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## Possibility of Charmed Hypernuclei

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We suggest that both two-body and many-body bound states of a charmed baryon and nucleons should exist. Estimates indicate binding in the $^1S_0$ state of $C_1N$ ($I=\frac{3}{2}$) and $SN$ ($I=1$). We further estimate the binding energy of $C_0, C_1$ in various finite nuclei.

Recent experiments[1] have established the existence of a new class of mesons and baryons possessing net charm. In particular, there is firm evidence for the charmed baryons $B_c = C_0$, $C_1$ ($\Lambda_c$ and $\Sigma_c$) and the charmed mesons, $D, D^*$. It is hence of fundamental interest to establish the nature of the interactions of these charmed particles with more familiar hadrons.

In this Letter we ask whether the charmed baryons will bind to a single nucleon or to finite nuclei, producing charmed analogs of the deuteron or of hypernuclei. The estimated short lifetime of even the lowest-mass charmed baryon, $\tau(C_0) \sim 10^{-11}-10^{-14}$ sec, may make it difficult to establish the existence of such analog bound states. These questions are discussed here in a nonrelativistic framework in which the $B_c$-nucleon interaction is represented by a sum of single-boson-exchange potentials. The $B_c$-nucleus potential is obtained by averaging the $B_c$-$N$ interaction over the nuclear density. The large mass of the charmed baryon alters the dynamics appreciably, resulting in a reduction in the $B_c$ kinetic energy for both two- and many-body systems. Thus, in the $B_c$-nucleus case, where we find potentials comparable to the nucleon-nucleus potential in depth, a rich spectroscopy emerges encompassing many bound levels.

The $B_c$-$N$ potential, $V(r)$, is given by a sum of meson exchange potentials for $r \geq r_c$ but taken phenomenologically as an infinite hard core for $r \leq r_c$. Such models have frequently been applied to nucleon-nucleon ($NN$) scattering in the low-energy region.[2] Recently, a one-boson-exchange model has been constructed[3] which successfully reproduces both low-energy $NN$ and hyperon-nucleon ($\Lambda N$ and $\Sigma N$) data. This model includes the coupling of SU(3) nonets of pseudoscalar, scalar, and vector mesons to the $\frac{1}{2}^+$ baryon octet of SU(3). We use here a natural generalization of this model which considers the coupling of mesons obtained by mixing the $\{15\}$ and $\{1\}$ representations of SU(4) with members of the $\{20\}$ representation of $\frac{1}{2}^+$ baryons. The details are given by Dover, Kahana, and Trueman.[4]

We have restricted out attention to the four lightest singly charmed baryons $C_0$, $C_1$, $A$, and $S$ with isospin $T=0$, 1, $\frac{1}{2}$, and $\frac{1}{2}$, respectively. For the $C_0$ (also called $\Lambda_c$), we take the presumed experimental mass of $M_{C_0} = 2.26$ GeV; for $C_1$ (also called $\Sigma_c$), we take $M_{C_1} = 2.42$ GeV. For $A$ and $S$, we use theoretical estimates[5] of $M_A = 2.47$ GeV and $M_S = 2.57$ GeV. Since only the $C_0$ baryon is stable under strong interactions it is likely that the $C_0$-$N$ and $C_0$-nucleus systems are of greatest interest, although the $C_1$ nuclei are most strongly bound.

Using these potentials we have searched for $B_cN$ bound states by solving the nonrelativistic Schrödinger equation. The meson-nucleon and meson-$B_c$ coupling constants we use are listed in Table II of Ref. 4 and Table I of Ref. 3. We took $r_c = 0.522$ fm for $^1S_0$ states.

We start by considering the $^1S_0$ state of $B_cN$

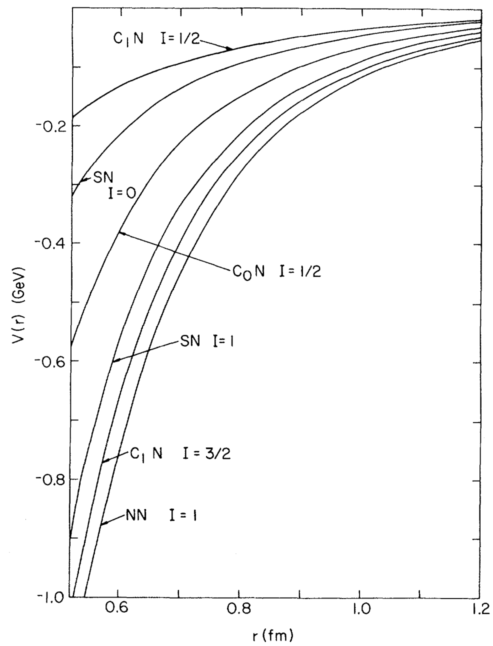


FIG. 1. Two-body charmed-baryon–nucleon potentials in the $^1S_0$ state. The model of Nagels, Rijken, and Swart (Ref. 3) was used, extended to SU(4) symmetry as in Ref. 4.

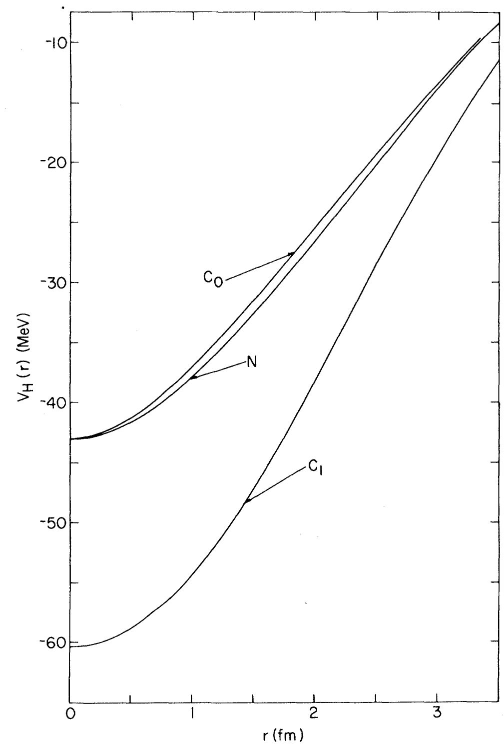


FIG. 2. The Hartree potential $V_H(r)$ for $N$, $C_0$, and $C_1$ in $^{16}O$, calculated according to Eq. (1), using the spin-isospin averaged interaction $v_0(r)$ of Eq. (2).

systems. In analogy to the situation for the deuteron, if a $^1S_0$ $B_cN$ state is bound, we might also expect the corresponding $^3S_1$ state to be bound, as a result of tensor coupling of $^3S_1$ and $^3D_1$. However, the $\pi$-meson component of the latter, proportional to $(m_\pi/M_{B_c})^2$, is weakened and this question must be answered in a separate coupled-channel calculation. The diagonal potentials $V(r)$ for $B_cN \to B_cN$ are shown in Fig. 1 for the $^1S_0$ configuration of $NN$, $C_0N$, $C_1N$, and $SN$. The isovector $\rho$- and $\pi$-exchange potentials are always separately *attractive* for the maximum $I$ and *repulsive* for minimum $I$. The isospin dependence of the $B_cN$ potential is appreciable; thus binding is most likely in states of maximum $I$.

In fact we find that $^1S_0$ $C_1N$ $(I=\frac{3}{2})$ and $SN$ $(I=1)$ states are bound by 1.76 and 0.08 MeV, respectively. The somewhat stronger potential obtained for the unbound $^1S_0$ $NN$ system is here compensated for by the lowered kinetic energy of $B_cN$. The larger reduced mass of $B_cN$ is also responsible, however, for a lowering of the recoil-dependent $\pi$-meson exchange potential. Thus we find that neither the $^3S_1$ $(I=\frac{1}{2})$ $C_1N$ nor $^1S_0$ $C_0N$ states are bound in approximate calculations including tensor forces in the former case and off-diagonal $I=1$ exchange $(C_1N\text{-}C_0N)$ forces in the latter. A slight increase in the scalar meson coupling from $g_\epsilon/\sqrt{4\pi}=5.03$ to 5.23 does produce binding in $^1S_0$ $C_0N$ when coupling to $C_1N$ is approximately included. Our conclusions about the two-body system are model dependent, but it is clear that no such states will ever be strongly bound.

In contrast, the large mass of the charmed baryon should lead to many bound states in the many-body system, some of which will be deeply bound. In particular the spin and isospin averaging which occurs in a finite nucleus leads to strong $C_0$-, $C_1$-nucleus interactions comparable to $N$-nucleus. Thus charmed nuclear states stable against strong decay will exist. For spin-, isospin-saturated nuclei ($J=0$, $I=0$), the Hartree potential is[6]

$$V_H(r)=\int\rho(\vec{r}\,')v_0(\vec{r}-\vec{r}\,')d^3r', \tag{1}$$

neglecting spin-orbit and Coulomb contributions. Only Wigner forces are generated, and these arise from exchange of $\omega$, $\varphi$, $\epsilon$ mesons. The hard core is removed by use of the Moszkowski-Scott separation procedure.[7] This method consists of finding a cutoff radius $r_0$ for which the $^1S_0$ phase shift vanishes near zero kinetic energy. We may

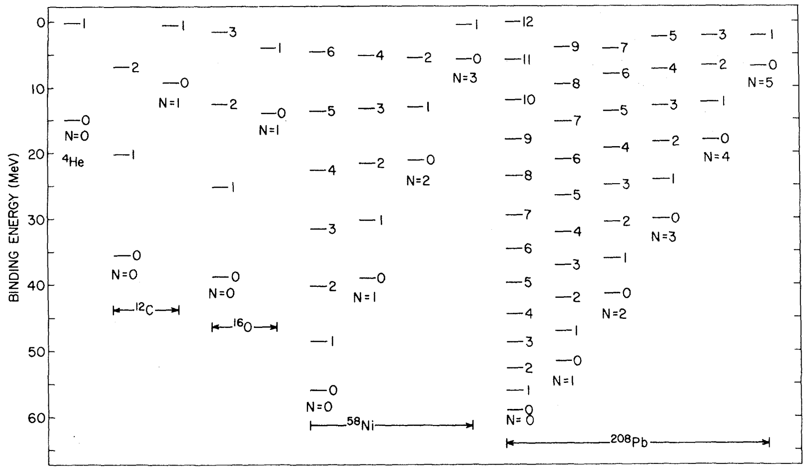


FIG. 3. The charmed hypernuclear spectrum for a $C_1$ bound in various nuclei. The states are labeled by $L$ value; each will be further split into two components by the $\vec{L}\cdot\vec{S}$ potential. Note that some states of high $N$ and $L$ are bound. The spectrum for $C_0$ is very similar in spacing, but fewer states are bound ($L \leqslant 9$ for $N = 0$ in $^{208}$Pb, for example).

then take

$$v(r) = 0, \quad r < r_0, \qquad v(r) = v_0(r), \quad r \geqslant r_0, \tag{2}$$

with

$$v_0(r) = \sum \mu_i g_i^2 e^{-x_i}/x_i. \tag{3}$$

In Eq. (3) $x_i = \mu_i r$ where $\mu_i$ is a meson mass. Also $g_\omega^2/4\pi = 15.29$, 10.31, 11.74 while $g_\varphi^2/4\pi = 1.81$, 1.56, 1.36 for the $NN$, $C_0N$, $C_1N$ systems, respectively. The scalar meson $\epsilon$ is universally coupled to the $\{20\}$ multiplet of SU(4) with $g_\epsilon^2/4\pi = 25.32$. The large width of the $\epsilon$ is taken into account by using a distributed mass spectrum.[3] For *each* $B_cN$ channel, one must in general determine $r_0$ independently. As a first approximation, we have used a *single* value of $r_0$, determined from the $^1S_0$ channel of maximum $I$. For the $NN$, $C_0N$, and $C_1N$ systems, we obtained $r_0 = 1.025$, 1.18, and 0.93 fm, respectively.

The $B_c$-nucleus potentials were calculated for $^4$He, $^{12}$C, $^{16}$O, $^{58}$Ni, and $^{208}$Pb. Some typical results are shown in Fig. 2 for $^{16}$O. The $B_c$-nucleus potentials are seen to be comparable in depth to the nucleon-nucleus potential. In Fig. 3 we display the calculated spectrum of $C_1$-nucleus shell-model states. The high level density and deep binding resulting from the high charmed-baryon mass is evident.

Several points are worth mentioning: (1) The spacings between both radial and orbital excitations vary inversely as the reduced mass, $M_R \approx M_{B_c}$ for heavy systems. Since $M_{B_c}/M_N \geqslant 2.4$, levels in a charmed nucleus will be much more closely spaced than in an ordinary nucleus. (2) Also for a heavy nucleus, the lowest state ($N = L = 0$) occurs close to the bottom of the well. (3) The close packing of $B_c$ levels implies that states of rather high $L$ may be bound in heavy charmed nuclei. For example a $C_1$ in a $j = \frac{25}{2}$ state coupled to a $i_{13/2}^{-1}$ neutron hole can produce states with spin up to 19 in $^{208}\mathrm{Pb}_{C_1}$.

We have already noted that $C_1N$ or $C_1$-nucleus states are unstable with respect to strong interactions, in fact decaying to $C_0$-nuclear states by single $\pi$ emission. One expects widths of order 1 MeV for such states, appreciably wider than the lowest $C_0$-nuclear states which decay only weakly. The observation of any such states will be similar, however, since $C_0$ tracks before decay will be very short. A most likely mode of detection of the bound states will probably parallel closely the observation of the free charmed baryons ($C_0$, $C_1$) themselves. In an emulsion, for example, an event which can be sufficiently detailed to reveal a charmed baryon mass lower than that a free $C_0$ + nuclear system would signal binding. Presumably a good sign of charmed nuclear states would be the observation of a strange hypernucleus as an intermediate step after the production of a charmed baryon, even though this is an event that may occur with low probability. Clearly the most favorable systems to study are $B_c$ nucleus for which binding is appreciable. Thus emissions placed in beams where slow charmed baryons or mesons are produced are likely to exhibit bound $B_c$ nuclei.

Finally one notes the strangeness analog states discussed by Kerman and Lipkin[8] will have their counterparts in charm analog states, the latter resulting from the extension of the Sakata model to SU(4).

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## Hydrogen Surface Contamination and the Storage of Ultracold Neutrons

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The results of hydrogen profile measurements on samples prepared by the same methods used for these materials when making ultracold-neutron (UCN) storage bottles show that the surface hydrogen is sufficient to account for the bulk of the anomalous shortening of UCN storage times.

Ultracold neutrons (UCN) are neutrons with such low energy ($\sim 10^{-7}$ eV) that they undergo total reflection for all angles of incidence from many materials. As such, these neutrons can be stored in "bottles" for substantial periods of time and may provide a source of neutrons upon which a variety of important fundamental measurements can be made, e.g., the accurate determination of the neutron lifetime[1] and a more sensitive search for its electric dipole moment.[2]

The first direct measurements of ultracold-neutron storage times[3] revealed serious discrepancies between the observed storage times, $\tau_s$, and those expected for an ideally flat, pure surface. While neutrons below a critical energy undergo total reflection, there are mechanisms which result in the "loss" of these neutrons. During a reflection, the neutron can be lost via nuclear capture or by being inelastically scattered to an energy above the critical energy.

Further experiments[4-8] have shown similar anomalies and much effort has been devoted to understanding of these results.[6,7,9-11] Several authors[3,6,7,12] have concluded that since the observed storage times appear to be independent of temperature, surface impurities cannot play an important role in the UCN losses. The only published data on this point, Fig. 1 (Refs. 7a and 7c), is indeed consistent with UCN losses being independent of temperature, but it does not rule out variations in $\tau_s$ up to a factor of 2 over the temperature range covered (dotted line, Fig. 1).

It is difficult to predict a unique temperature dependence of $\tau_s$ if impurities are responsible for the anomalies. If the impurity concentration was temperature independent and if the losses mainly result from inelastic scattering from a coherent scatterer or harmonically bound nucleus, the loss rate will be proportional to the Planck oscillator occupation number, i.e., $\sim T$ for large $T$ and $\sim \exp[-\omega_0/T]$ for small $T$ (with respect to the oscillator frequency $\omega_0$).[13] For incoherent