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Strange tribaryons as $\bar{K}$-mediated dense nuclear systems

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Abstract

We discuss the implications of recently discovered strange tribaryons in ^{4}He(stopped-K^-, p) S^0(3115) and ^{4}He(stopped-K^-, n) S^1(3140) within the framework of deeply bound $\bar{K}$ states formed on shrunk 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 [1–6], 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 Λ_{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, ppK^- with a total binding energy of $-E_K = 48$ MeV, $ppnK^-$ with $-E_K = 118$ MeV, and $pppK^-$ with $-E_K = 97$ MeV. The calculated rms distances in the $ppnK^-$ system are: $R_{N-N} = 1.50$ fm and $R_{\bar{K}-N} = 1.19$ fm, whereas $R_{\bar{K}-N} = 1.31$ fm in Λ_{1405}. The NN distance in $ppnK^-$ is substantially smaller than the normal inter-nucleon distance (~ 2.2 fm), and the average nucleon density, $\rho_N^{\text{ave}} = 3.1 \times \rho_0$, is much larger than

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the normal density, $\rho_0 \sim 0.17\ \mathrm{fm}^{-3}$. 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 bench marks to examine the above problems.

An experimental search for the $ppnK^-$ in the ^{4}He(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],

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

where a distinct narrow peak appeared at a mass of $M = 3117 \pm 5\ \mathrm{MeV}/c^2$ with a total binding energy (separation energy of K^-) of $S_K = -E_K \equiv -(M - M_p - 2M_n - M_{K^-})c^2 = 194 \pm 5$ 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 [1]. On the other hand, an exotic $T = 1$ $pppK^-$ state with $S_K = 97$ MeV had been previously predicted [2,5], and the observed $\mathrm{S}^0(3115)$ can be identified as its isobaric analog state.

An indication of another species was observed in the neutron-emission channel [8],

$$^4\mathrm{He}(\text{stopped-}K^-, n)[N'NN\bar{K}]^{Q=1}_{T=0}\text{: } \mathrm{S}^1(3140), \quad (2)$$

in which a peak corresponding to a total mass of $M = 3141 \pm 6\ \mathrm{MeV}/c^2$ with a total binding energy of $S_K + BE(^3\mathrm{He}) = -E_K \equiv -(M - 2M_p - M_n - M_{K^-})c^2 = 169 \pm 6$ 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 $\mathrm{S}^0(3115)$ and $\mathrm{S}^1(3140)$ (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 Letter 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, $[N'NN\bar{K}]^Q_{(T,T_3)}$, 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 $T_{N'NN} = 3/2$ has to occupy an excited orbital (the lowest one is $0\mathrm{p}_{3/2}$). 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 $T_{N'NN} = 1/2$ and spin and parity (including that of $\bar{K}$) of $J^\pi = 1/2^-$, in which the attractive interaction is represented by

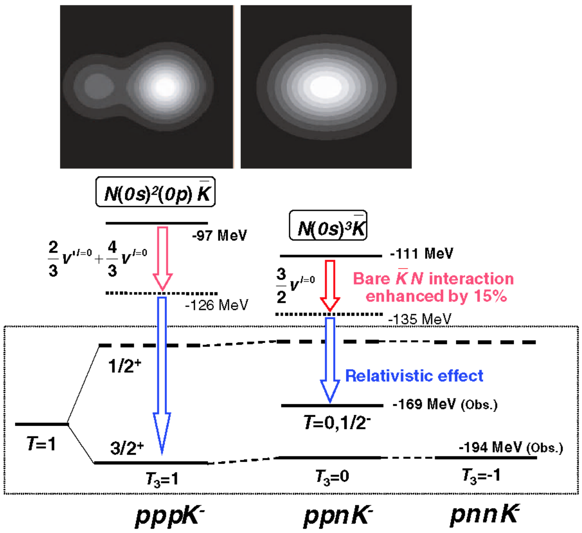

Fig. 1. Spin–isospin structure of the strange tribaryon system $[(NNN)\bar{K}]^Q_{(T,T_3)}$. The previously calculated nucleon-density contours and energy levels with E_K values are shown on top. The observed $\mathrm{S}^0(3115)$ and $\mathrm{S}^1(3140)$ are identified as the $(T, T_3) = (1, -1)$ and $T = 0$ states, respectively.

$$V^{(T=0)}_{N'NN-\bar{K}} = \frac{3}{2} v^{I=0}. \tag{3}$$

In this case, the bound "nucleus" is a shrunk $(ppn)_{T=1/2}$. 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 $T_{\bar{K}N'} = 1$, $T_{NN} = 1$, $T_{N'NN} = 3/2$ and with $T_{\bar{K}N'} = 0$, $T_{NN} = 1$, $T_{N'NN} = 1/2$, in which the attractive interaction is represented by

$$V^{(T=1)}_{N'NN-\bar{K}} = \frac{2}{3} v^{I=0'} + \frac{4}{3} v^{I=0}. \tag{4}$$

In the case that $v^{I=0'} \sim v^{I=0}$, the attractive interaction for the $T = 1$ state amounts to $\sim 2\, v^{I=0}$, 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 $T_{N'NN} = 3/2$ 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 ($[pnn]_{T=1/2}$), but is close to a non-existing ppp. Fig. 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 $[ppnK^-]_{T=0}$ and $[pppK^-]_{T=1}$ (upper part) and the observed energy levels in this framework (lower part). Now that $[NNN\bar{K}]^{Q=0}_{(T,T_3)=(1,-1)}$ has been observed as $\mathrm{S}^0(3115)$, another isospin partner $\mathrm{S}^{Q=1}_{(T,T_3)=(1,0)}$ should also exist, and is expected to appear in a spectrum of ^{4}He(stopped-K^-, n) with a marginal strength [8].

It should be noted that the larger number of the attractive $\bar{K}N^{I=0}$ pairs 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 ($0\mathrm{p}_{3/2}$). 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 possibility of a medium modification of the bare $\bar{K}N$ interactions.

3. Relativistic effect on $\bar{K}$ binding

The calculations so far made were based on a non-relativistic (NR) treatment of many-body systems. For very deeply bound $\bar{K}$, however, relativistic corrections are indispensable. In this respect it is important to recognize that the $\bar{K}$ bound system is a very peculiar one in which the $\bar{K}$ is bound by a non-existing fictitious nucleus, namely, a *shrunk nuclear core* with a large internal energy (compression energy, ΔE_{core}). Thus, to avoid confusion, it is convenient to divide the total $\bar{K}$ potential ("separating" potential) into the core part and a $\bar{K}$-core "binding" potential as:

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

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

$$-S_K = \langle \Delta E_{\mathrm{core}} \rangle + E_{\bar{K}\text{-core}}, \tag{6}$$

where $\langle \Delta E_{\mathrm{core}} \rangle$ is an expectation value of the core compression energy with respect to $\bar{K}$ distribution. The calculated shrunk-core energy ($\Delta E_{\mathrm{core}}(r)$), $\bar{K}$-core potentials ($U_{\bar{K}\text{-core}}(r)$) and $\bar{K}$-core binding energies ($-E_{\bar{K}\text{-core}}$) in the $T = 0$ state are shown in Fig. 3.

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

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

where $\varepsilon_{\mathrm{KG}}$ ($\equiv E_{\bar{K}\text{-core}}$) is the energy of $\bar{K}$ without its rest-mass energy, and m_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 *provided by the shrunk nuclear core*. When we make a transformation of the KG energy as

$$\left(\varepsilon_{\mathrm{KG}} + \frac{\varepsilon^2_{\mathrm{KG}}}{2 m_K c^2} \right) \to \varepsilon_{\mathrm{S}}. \tag{8}$$

Eq. (7) becomes equivalent to a Schrödinger-type equation with an energy solution of ε_{S}. Thus, the KG energy can be estimated from a Schrödinger solution, which we obtain in the NR calculation, by using

$$\varepsilon_{\mathrm{KG}} = m_K c^2 \left(\sqrt{1 + \frac{2\,\varepsilon_{\mathrm{S}}}{m_K c^2}} - 1 \right). \tag{9}$$

This "exact" relation means that, when the Schrödinger energy (ε_{S}) drops down to $-m_K c^2/2$, the relativistic energy becomes $-m_K c^2$, 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 Λ_{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 $E_{\bar{K}\text{-core}}$.

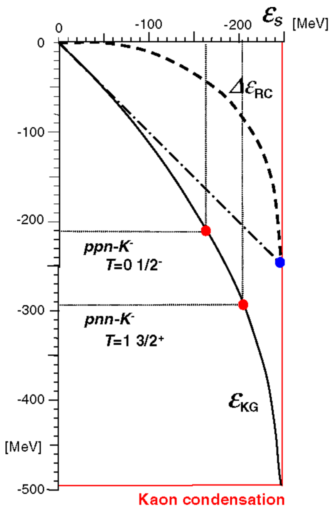

Fig. 2. Relation between the Schrödinger energy (ε_S) and the KG energy (ε_{KG}). $\Delta\varepsilon_{RC} \equiv \varepsilon_{KG} - \varepsilon_S$ is the relativistic correction.

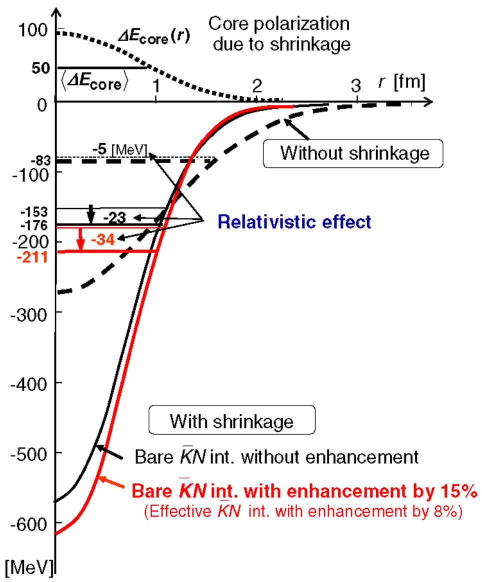

Fig. 3. Calculated $\bar{K}$-core potentials, $U_{\bar{K}\text{-core}}(r)$, for the $T = 0$ $ppnK^-$ state in three cases: without shrinkage, with shrinkage with the original bare $\bar{K}N$ interaction and with an enhanced bare $\bar{K}N$ interaction by 15%. The relativistic effects on the $\bar{K}$-core binding energies ($E_{\bar{K}\text{-core}}$) are indicated by arrows. Also shown is the nuclear core energy, $\Delta E_{\text{core}}(r)$, with an average value ($\langle \Delta E_{\text{core}} \rangle$).

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

$$U_{\bar{K}\text{-core}}(r) = (V_0 + i\,W_0) \exp\left[-\left(\frac{r}{a_K}\right)^2\right]. \qquad (10)$$

In this case, the relativistic correction is -23 MeV, yielding $E_K = -134$ 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 $\langle \Delta E_{\text{core}} \rangle$ 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 $\langle \Delta E_{\text{core}} \rangle$ (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

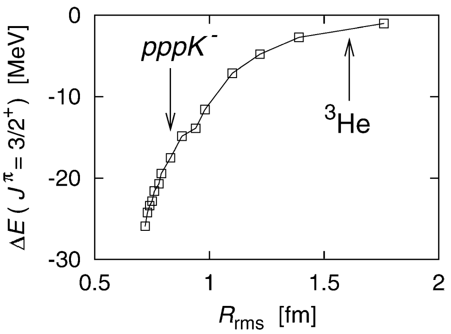

Fig. 4. Spin–orbit energy, $\Delta E(J^\pi = 3/2^+)$, in $pppK^-$ ($J^\pi = 3/2^+, T = 1$) from the NN–LS contribution as a function of $R_{\rm rms}$, the root-mean-square radius of the nuclear system.

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 $\boldsymbol{l} \cdot \boldsymbol{s}$ potential,

$$V_{ls}(r) = -(\vec{l}\vec{s})\frac{\hbar^2}{2M}\frac{2}{r}\frac{\mathrm{d}}{\mathrm{d}r}\ln\left[1 + \frac{U_{\rm nucl}(r)}{2Mc^2}\right], \qquad (11)$$

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 [1,3–5]. Fig. 4 shows the behavior of the NN–LS contribution in the $pppK^-$ when the rms radius ($R_{\rm rms}$) 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 ~ 15 MeV in the shrunk system, whereas it is only ~ 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, $\langle\rho(r)\rangle \sim 3 \times \rho_0$, approaches the nucleon compaction limit ($\rho_{\rm c} \sim 2.7\,\rho_0$), 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:

$$\frac{\bar{K}N}{\bar{K}N^{\rm bare}} \sim 1.15. \qquad (12)$$

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 $-618 - i\,11$ 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 $M = 2284\ \mathrm{MeV}/c^2$ ($S_K = 86$ MeV) and $\Gamma = 58$ MeV, in contrast to the original non-relativistic values, $M = 2322\ \mathrm{MeV}/c^2$ ($S_K = 48$ MeV) and $\Gamma = 61$ MeV. It is important to find this state (or equivalently, $(pn)_{T=1}K^-$) 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,

Table 1
Calculated energy differences (in MeV) in the $(T, T_3) = (1, \pm 1)$ states of $NNN\bar{K}$. E_C: Coulomb energy; ΔE_{mass}: mass difference of the constituent particles; ΔE_{sum}: total energy difference

(T, T_3)	$(1, 1)$	$(1, -1)$
Q	2	0
Charge	70% $pppK^-$	31% $pnnK^-$
States	30% $ppn\bar{K}^0$	69% $nnn\bar{K}^0$
E_C(total)	−0.6	−0.7
$E_C(NN)$	3.7	0
$E_C(\bar{K}N)$	−4.3	−0.7
ΔE_{mass}	−2.4	2.4
ΔE_{sum}	−3.5	1.7

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

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

The results of calculation are summarized in Table 1. The charge-state configuration for each isospin state is also shown.

First, we calculated the Coulomb energy, $E_C(T_3)$, for each particle pair using the total wavefunction. 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 total energy differences, $\Delta E_{\text{sum}}(T_3)$, are shown in the table. As a reference, a naive estimate of ΔE_{mass} as the deviation of the sum of the constituent particle masses $(p, n, \bar{K}_0$ and $K^-)$ from the central value is also shown.

The case of the condensed $T = 1$ $NNN\bar{K}$ can be distinguished experimentally from the case of a conventional system with the ordinary density, if the two (or three) isobaric analog states are produced and identified. In the next section we propose some experimental methods.

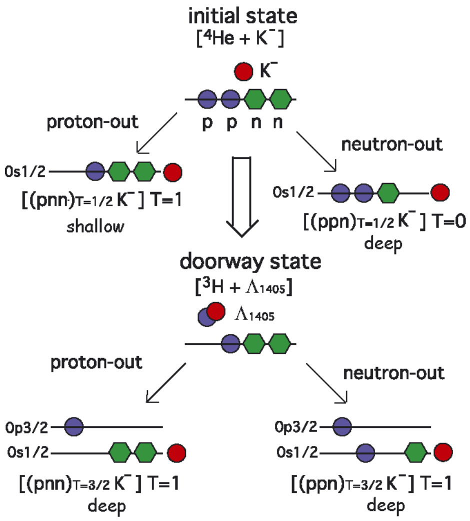

Fig. 5. Population mechanism of the $[N'NN\bar{K}]_{T=1}$ through the Λ_{1405} doorway in ^{4}He(stopped-K^-, x).

7. Role of Λ_{1405} in the S^0 population

In the ^{4}He(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 Λ_{1405} ($\equiv \Lambda^*$) as a doorway; the formation of Λ^* 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 $\bar{K}$ states:

$$K^- + \text{“}p\text{”} \to \Lambda^*, \tag{14}$$

$$\Lambda^* + \text{“}pnn\text{”} \to [(pnn)_{T=3/2}K^-]_{T=1} + p, \tag{15}$$

$$\to [(ppn)_{T=3/2}K^-]_{T=1} + n, \tag{16}$$

where the proton from $\Lambda^* = pK^-$ falls onto the 0p orbital, as shown in Fig. 5. Likewise, the doorway Λ^* leads to many other $\bar{K}$ bound states, such as $\Lambda^* p \to ppK^-$, as emphasized in [2,13].

8. Future experiments

A direct reaction to produce $pppK^-$ via ^{3}He(K^-, π^-) and ^{3}He(π^+, K^+) was proposed in [2]. Spectral functions (“effective nucleon numbers”) in the missing mass, as shown in Fig. 6, were calculated based on the Λ_{1405} doorway model. The spectrum shows

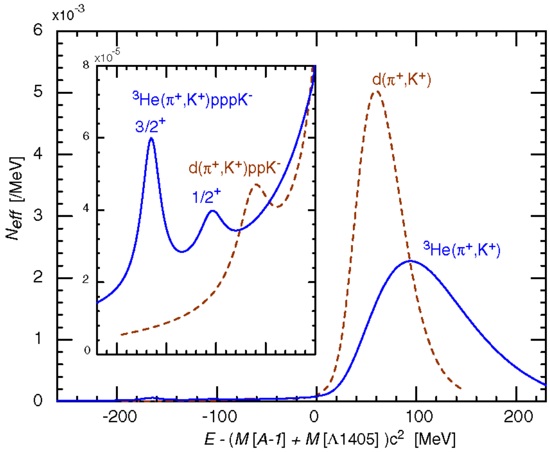

Fig. 6. Spectra (N_{eff}) of (π^+, K^+) reactions at $p_\pi = 1.5$ GeV$/c$ on d and ^{3}He as functions of $E_K - [M(A-1) + M(\Lambda_{1405})]c^2$, calculated based on the Λ^* doorway model. Not only the $J^\pi = 3/2^+$ state, but also its spin–orbit partner $(1/2^+)$ with a splitting of 60 MeV are incorporated.

the spin–orbit pair ($J^\pi = 3/2^+$ 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

$$p + \text{“}n\text{”} \to \Lambda^* + K^0 + p, \tag{17}$$

$$\Lambda^* + \text{“}ppn\text{”} \to pppK^- + n. \tag{18}$$

In the second process, the energetic “doorway particle”, Λ_{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 (“$\bar{K}$-transfer” reactions). Λ^* compound processes induced by (K^-, π^-) and (π^+, 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 [6]. In all of these reactions one can identify a $\bar{K}$ cluster by invariant-mass spectroscopy following its decay, such as

$$pppK^- \to p + p + \Lambda. \tag{19}$$

Once a $\bar{K}$ cluster is identified in a missing-mass spectrum, the momentum correlation of its decay particles can be used to determine the size of the system.

9. Concluding remarks

In the present Letter we have shown that the observed strange tribaryons, 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” $(N'NN)_{T=3/2,J=3/2}$ with which the attractive $I = 0$ $\bar{K}N$ 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 $\bar{K}N$ 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 was obtained by inappropriately applying Lutz’s procedure [18] to deeply bound discrete states, where the $\bar{K}$ self-mass should be introduced not only in the intermediate energies but also in the starting energy of the g-matrix equation (a natural extension of the Bethe–Goldstone linked-cluster expansion).

Note added in proof

We have noticed two theoretical papers concerning the strange tribaryon systems. One is based on a 9-quark model by Maezawa, Hatsuda and Sasaki [19]. Another paper by Wycech and Green [20] emphasizes the coupling of $\Sigma(1385)$.

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