

A further contribution is expected from the nucleon-nucleon two-body spin-orbit interaction ($NN \ LS$ interaction), because it is known to be attractive enough to cause the $^3P_2$ pairing in dense neutron matter [9]. We calculated the expectation value of the sum of the $NN \ LS$ interaction among the nucleons. Here, we used the effective $LS$ interaction derived from the Tamagaki potential (OPEG) with the $g$-matrix method, similarly to our previous studies [14][15]. Fig. 4 shows the behavior of the $NN \ LS$ contribution in the $pppK^-$ when the rms radius ($R_{\text{rms}}$) is varied. The contribution of the $NN \ LS$ interaction increases rapidly as the system becomes small. The magnitude of $\Delta E(J^\pi = 3/2^+)$ is found to increase to $\sim 15$ MeV in the shrunk system, whereas it is only $\sim 1$ MeV in the normal-size nuclei.

Thus, the spin-orbit interactions of both kinds make the $T = 1$ energy even lower. It is important to find a spin-orbit partner, $J^\pi = 1/2^+$, which will give an experimental value of the spin-orbit splitting in such a dense nuclear system. From this one can obtain information about the size of the shrunk $T = 1$ state. According to our calculation, the spin-orbit splitting energy is $E(1/2^+) - E(3/2^+) = -3 \Delta E(3/2^+) \sim 60$ MeV.

5. Medium-modified $\bar{K}N$ interaction

There are still large discrepancies in $S_K$ between theory and observation, even after a relativistic correction. They can now be ascribed to a medium-modified bare $\bar{K}N$ interaction that may occur in such a dense nuclear medium. In the present case, the average nucleon density, $<\rho(r)> \sim 3 \times \rho_0$, approaches the nucleon compaction limit ($\rho_0 \sim 2.7 \rho_0$), where a chiral symmetry restoration may occur. Similar to the case
of the observed pionic bound states [10,11], the $\bar{K}N$ interaction is related to the order parameter of the quark condensate and is expected to be enhanced as chiral symmetry is restored. Thus, we tuned the bare $\bar{K}N$ strength by a small factor, and recalculated the total binding energies to find the most suitable enhancement parameter. Since this modification causes a change in the relativistic correction, we iterated all of these corrections consistently. The following enhancement was found to account for both $S^0(3115)$ and $S^1(3140)$ simultaneously:

$$\frac{\bar{K}N}{K^N_{\text{bare}}} \sim 1.15.$$  \hfill (12)

The final results after this tuning are also presented in Fig. 3. The $\bar{K}$-core potential strength, $U_{\bar{K}-\text{core}}$, is now $-618-i11$ MeV with $a_K = 0.920$ fm for the $T = 0$ state.

Using the enhanced $\bar{K}N$ interaction strength and also taking into account the relativistic effect, we recalculated the binding energy and width of the most basic $\bar{K}$ nuclear system, $ppK^-$. The results are $M = 2284$ MeV/$c^2$ ($S_K = 86$ MeV) and $\Gamma = 58$ MeV, in contrast to the original non-relativistic values, $M = 2322$ MeV/$c^2$ ($S_K = 48$ MeV) and $\Gamma = 61$ MeV. It is important to find this state experimentally so as to establish a solid starting gauge for more complex $\bar{K}$ bound systems.

6. Energy difference among the isotriplet states

Although the observed $T = 0$ and $T = 1$ states support the theoretical expectation for nuclear shrinkage, a direct experimental verification, if possible, would be vitally important. We examine the energy difference of the isobaric analog states of the isotriplet, which is related to the strong-interaction mass term and the Coulomb displacement energy as

$$\Delta E_{\text{sum}}(T_3) = \Delta E_{\text{mass}}(T_3) + E_C(T_3).$$  \hfill (13)

We estimated the mass term, $\Delta E_{\text{mass}}$, as the deviation of the sum of the constituent particle masses ($p, n, K_0$ and $K^-$) from the central value by using the calculated isospin eigenstates in terms of the charge states (the $(T, T_3) = (1, 1)$ and $(1, -1)$ isospin eigenstates have mixtures of 70% $pppK^-$ + 30% $ppnK_0$ and of 30% $pnnK^- + 70% nnnK_0$, respectively). The values of $\Delta E(T_3)$ are listed in Table 1.

| $(T, T_3)$ | $(1, 1)$ | $(1, 0)$ | $(1, -1)$ |
|------------|----------|----------|-----------|
| $Q$ Charge states | $pppK^-$ | $ppnK^-$ | $pnnK^-$ |
| $E_{\text{mass}}$ | $-2.4$ | $-0.7$ | $2.4$ |
| $E_C(\text{total})$ | $-0.6$ | $-1.3$ | $-0.7$ |
| $E_C(\bar{K}N)$ | $3.7$ | $1.1$ | $0$ |
| $E_{\text{sum}}$ | $-4.3$ | $-2.4$ | $-0.7$ |

For estimating $E_C$ we calculated the Coulomb energy of each particle pair using the total wavefunction. The results, presented in Table 1, show that the Coulomb energies of the $NN$ and $\bar{K}N$ pairs are 0 and $-0.7$ MeV, respectively, for the $T_3 = -1$ state, whereas they increase in magnitude to 3.7 and $-4.3$ MeV, respectively, in the $T_3 = 1$ state, which are, however, nearly cancelled by each other. Thus, the total Coulomb energies for the two isospin states remain nearly zero. On the other hand, a naive estimate of the Coulomb energy assuming a uniformly charged sphere for a fictitious $pppK^-$ would give $E_C(T_3 = 1) - E_C(T_3 = -1) \sim (3/5)Q^2e^2/R \sim 2.1$ MeV, if one takes $R_{\text{rms}} = 1.61$ fm, the ordinary nuclear radius for $A = 3$. The different cases can be distinguished experimentally, if the two (or three) isobaric analog states are produced and identified. In the next section we propose some experimental methods.

8. Role of $\Lambda_{1405}$ in the $S^0$ population

In the $^4\text{He}(\text{stopped-}K^-, p)$ reaction with
Auger-proton emission, the three nucleons in the target $^4$He are expected to remain in the $0s$ orbital. Then, why can the $T = 1$ state with a shrunk core of $T = 3/2$ be populated? The key to understand this process is the role of $\Lambda^*$ ($\equiv \Lambda^*$) as a doorway; the formation of $\Lambda^*$ in the $K^-$ absorption at rest by $^4$He is known to occur with a substantial branching ratio [12]. This doorway state can lead to core excited $K^-$ states:

$$K^- + "p" \rightarrow \Lambda^*,$$  
$$\Lambda^* + "pnn" \rightarrow [(pnn)_{T=3/2}K^-]_{T=1} + p,$$  
$$\rightarrow [(pnn)_{T=3/2}K^-]_{T=1} + n,$$

where the proton from $\Lambda^* = pK^-$ falls onto the $0p$ orbital, as shown in Fig. 5. Likewise, the doorway $\Lambda^*$ leads to many other $K$ bound states, such as $\Lambda^*p \rightarrow ppK^-$, as emphasized in [21,13].

9. Future experiments

A direct reaction to produce $pppK^-$ via $^3$He($K^-,\pi^-$) and $^3$He($\pi^+,K^+$) was proposed in [2]. Spectral functions (“effective nucleon numbers”) in the missing mass, as shown in Fig. 6, were calculated based on the $A_{1405}$ doorway model. The spectrum shows the spin-orbit pair ($J^* = 3/2^+$ and $1/2^+$) states with a calculated splitting of 60 MeV. Such a pair, if observed, will give important information on the size of the system.

Another way to produce $pppK^-$ is to use a cascade reaction in a light target (say, $^4$He), such as

$$p + "n" \rightarrow \Lambda^* + K^0 + p,$$  
$$\Lambda^* + "pnn" \rightarrow pppK^- + n.$$  

In the second process, the energetic “doorway particle”, $A_{1405}$, produced by an incident proton of sufficiently large kinetic energy, knocks on one of the remaining nucleons and/or hits the remaining nucleus to form a kaonic bound system (“$K$-transfer” reactions). $\Lambda^*$ compound processes induced by ($K^-,\pi^-$) and ($\pi^+,K^+$) may also produce kaonic systems. Recently, it has been pointed out that a fireball in heavy-ion collisions can be a source for kaonic systems [3]. In all of these reactions one can identify a $K$ cluster by invariant-mass spectroscopy following its decay, such as

$$pppK^- \rightarrow p + p + \Lambda.$$  

10. Concluding remarks

In the present paper we have shown that the observed strange triaryons, $S^0(3115)$ and $S^1(3140)$, can be understood as the $(T, T_3) = (1, -1)$ and $T = 0$ $N'NN\bar{K}$ bound states with shrunk nuclear cores of $T = 3/2, J = 3/2$ and $T = 1/2, J = 1/2$, respectively. The fact that $S^0(3115)$ lies below $S^1(3140)$ strongly supports the prediction that the three nucleons are in a “non-existing nucleus” ($N'NN\bar{K}$) with which the attractive $I = 0$ $KN$ attraction is maximal. The spin-orbit splitting, enhanced in a condensed nucleus, helps to further lower the $T = 1$ state. The observed binding energies, which are substantially larger than the predicted non-relativistic values, are partially accounted for by the relativistic effect on the $\bar{K}$ and partially by invoking an enhanced bare $KN$ interaction. The enhancement may indicate a partial restoration
of chiral symmetry and/or a transition to a 11-quark-gluon phase. The observed deep $\bar{K}$ binding indicates that the system is approaching the kaon condensation regime [14,15]. These discoveries have demonstrated that narrow deeply bound states of $\bar{K}$ exist, as we have predicted, in contrast to the prevailing belief and claim for a shallow $\bar{K}$ potential [16,17], which may apply to unbound continuum systems.

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