NEW POSSIBILITIES FOR EXOTIC HADRONS – ANTICHARMED STRANGE BARYONS☆

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A new candidate for an exotic hadron is presented: an anticharmed strange baryon denoted by $P_{\bar{c}s}$, a bound state of a nucleon and an F (now called D_s). Theoretical estimates of the binding energy due to the hyperfine interaction give values comparable to the binding of the H dibaryon, but color-electric repulsive effects which may be present in the H are expected to be smaller in the $P_{\bar{c}s}$. The $P_{\bar{c}s}$ may be more easily detected than the H because it has several distinctive signatures detectable against a multiparticle background.

The possible existence of exotic hadrons remains a principal question in hadron spectroscopy and the application of QCD to hadron physics [1].

One of the most interesting candidates for a bound exotic is the H dibaryon, a bound state of two Λ's. Although Jaffe's original calculation [2] and subsequent work [3] indicate a gain in hyperfine interaction energy by recoupling color and spins in the six-quark system over the two-Λ system, a lattice gauge calculation [4] indicates that the H is unbound and well above the ΛΛ threshold. Furthermore, although hyperfine binding calculations [3] indicate sensitivity of the hyperfine energy to flavor-SU(3) symmetry breaking, the lattice results are insensitive to the strange quark mass and SU(3) breaking [4].

This difference in the effects of SU(3) breaking suggests that different physics dominates lattice gauge and bag or potential model calculations. Which calculations include the important physics is not obvious. However, the lattice calculation shows a repulsive Λ–Λ interaction generated by quark exchange [5] which is not included in bag model calculations and could well prevent the six quarks from coming close enough together to feel the additional binding of the short range hyperfine interaction.

It is therefore of interest to look for other cases of hyperfine binding where such a repulsive exchange force may not be present. One case is an anticharmed strange baryon ($\bar{c}$uuds), which has a hyperfine binding roughly equal to that of the H, but which has no possibility of a quark exchange force in the lowest decay channel FN [6]. If this five quark system breaks up into an F and a nucleon, there is no possible quark exchange between the two hadrons without flavor exchange, and therefore no diagonal matrix element of the one-gluon-exchange interaction that could give rise to a short range repulsion.

Anticharmed baryons were suggested as good candidates for possible bound exotics at the 1980 baryon conference [7]. However, the nonstrange anticharmed baryon was not bound by the hyperfine interaction and the strange anticharmed baryon was very remote from experiment and not pursued seriously. Now, however, that charmed strange baryons have been produced in hadronic experiments which also produce charmed antiquarks, such experiments may also produce the anticharmed strange baryon if it is bound. We denote the anticharmed strange baryon as $P_{\bar{c}s}$ for pentaquark.

Examination of the color-spin hyperfine interaction in the H and $P_{\bar{c}s}$ using Jaffe's color-spin algebra [2] shows that both states have very similar properties and indicate that the $P_{\bar{c}s}$ may well be bound. The stability against breakup of an exotic multiquark system can be examined by checking whether hyperfine energy can be gained by recoupling the color and spins of the lowest lying two-hadron

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threshold [8]. We first calculate the hyperfine interaction for the $P_{\bar{c}s}$ and the H in the limit where the charmed quark has infinite mass and its hyperfine interaction energy is neglected, and SU(3) flavor symmetry is assumed for the light quarks. This calculation applies also for analogous states with heavier antiquarks which are equally attractive candidates for bound exotics. We use a variational approach with a wave function in which the two-body density matrix is the same for all pairs as in a baryon, and can then use the experimental N–Δ mass splitting to determine the strength of the hyperfine interaction energy [8].

A simplified form of the color-spin hyperfine interaction [2] can be used for systems containing only quarks and no active antiquarks:

$$V=-\tfrac{1}{2}v[C_6-C_3-\tfrac{8}{3}S(S+1)-16N]\,, \tag{1}$$

where v is a parameter defining the strength of the interaction, C_6 and C_3 denote the eigenvalues of the Casimir operators of the SU(6) color-spin and SU(3) color groups respectively, S is the total spin of the system and N is the number of quarks in the system. We consider the following states:

$$|\text{Nucleon}\rangle=|70,1,1/2,3\rangle\,, \tag{2a}$$

$$|\Delta\rangle=|20,1,3/2,3\rangle\,, \tag{2b}$$

$$|\Lambda\rangle=|70,1,1/2,3\rangle\,, \tag{2c}$$

$$|\mathrm{H}\rangle=|490,1,0,6\rangle\,, \tag{2d}$$

$$|\mathrm{P}_{\bar{c}s}\rangle=|210,3,0,4\rangle\,, \tag{2e}$$

$$|\mathrm{P}'_{\bar{c}s}\rangle=|105',3,0,4\rangle\,, \tag{2f}$$

where we have labeled the states by the quantum numbers $|d_6, d_3, S, N\rangle$ in a conventional notation [9] where d_6 and d_3 denote the dimensions of the color-spin SU(6) and color SU(3) representations in which the multiquark states are classified and primes distinguish between different representations with the same dimension. The $P_{\bar{c}s}$ is classified in the 210 of SU(6) to optimize the hyperfine interaction. We also define the state $P'_{\bar{c}s}$ in the 105′ of SU(6). For these two states we disregard the charmed antiquark which has no hyperfine interaction and consider the classification of only the state of the four quarks coupled to spin zero and a color triplet.

The hyperfine interaction (1) is easily evaluated for the states (2) by substituting the eigenvalues of the Casimir operators [9]:

$$C_6(70)=66,\quad C_6(20)=42\,, \tag{3a,b}$$

$$C_6(490)=144,\quad C_6(210)=304/3\,, \tag{3c,d}$$

$$C_6(105')=208/3\,, \tag{3e}$$

$$C_3(3)=16/3,\quad C_3(8)=12\,, \tag{4a,b}$$

We therefore obtain

$$V(\mathrm{N})=-8v,\quad V(\Delta)=8v\,, \tag{5a,b}$$

$$V(\Lambda)=-8v,\quad V(H)=-24v\,, \tag{5c,d}$$

$$V(\mathrm{P}_{\bar{c}s})=-16v,\quad V(\mathrm{P}'_{\bar{c}s})=0\,. \tag{5e,f}$$

$$M(\Delta)-M(\mathrm{N})=V(\Delta)-V(\mathrm{N})=16v\,, \tag{6a}$$

$$B(\mathrm{H})=V(\mathrm{H})-2V(\Lambda)=-8v$$

$$=-\tfrac{1}{2}[M(\Delta)-M(\mathrm{N})]\,, \tag{6b}$$

$$B(\mathrm{P}_{\bar{c}s})=V(\mathrm{P}_{\bar{c}s})-V(\Lambda)=-8v$$

$$=-\tfrac{1}{2}[M(\Delta)-M(\mathrm{N})]\,, \tag{6c}$$

where $B(X)$ denotes the difference in hyperfine energy between the state X and the relevant threshold. The gain in hyperfine interaction for the $P_{\bar{c}s}$ over the NF or ΛD threshold (degenerate in this symmetry limit) is equal to the gain for the H over the relevant ΛΛ threshold and is just half the Δ–N mass splitting.

At this level it appears that the $P_{\bar{c}s}$ is an equally attractive candidate for hyperfine binding as the H dibaryon [10,11].

The introduction of SU(3) symmetry breaking has been shown to reduce the binding of the H [3], and a similar effect occurs for the $P_{\bar{c}s}$. This is easily seen by noting that hyperfine binding energy of both the $P_{\bar{c}s}$ and the H is reduced by reducing the color-magnetic interaction of the strange quark. However the strange quark plays no role in the magnetic interactions of the ΛΛ, ΛD and NF final states and their hyperfine binding energies are unaffected by SU(3) symmetry breaking.

The broken-SU(3) hyperfine interaction can be written

$$V_{\mathrm{br}} = V(1-\delta) + \delta V_{\mathrm{n}} \qquad (7)$$

where δ is a parameter expressing the suppression of the strange quark hyperfine interaction and V_{n} is the hyperfine interaction (1) acting only in the space of the nonstrange quarks.

The eigenfunctions of the interaction V_{n} are states in which the colors of the three nonstrange quarks are coupled to either a singlet or an octet. We denote these respectively as $|\mathrm{n}_1; \mathrm{s}\rangle$ and $|\mathrm{n}_8; \mathrm{s}\rangle$. Their eigenvalues, obtained from eqs. (1)–(4) are

$$V_{\mathrm{n}}(\mathrm{n}_8; \mathrm{s}) = -2v , \qquad (8a)$$

$$V_{\mathrm{n}}(\mathrm{n}_1; \mathrm{s}) = -8v = \langle \mathrm{n}_1; \mathrm{s} | V | \mathrm{n}_1; \mathrm{s}\rangle = V(\mathrm{N}) , \qquad (8b)$$

where we have noted that the expectation values of the interactions V and V_{n} are equal for the state $|\mathrm{n}_1; \mathrm{s}\rangle$, since the diagonal matrix element of the interaction V between the color singlet cluster and the strange quark vanishes. The states $|\mathrm{n}_1; \mathrm{s}\rangle$ and $|\mathrm{n}_8; \mathrm{s}\rangle$ are linear combinations of the SU(6) eigenstates $\mathrm{P}_{\bar{c}s}$ and $\mathrm{P}'_{\bar{c}s}$. The SU(6) coefficients in the transformation between these two bases can be determined from the relation (8b). For any normalised linear combination of the two SU(6) eigenstates.

$$\langle \alpha \mathrm{P}_{\bar{c}s} + \beta \mathrm{P}'_{\bar{c}s} | V | \alpha \mathrm{P}_{\bar{c}s} + \beta \mathrm{P}'_{\bar{c}s}\rangle$$

$$= \alpha^2 V(\mathrm{P}_{\bar{c}s}) + \beta^2 V(\mathrm{P}'_{\bar{c}s}) = -16\alpha^2 v . \qquad (9)$$

Thus for the state $|\mathrm{n}_1; \mathrm{s}\rangle$, eqs. (8b) and (9) determine α^2 to be $\frac{1}{2}$. Choosing a simple real phase convention, we obtain

$$\mathrm{P}_{\bar{c}s} = (1/\sqrt{2})(|\mathrm{n}_1; \mathrm{s}\rangle + |\mathrm{n}_8; \mathrm{s}\rangle) , \qquad (10a)$$

$$\mathrm{P}'_{\bar{c}s} = (1/\sqrt{2})(|\mathrm{n}_1; \mathrm{s}\rangle - |\mathrm{n}_8; \mathrm{s}\rangle) , \qquad (10b)$$

Thus

$$\langle \mathrm{P}_{\bar{c}s} | V_{\mathrm{br}} | \alpha \mathrm{P}_{\bar{c}s}\rangle = -(16 - 11\delta) v , \qquad (11a)$$

$$\langle \mathrm{P}'_{\bar{c}s} | V_{\mathrm{br}} | \alpha \mathrm{P}'_{\bar{c}s}\rangle = 5\delta v , \qquad (11b)$$

$$\langle \mathrm{P}_{\bar{c}s} | V_{\mathrm{br}} | \alpha \mathrm{P}'_{\bar{c}s}\rangle = 3\delta v . \qquad (11c)$$

The eigenvalues of the 2×2 matrix (11) give the hyperfine energy in the broken SU(3) eigenstates. The lowest eigenvalue is seen to vary monotonically from the value $-16v$ in the symmetry limit where $\delta = 0$ to $-8v$ which is just the threshold energy in the case $\delta = 1$ where the strange quark hyperfine interaction is zero. A similar analysis to include the effect of the hyperfine interaction of a charmed antiquark with finite mass shows a very small change in the overall hyperfine energy relative to the NF threshold.

Detailed comparisons of the symmetry breaking for the $\mathrm{P}_{\bar{c}s}$ and the H as a function of δ and of the effect of finite charmed quark mass are discussed by Gignoux et al. [10,11] and show that the reduction of binding by symmetry breaking is less serious for the $\mathrm{P}_{\bar{c}s}$ than for the H. They obtained the same matrix (11) in a treatment using a flavor SU(3) classification instead of color-spin SU(6) and a diquark–diquark decomposition of the four-quark state instead of the 3–1 decomposition used here with the states $|\mathrm{n}_1; \mathrm{s}\rangle$ and $|\mathrm{n}_8; \mathrm{s}\rangle$. The treatments are equivalent since there is a one-to-one correspondence imposed by the Pauli principle between flavor SU(3) and color-spin SU(6) representations for spatially symmetric states containing only quarks and no antiquarks.

It does not seem possible at this time to give a credible quantitative estimate for the mass of the $\mathrm{P}_{\bar{c}s}$ since there are too many uncertain factors. The gain in potential energy from the hyperfine interaction is probably not enough to overcome the kinetic energy needed to keep the constituents confined in our trial wave function. Whether a better wave function and possibly a molecular type two-cluster wave function [12,13] might be barely bound may be calculable with sufficiently large lattices. The uncertainties in the color electric interaction and possible repulsive interactions must also be considered.

A rough estimate of the binding of a molecular type wave function extending over a distance large compared with the range of the hyperfine interaction is obtained from a model describing such states by a two-body Schrödinger equation with a short range potential proportional to the parameter $B(X)$, the gain in hyperfine energy by recoupling color and spin [14]. This equation has a bound state if $B(X)$ is greater than a critical value depending upon hadron masses and the strength of the interaction. The unknown proportionality factor between $B(X)$ and the interaction strength can be determined by comparison with the δ and S* mesons under the assumption that the latter are barely bound $K\bar{K}$ molecular states [14,13]. We then find that the H and $P_{\bar{c}s}$ will be more strongly bound than the δ and S* mesons if they satisfy the conditions

$$B(\mathrm{H})/B(\delta) > M(\mathrm{K})/M(\Lambda) = 0.44\,, \tag{12a}$$

$$B(\mathrm{P}_{\bar{c}s})/B(\delta) > M(\mathrm{K})[M(\mathrm{N})+M(\mathrm{F})]/2M(\mathrm{N})M(\mathrm{F}) = 0.39\,, \tag{12b}$$

where $B(\delta)$ is defined by analogy with eqs. (6) and the right-hand sides are the ratios of the relevant reduced masses.

From previous calculations [12]

$$B(\delta) = B(\mathrm{S}^*) = -\tfrac{3}{7}[M(\mathrm{K}^*) - M(\mathrm{K})] = -\tfrac{4}{7}[M(\Delta) - M(\mathrm{N})]\,, \tag{12c}$$

in the SU(3) symmetry limit and where the meson and baryon masses are related by assuming the same two-body radial density matrices for meson and baryon wave functions. With this value the conditions (12a,b) are well satisfied,

$$B(\mathrm{H})/B(\delta) = B(\mathrm{P}_{\bar{c}s})/B(\delta) = 7/8 > 0.44 > 0.39\,, \tag{12d}$$

The conditions (12a) and (12b) will still be satisfied if corrections from SU(3) breaking and other effects are less than a factor of two. Thus the H and $P_{\bar{c}s}$ are excellent candidates for weakly bound molecular states.

There is therefore interest both in experimental searches for the $P_{\bar{c}s}$ and in lattice gauge calculations. The simplest lattice calculation with an infinitely heavy charmed antiquark and four light quarks uuds, can easily be done in parallel with the more complicated H calculation both in the symmetry limit where all light quarks have the same mass and with SU(3) symmetry breaking. Comparing the results for these cases may provide considerable insight into our understanding of the physics of QCD in multiquark systems even if the $P_{\bar{c}s}$ is not found as a physical bound state in experiment. There is however a difficulty in treating loosely bound molecular states on the lattice, since these are sensitive both to the details of the short range hyperfine interaction and the long range part of the wave function. A proper treatment of both these effects may require a rather large lattice.

The experimental search for the $P_{\bar{c}s}$ has a completely different character from the search for the H, which has no easily detected signature, requires a sophisticated single-purpose experiment for its production and identifies it only by missing strangeness and missing mass. The $P_{\bar{c}s}$ has several distinctive signatures in decay modes detectable against a multiparticle background. Any multipurpose or exploratory experiment which produces charmed pairs in the presence of strange baryons can be used also to look for the $P_{\bar{c}s}$. In particular, an experiment which produces charmed strange baryons will also produce charmed antiquarks, and there may be a possibility that a strange baryon will pick up a charmed antiquark as well as a charmed quark in a hadronization process.

Some examples of decay modes which provide a good signature are

$$\mathrm{P}_{\bar{c}s} \to \mathrm{p}\phi\pi^-\,, \quad \mathrm{P}_{\bar{c}s} \to \mathrm{p}\eta\pi^-\,, \tag{13a,b}$$

$$\mathrm{P}_{\bar{c}s} \to \Lambda\mathrm{K}^+\pi^-\pi^-\,, \tag{13c}$$

$$\mathrm{P}_{\bar{c}s} \to \mathrm{pF}^-\,, \quad \mathrm{P}_{\bar{c}s} \to \Lambda\mathrm{D}^-\,, \tag{13d,e}$$

where the two decay modes (13d) and (13e) can occur if the $P_{\bar{c}s}$ is a resonance above the relevant threshold, or if it is a weakly bound "molecule". In the latter case the charmed meson would be "off shell" and have an apparently lower mass.

Note that there must always be associated production of a charmed particle together with the $P_{\bar{c}s}$. This can also be used to optimize signal to noise.

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