How the ''H particle'' unravels the quark dynamics

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It is shown that the short-range part of the Goldstone boson exchange interaction between the constituent quarks, which explains baryon spectroscopy and the short-range repulsion in the NN system, induces a strong short-range repulsion in the flavor-singlet state of the $S=-2$ system with $J^P=0^+$. This suggests that a deeply bound H-particle should not exist. We compare our approach with other models employing different hyperfine interactions between quarks in the nonperturbative regime of QCD. [S0556-2821(98)04007-7]

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Soon after the suggestion that the hyperfine splitting in hadrons should be due to the color-magnetic interaction between quarks [1,2] it has been noted by Jaffe [3] that the dibaryon $uuddss$ with $J^P=0^+$, $I=0$, called the H particle, is stable against strong decays. Its mass turned out to be about 80 MeV below the $\Lambda\Lambda$ threshold. The reason is that a flavour-singlet state in the $6q$ system is allowed in this case and the color-magnetic interaction gives more attraction for the most favorable configuration $[33]_{CS}$ than for two well-separated Λ-hyperons. In Jaffe's picture the H particle should be a compact object, in contrast with the molecular-type structure of the deuteron.

Since Jaffe's prediction many calculations have appeared in a variety of models [4]. They give a wide range of predicted masses, depending on the model. Realistic calculations usually predict a well-bound H particle. In particular, the quark-cluster calculations suggest that the implications of the color-magnetic interaction are radically different in two-nucleon and in coupled YY-YN systems. While in the former case the color-magnetic interaction between quarks gives rise to a strong short-range repulsion in the NN system, in the latter, there appears either a soft attraction or a soft repulsion at short-range [5] when the linear combination of the coupled channels is close to a flavor-singlet state. This soft short-range interaction, reinforced by the medium and long-range attraction coming from the meson exchange between lambdas, provides a bound state with a binding energy of the order 10–20 MeV [6] or even 60–120 MeV [7]. However, a simple quark-cluster variational basis, used in these calculations, is rather poor at short-range. While it is not so important for the baryon-baryon systems with strong repulsion at short-range, this shortcoming becomes crucial for the $\Lambda\Lambda-N\Xi-\Sigma\Sigma$ system, with the quantum numbers of H. As soon as a simple quark-cluster variational basis is properly extended, a very deeply bound state with the binding energy of about 250 MeV is found [8].

The existence or non-existence of the H particle has to be settled by experiment. For approximately 20 years several experiments have been set for ''hunting'' the H particle. The very recent high-sensitivity search at Brookhaven [9] gives no evidence for the production of deeply bound H, the production cross section being one order of magnitude below the theoretical estimates.

It has recently been suggested that in the low-energy regime, light and strange baryons should be considered as systems of three constituent quarks with a QQ interaction (Q is a constituent quark, to be contrasted with a current quark q) that is formed of a central confining part and a chiral interaction that is mediated by Goldstone bosons between constituent quarks [10]. Indeed, at low temperatures and densities, the underlying chiral symmetry of QCD is spontaneously broken by the QCD vacuum. This implies that the valence quarks acquire a constituent (dynamical) mass, which is related to the quark condensates $\langle\bar{q}q\rangle$, and at the same time the Goldstone bosons π, K, η appear, which couple directly to the constituent quarks [11]. It has been shown that the hyperfine splittings as well as the correct ordering of positive and negative parity states in spectra of baryons with valence u,d and s quarks are produced in fact, not by the color-magnetic part of the one-gluon exchange interaction (OGE), but by the short-range part of the Goldstone boson exchange (GBE) interaction [10,12,13]. This short-range part of the GBE interaction has just opposite sign as compared to the Yukawa potential tail and is much stronger at short interquark separations. There is practically no room for the OGE interaction in light baryon spectroscopy and any appreciable amount of color-magnetic interaction, in addition to GBE, destroys the spectrum [14]. The same short-range part of the GBE interaction, which produces good baryon spectra, also induces a short-range repulsion in the NN system [15]. Thus it is interesting to study the short-range interaction in the $\Lambda\Lambda$ system and the stability of the H particle in the GBE model.

For qualitative insight it is convenient first to consider a schematic quark-quark interaction which neglects the radial dependence of the GBE interaction. In this model, the short-range part of the GBE interaction between the constituent quarks is approximated by

$$V_\chi = -C_\chi \sum_{i<j} \lambda_i^F \cdot \lambda_j^F \vec{\sigma}_i \cdot \vec{\sigma}_j, \tag{1}$$

where λ_i^F with an implied summation over F (F=1,2,...,8) are the quark flavor Gell-Mann matrices and $\vec{\sigma}$ the spin ma-

TABLE I. Expectation values of the operator (1) in C_χ units corresponding to the states (5).

$[f]_O[f]_{FS}[f]_{OC}$	$\langle V_\chi\rangle/C_\chi$
$[6]_O[33]_{FS}[222]_{OC}$	-24
$[42]_O[411]_{FS}[3111]_{OC}$	-24
$[42]_O[33]_{FS}[222]_{OC}$	-24
$[42]_O[2211]_{FS}[42]_{OC}$	8

trices. The minus sign of the interaction (1) is related to the sign of the short-range part of the GBE interaction (which is opposite to that of the Yukawa potential tail), crucial for the hyperfine splittings in baryon spectroscopy.

In an harmonic oscillator basis, $\hbar\omega$ and the constant C_χ implied by the schematic model (1), can be determined from the $\Delta-N$ and $N(1440)-N$ splittings to be $C_\chi=29.3$ MeV and $\hbar\omega\sim250$ MeV [10].

The color- and flavor-singlet $uuddss$ states are described by the $[222]_C$ and $[222]_F$ Young diagrams respectively. For the S-wave relative motion of two s^3 clusters, the spatial symmetries of the $6Q$ system are $[6]_O$ and $[42]_O$ and for the spin S=0 the corresponding spin symmetry is $[33]_S$. The antisymmetry condition requires $[f]_{FS}=[\tilde{f}]_{OC}$, where $[\tilde{f}]$ is the conjugate of $[f]$. Thus, among the states given by the inner products

$$[33]_S\times[222]_F=[33]_{FS}+[411]_{FS}+[2211]_{FS}+[1^6]_{FS}, \tag{2}$$

$$[6]_O\times[222]_C=[222]_{OC}, \tag{3}$$

$$[42]_O\times[222]_C=[42]_{OC}+[321]_{OC}+[222]_{OC}+[3111]_{OC}+[21111]_{OC} \tag{4}$$

only the four states are allowed:

$$|1\rangle=|[6]_O[33]_{FS}[222]_{OC}\rangle$$

$$|2\rangle=|[42]_O[33]_{FS}[222]_{OC}\rangle$$

$$|3\rangle=|[42]_O[411]_{FS}[3111]_{OC}\rangle$$

$$|4\rangle=|[42]_O[2211]_{FS}[42]_{OC}\rangle. \tag{5}$$

For each of these states the expectation value of the interaction (1) can be easily calculated in terms of the Casimir operators eigenvalues for the groups $SU(6)_{FS}$, $SU(3)_F$ and $SU(2)_S$ using the formula given in Appendix A of Ref. [15]. The corresponding matrix elements are given in Table I. One can see that the interaction (1) is attractive for the states $|1\rangle-|3\rangle$ and repulsive for $|4\rangle$. This suggests that it is a good approximation to restrict the basis to $|1\rangle$, $|2\rangle$ and $|3\rangle$ for the diagonalization of a more realistic Hamiltonian. Keeping in mind that the spatial symmetry $[6]_O$ is compatible with the s^6 configuration, one can roughly evaluate the energy of the lowest $6Q$ configuration relative to the 2Λ threshold as

$$\begin{aligned}&\langle s^6[6]_O[33]_{FS}|H_0+V_\chi|s^6[6]_O[33]_{FS}\rangle\\&\quad-2\langle s^3[3]_O[3]_{FS}|H_0+V_\chi|s^3[3]_O[3]_{FS}\rangle\\&\quad=4C_\chi+3/4\hbar\omega=305\ \text{MeV},\end{aligned} \tag{6}$$

where H_0 is the kinetic energy in the $6Q$ or $3Q$ system. While here and below we use notations of the shell model, it is always assumed that the center of mass motion is removed. In deriving the kinetic energy, $3/4\hbar\omega$, we have neglected the mass difference between u, d and s constituent quarks. The pair-wise color electric confinement contribution is exactly the same for s^6 configuration and for two well separated s^3 clusters, so it cancels out.

This simple estimate shows that the lowest ''compact'' flavor-singlet $6Q$ state with quantum numbers $J^P=0^+,I=0,S=-2$ lies a few hundred MeV above the $\Lambda\Lambda$ threshold.

In a more quantitative calculation we use the Hamiltonian [10,12]:

$$H=\sum_{i=1}^{6} m_i+\sum_i \frac{\vec{p}_i^{\,2}}{2m}-\frac{(\Sigma_i\vec{p}_i)^2}{12m}+\sum_{i<j} V_{conf}(r_{ij})+\sum_{i<j} V_\chi(\vec{r}_{ij}) \tag{7}$$

where the confining interaction is

$$V_{conf}(r_{ij})=-\frac{3}{8}\lambda_i^c\cdot\lambda_j^c\ C\ r_{ij} \tag{8}$$

and the spin-spin component of the GBE interaction between the constituent quarks i and j reads

$$V_\chi(\vec{r}_{ij})=\left\{\sum_{F=1}^{3} V_\pi(\vec{r}_{ij})\lambda_i^F\lambda_j^F+\sum_{F=4}^{7} V_K(\vec{r}_{ij})\lambda_i^F\lambda_j^F+V_\eta(\vec{r}_{ij})\lambda_i^8\lambda_j^8+V_{\eta'}(\vec{r}_{ij})\lambda_i^0\lambda_j^0\right\}\vec{\sigma}_i\cdot\vec{\sigma}_j, \tag{9}$$

where $\lambda^0=\sqrt{2/3}\mathbf{1}$ ($\mathbf{1}$ is the 3×3 unit matrix). The interaction (9) includes π,K,η and η' exchanges. In the large-N_c limit, where the axial anomaly vanishes [16], the spontaneous breaking of the chiral symmetry $U(3)_L\times U(3)_R\to U(3)_V$ implies a ninth Goldstone boson [17], which corresponds to the flavor singlet η'. Under real conditions, where $N_c=3$, a certain contribution from the flavor singlet remains and the η' must thus be included in the GBE interaction.

In the simplest case, when both the constituent quarks and mesons are point-like particles and the boson field satisfies the linear Klein-Gordon equation, one has the following spatial dependence for the meson-exchange potentials [10]:

$$V_\gamma(\vec{r}_{ij})=\frac{g_\gamma^2}{4\pi}\frac{1}{12m_im_j}\left\{\mu_\gamma^2\frac{e^{-\mu_\gamma r_{ij}}}{r_{ij}}-4\pi\delta(\vec{r}_{ij})\right\}\times(\gamma=\pi,K,\eta,\eta') \tag{10}$$

where μ_γ are the meson masses and $g_\gamma^2/4\pi$ are the quark-meson coupling constants given below.

Equation (10) contains both the traditional long-range Yukawa potential as well as a δ-function term. It is the latter that is of crucial importance for baryon spectroscopy and the short-range NN interaction since it has a proper sign to provide the correct hyperfine splittings in baryons and is becoming highly dominant at short range. Since one deals with structured particles (both the constituent quarks and pseudoscalar mesons) of finite extension, one must smear out the δ-function in Eq. (10). In Ref. [18] a smooth Gaussian term has been employed instead of the δ-function

$$4\pi\delta(\vec{r}_{ij})\Rightarrow\frac{4}{\sqrt{\pi}}\alpha^3\exp(-\alpha^2(r-r_0)^2), \qquad (11)$$

where α and r_0 are adjustable parameters.

The parameters of the Hamiltonian (7)–(9) are [18]

$$\frac{g^2_{\pi q}}{4\pi}=\frac{g^2_{\eta q}}{4\pi}=0.67,\ \frac{g^2_{\eta' q}}{4\pi}=1.206,$$

$$r_0=0.43\ \text{fm},\ \ \alpha=2.91\ \text{fm}^{-1},\ \ C=0.474\ \text{fm}^{-2},$$

$$\mu_\pi=139\ \text{MeV},\ \ \mu_\eta=547\ \text{MeV},$$

$$\mu_{\eta'}=958\ \text{MeV},\ \ \mu_K=495\ \text{MeV}. \qquad (12)$$

The Hamiltonian (7)–(12) with constituent masses $m_{u,d}=340$ MeV and $m_s=440$ MeV provides a very satisfactory description of the low-lying N and Δ spectra in a fully dynamical nonrelativistic 3-body calculation [18] as well as of the strange baryon spectra [19]. However, this parametrization should be considered as an effective one only. Indeed, the volume integral of the GBE interaction should be zero [13], while in the parametrization above this is not so because of the off-shift r_0 of the ‘‘contact’’ term. In Ref. [13] a very good fit of the nonstrange and strange baryon spectra has been obtained in a fully dynamical calculation without such an off-shift. There a relativistic kinematics for the constituent quarks has been used. Thus one should consider the above off-shift only as an artifact of the fit of baryon masses with the nonrelativistic kinematics used in Ref. [18].

In principle it would be better to use the parametrization of Ref. [13]. However, in applying the quark cluster approach to two-baryon systems we are restricted to use a nonrelativistic kinematics and a s^3 wave function for the ground state baryons. With such an approximation the nonrelativistic model of Ref. [18] works well. For example, the quantity $\langle\Lambda|H|\Lambda\rangle$ reaches its minimum of 1165.4 MeV at an harmonic oscillator parameter value of $\beta=0.449$ fm, i.e., only about 40 MeV above the result obtained in the dynamical 3-body calculations of [19]. On the other hand, the nonrelativistic s^3 ansatz is not compatible with the model of Ref. [13]. Since in this paper we study only qualitative effects, related to the spin-flavor structure and sign of the short-range part of the GBE interaction, we consider the nonrelativistic approach as a reasonable framework.

We calculate the potential in the flavor-singlet $S=-2$ two-baryon system at zero separation between clusters in the adiabatic (Born-Oppenheimer) approximation defined as

$$V(R)=\langle H\rangle_R-\langle H\rangle_\infty, \qquad (13)$$

where R is a collective coordinate which is the separation distance between the two s^3 clusters, $\langle H\rangle_R$ is the lowest expectation value of the Hamiltonian describing the $6Q$ system at fixed R and $\langle H\rangle_\infty=2m_\Lambda$, i.e., the energy of two well separated lambdas, obtained with the same Hamiltonian.

It has been shown by Harvey [20] that when the separation R between two s^3 clusters approaches zero, then only two types of $6Q$ configurations survive: $|s^6[6]_O\rangle$ and $|s^4p^2[42]_O\rangle$. Thus, in order to extract an effective potential at zero separation between clusters in the adiabatic approximation, we diagonalize the Hamiltonian (7)–(12) in the basis of the first three states defined by Eq. (5). All the necessary matrix elements are calculated with the help of the fractional parentage technique, also used in a study of the short-range NN interaction in Ref. [15].

We find the lowest eigenvalue of the flavor-singlet state $J^P=0^+$ to be 847 MeV above the $\Lambda\Lambda$ threshold. According to Eq. (13), there is a strong short-range repulsion in a two-baryon flavor-singlet $S=-2$ system in the $^1S_0$ wave. It then definitely suggests that within the physical picture under discussion a compact (well bound) H particle should not exist.

The value of the repulsion given above depends on the way the kinetic energy of the $6Q$ system was calculated. For simplicity, in the kinetic energy term only, we considered that the u, d and s quarks have the same mass $\overline{m}=(4m_u+2m_s)/6$. We have also carried calculations in the extreme limits $\overline{m}=m_u$ and $\overline{m}=m_s$ and obtained 1050 MeV and 531 MeV above the $\Lambda\Lambda$ threshold, respectively. These extreme values just prove that the strong repulsion persists in any case.

This result is in a sharp contrast with results derived from the models based on the color-magnetic interaction. We consider it as an additional evidence in favor of the GBE model. Indeed, a deeply bound H particle is definitely excluded by experiment [9] and the color-magnetic interaction, at variance with GBE interaction, implies a deeply bound state (see the introduction).

There are suggestions that the instanton-induced (’t Hooft) interaction in QQ pairs could be important for the hyperfine splittings in baryons [21]. Assuming that this interaction is responsible for the essential part of the $\Delta-N$ hyperfine splitting, the deeply bound H particle should also disappear [22]. The ’t Hooft interaction is very strong and attractive in the color-singlet $q\overline{q}$ pseudoscalar channel and could be indeed responsible for the chiral symmetry spontaneous breaking in the QCD vacuum [23] and be the most important interaction in mesons. Thus to the extent that the ’t Hooft interaction contributes to $q\overline{q}$ pseudoscalar pairs, it is automatically included in the GBE interaction in QQ pairs (the ’t Hooft interaction could be responsible, at least in part, for the pole in the t-channel). However, the ‘‘direct’’ ’t Hooft interaction in qq pairs is rather weak. There are also indications from lattice QCD that the ‘‘direct’’ instanton-induced interaction in qq pairs cannot be important for the $\Delta-N$ splitting. For example, the $\Delta-N$ splitting disappears after cooling [24] (only instantons survive the cooling procedure), while it is appreciable before cooling. There is also evidence from lattice QCD that the hyperfine splittings are related mostly to $q\overline{q}$ excitations in baryons, but not to forces mediated by gluonic fields in qq pairs [25]. Simple symmetry arguments also show that the ''direct'' 't Hooft interaction in QQ pairs cannot provide a correct ordering of the lowest positive and negative parity states in light and strange baryon spectra [10] [for baryon spectra obtained in such a model in a nonperturbative calculation see second paper of Ref. [21]. From that paper, one can see, indeed, that the lowest positive parity excitations in all parts of the spectrum—$N(1440)$, $\Delta(1600)$,$\Lambda(1600)$,$\Sigma(1660)$, . . . — lie much above the negative parity excitations.]

One should also mention the QCD sum rule estimate for the H particle [26]. There it was shown that there is no qualitative difference between the NN and the $\Lambda\Lambda$ systems (including the flavor singlet channel), which strongly supports our point of view.

Here we have considered the $6Q$ $S=-2$ system in a flavor singlet state only (the ''H particle'' channel) and found that there appears a strong short-range repulsion in the $^1S_0$ partial wave. This strong short-range repulsion implies that a deeply bound (on nuclear scale) H particle should not exist. The same analysis can be extended to the $\Lambda\Lambda$ system in all allowed flavor states. Then, similarly to the NN system [15], there will appear a strong short-range repulsion coming from the same short-range part of the GBE interaction. There is however an attraction in the $\Lambda\Lambda-\Sigma\Sigma-N\Xi$ system at medium- and long-range, coming from the Yukawa potential tail of the GBE interaction as well as from correlated two-pseudoscalar-meson exchange. At the moment, one cannot exclude that this interaction could weakly bind $\Lambda\Lambda$ in a molecule-like system of nuclear nature. However, in its origin this attraction should be similar to the attraction in the $^1S_0$ partial wave of the NN system, which is too weak to bind the system. A firm prediction of the existence or nonexistence of a weakly bound $\Lambda\Lambda$ system of nuclear nature can only be made in a fully dynamical calculation.

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