July 17, 1986

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THE LIFETIME OF A DEEPLY BOUND H DIBARYON AND THE CYGNUS X-3 EVENTS

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Received 3 March 1986

The explanation of the Cygnus X-3 events which involves a strange-quark star emitting the doubly strange H dibaryon requires the H to have a lifetime of greater than ten years. We calculate the $\Delta S = 2$ $H \rightarrow NN$ for an H bound below ΛN threshold, and find $\tau_H \leq 6$ days.

The recent observation [1] of high-energy muons arriving from the direction of Cygnus X-3 and correlated with its 4.8 h orbital period is a puzzling and intriguing phenomenon. The muon flux indicates that the particles coming from Cygnus must be hadronic, producing the muons by interaction either in the atmosphere or within the rock above the detector. In order to arrive in phase, the particles must have [2] $\gamma \sim \text{few} \times 10^3$, while not to have decayed in flight requires a lifetime $\tau \gtrsim 10$ yr. These properties are not matched by known particles, and, if the observations are in fact not spurious (which still remains a reasonable possibility), the Cygnus X-3 signal would indicate interesting new physics. One proposed explanation [2] of the events suggests that Cygnus is a form of matter of high strangeness, and that the emitted particles are the neutral $S = -2$ dibaryons proposed by Jaffe [3]. This particle, usually called the H, is a bound state of uuddss, with overall spin and isospin zero, and is sometimes called the "dilambda" (this term is somewhat misleading as there are other components to its wavefunction). Although in his original article Jaffe suggested only that the H was bound enough to be stable against strong decay, i.e. $m_H < 2m_\Lambda$, the Cygnus hypothesis would require that it also be below threshold for $\Delta S = 1$ weak decay ($m_H < m_N + m_\Lambda$) in order that its lifetime have the possibility of being as long as ten years. The decay of the H would then proceed via $H \rightarrow NN$, and would be governed by $\Delta S = 2$ weak interactions. It is the purpose of this letter to present a calculation of the life-

time of an H dibaryon bound deeply enough that it must decay with $\Delta S = 2$. We find a lifetime of less than $5 \times 10^5 \text{ s} \approx 6 \text{ d}$. Such a value is difficult to reconcile with the requirements of Cygnus X-3. We will, in a separate work [4], discuss the $\Delta S = 1$ decay of a more conventional H with $m_N + m_\Lambda \leq m_H < 2m_\Lambda$, and also at that time provide further detail on our methods.

Experience with $\Delta S = 2$ interactions in the kaon system can provide a guide for $\Delta S = 2$ H decay ^{#1}. Thus the dominant short-distance process is the box diagram, pictured in fig. 1. In addition one may in general have long-distance, dispersive contributions; i.e. the action of two $\Delta S = 1$ interactions with a low-energy intermediate state between them. Formally the box diagram should be larger than the dispersive component by a factor of m_c^2/μ^2 (where μ is some typical hadronic mass scale, e.g. $\mu \sim 300 \text{ MeV}$), sim-

^{#1} For a review see ref. [5].

Fig. 1. The box diagram for the transition $ss \rightarrow dd$.

ply due to the dimensional parameters involved. However, in $K^0\bar{K}^0$ mixing both types of terms seem to be roughly the same magnitude [6]. This equality may be due to the helicity suppression in the hadronic matrix element of the box diagram [7]. Such a suppression is not, however, present in H decay and we thus would expect the box diagram to be larger than or comparable to the dispersive effects in this case. In addition, because estimates of the dispersive component are bound to be less reliable than the relatively straightforward calculation of the box diagram we will focus our discussion on the latter.

For our calculation we will utilize the P matrix of Jaffe and Low [8]. As many readers are no doubt unaware of P -matrix methods, we provide here a brief overview at a crude level, to be made precise below. The effective $\Delta S = 2$ hamiltonian which arises from the box diagram is given by

$$H^{\Delta S=2} = (G_F^2/16\pi^2)m_c^2 \cos^2\theta_c \sin^2\theta_c \times \bar{\eta} \bar{d} \gamma^\mu (1 + \gamma_5) s \bar{s} \bar{\gamma}_\mu (1 + \gamma_5) s, \quad (1)$$

where $\eta \sim 0.7$ is the well-known factor arising from QCD radiative corrections [9] and we have dropped the top-quark contribution, since small KM angles render it less important. The hamiltonian, when acting on the H, will produce a state with the quantum numbers of NN in a 1S_0 or 3P_0 configuration. Such a state can fall apart into two neutrons. The P matrix allows us to handle this situation by describing the properties of an artificially confined six-quark state (i.e. NN) which can mix with the H. Thus roughly we have

$$\Gamma_H = |\langle ''NN'' | H_w^{\Delta S=2} | H \rangle / (m_H - m_{''NN''})|^2 \Gamma_{''NN''}. \quad (2)$$

Our primary calculation is that of the weak matrix element, and this in turn is used in the P matrix formalism in order to produce the calculated H width. A point to be emphasized is that the P matrix is a rigorous method of handling the S matrix and scattering data. Phase shift data can always be described by a P matrix for any matching radius b (thus, e.g., for $b = 6.5 \text{ GeV}^{-1}$ and the 1S_0 state, Mulders [10] finds $m_{''NN''} = 2.25 \text{ GeV}$, $\Gamma_{''NN''} = 0.43 \text{ GeV}$; see below for details). Model dependence enters only in the identification of P matrix poles with quark model eigenstates.

First we consider the parity-conserving decay of the H to the 1S_0 NN final state. The six-quark con-

figuration which breaks apart into this channel has all six quarks in identical $S_{1/2}$ spatial states, in which case the quantum numbers of the resulting color singlet hadrons are severely restricted. Jaffe [3] has analyzed this case, introducing the useful classification of $SU(6)$ color-spin. In particular the only even-parity $J = 0$ states fall into $SU(3)_{\text{FLAVOR}}$ multiplets of 1, 27 and 28 (corresponding to 490, 189 and 1 of color-spin, respectively) with each multiplet occurring only *once*. This is extremely useful information, as it implies that if one writes down *any* color singlet, spin-singlet nonvanishing combination of six quarks in a flavor singlet, then this is the unique wavefunction of the H. Similar statements apply for members of the 27 and 28. In practice there is often more than one way to construct a given wavefunction; however the uniqueness of the flavor decomposition requires that any two wavefunctions with the same quantum numbers be in fact identical. We have verified this identity in several cases^{‡2}. The wavefunction of the H, and the neutral $S = 0$ member of the 27 can be given using quark creation operators, and the notation

$$(abcdef) \equiv \epsilon^{\alpha\beta\gamma} \epsilon^{\eta\sigma\rho} \epsilon^{mn} \epsilon^{pq} \epsilon^{st} a_{\alpha m}^\dagger b_{\beta n}^\dagger c_{\gamma s}^\dagger d_{\eta p}^\dagger e_{\sigma q}^\dagger f_{\rho t}^\dagger, \quad (3)$$

with $a, b, \dots, f$ being flavor quantum numbers, $\alpha, \beta, \dots, \rho$ colors, and $m, n, \dots, t$ spins. Here $\epsilon^{\alpha\beta\gamma}$ is the usual antisymmetric tensor and ϵ^{mn} is its two variable counterpart

$$\epsilon^{mn} = (i\sigma_2)_{mn}.$$

All creation operators create a quark in the ground-state bag model mode. The H wavefunction can then be written as

$$\begin{aligned} |H\rangle = \frac{1}{144} \{ & (udsuds) + (usdusd) + (dsudsu) \\ & - (ussudd + udduss) - (dssduu + duudss) \\ & - (usudsd + dsdusu) \} |0\rangle, \end{aligned} \quad (4)$$

^{‡2} The flavor singlet H can be formed from a flavor octet $\otimes$ octet product. The case of two flavor singlets, i.e. $\epsilon^{abc} \epsilon^{def} (abcdef)$ is the notation of eq. (3), vanishes because there is no $SU(3)$ singlet combination of three ground-state quarks. We have also verified the equivalence of color singlet $\otimes$ singlet and color octet $\otimes$ octet. This latter equivalence has important ramifications for the production rate of the H and for the concept of "hidden color" in nuclear states.

and the 1S_0 NN six-quark state is contained in

$$|27\rangle = (1/12\sqrt{10})(\text{udd udd})|0\rangle. \quad (5)$$

Using these wavefunctions, the calculation of the matrix element proceeds along standard lines [5], except for the unusual complexity of the spin and color sums. In particular we find

$$\begin{aligned} \langle 27|H_w^{\Delta S=2}|H\rangle &= (G_F^2/16\pi^2)m_c^2 \cos^2\theta_1 \sin^2\theta_1 \eta \\ &\times \frac{34\,560}{144 \cdot 12\sqrt{10}} d^3x(uu' + \ell\ell')^2, \end{aligned} \quad (6)$$

where the overlap is given in the notation of ref. [5] and the primes signify the s-quark wavefunctions. Numerically, using R_H obtained via the bag-model virial theorem $R_H \approx 21^{1/3} R_N \approx 6.3 \text{ GeV}^{-1}$, and $m_s = 0.3 \text{ GeV}$, this amounts to

$$\begin{aligned} \int d^3x(uu' + \ell\ell')^2 &= 0.23/R_H^3 = 0.88/V_H \\ &= (0.097 \text{ GeV})^3, \end{aligned} \quad (7)$$

and

$$\langle 27|H_w|H\rangle = 3.3 \times 10^{-16} \text{ GeV}. \quad (8)$$

It is also possible for the H to mix with a 3P_0 NN state through the parity violating part of the weak hamiltonian. In the six-quark picture, this would be contained in a configuration with one quark in a P wave orbital. We have estimated the contribution of the P state to H decay and find it small for three reasons: (i) P wave phase space, (ii) a larger energy denominator, and (iii) a smaller wavefunction overlap in the hadronic matrix element. Overall these put the P wave contribution at less than 5% of that of the S wave and thus we neglect it in our remaining discussion. Note too that since S and P waves do not interfere in the total decay rate, the calculation of the S wave decay alone provides an upper bound on the lifetime.

In order to decide how to use this result — eq. (6) — we turn to the P matrix. While the H is meant to be a true narrow resonance, the $|27\rangle$ NN six-quark state is not a resonance at all, as it can readily fall apart into two neutrons. It is “bound” in the bag model (or any other quark model) by imposing boundary conditions by hand. The P matrix tells us how to handle this situation. Poles in the P matrix correspond to

eigensolutions of particles artificially confined to a spherical square well of radius b . These can be related to the S matrix with $l=0, S=0$, by

$$S = \exp(-ikb)[(1+i\bar{P})/(1-i\bar{P})]\exp(-ikb), \quad (9)$$

where, in the elastic region,

$$S = \exp(2i\delta), \quad \bar{P} = \cot(kb + \delta), \quad (10)$$

with $k^2 = (s - 4m^2)/4$. Note that even in the *absence* of any interaction (i.e. with $\delta=0$) there is a tower of poles in the P matrix

$$\bar{P}|_{\text{no interaction}} = \frac{1}{kb} + \frac{2k}{b} \sum_{n=1}^{\infty} \frac{1}{k^2 - n^2\pi^2/b^2}. \quad (11)$$

Jaffe and Low call these “compensation poles”. The lowest of these is the NN state which we are dealing with. In the real world, the locations and residues of the P matrix poles have been determined from nucleon–nucleon scattering data by Mulders [10]. The lowest NN state has mass 2.25 GeV and residue 1.35 GeV^2 . (Here we quote the fit with $b=6.5 \text{ GeV}$ as this is the smallest value which can contain an H of the required radius.)

A narrow resonance like the H also appears as a pole in the P matrix, but with a small residue. It is the residue which determines the resonant width. Thus the S matrix for a Breit–Wigner resonance is

$$S = 1 + 2im\Gamma/(s - m^2 - im\Gamma), \quad (12)$$

which has the property

$$S|_{s=m^2} = -1, \quad \partial S/\partial s|_{s=m^2} = -2i/m\Gamma. \quad (13)$$

The associated P matrix will have the form

$$\bar{P} = c + \bar{r}/(s - s_0) \quad (14)$$

in the neighborhood of a narrow pole. This implies that

$$\partial S/\partial s \approx + (2i/\bar{r})S \{[\bar{r}^2/(s - s_0)^2]/(1 + \bar{P}^2)\} \rightarrow -2i/r, \quad (15)$$

where the last line is evaluated near the P matrix pole. This allows us to identify (in the narrow resonance limit)

$$\Gamma = \bar{r}/m. \quad (16)$$

In the absence of the weak interaction, the H will have no residue in the NN channel as it is below thresh-

old for decay. Adding the $\Delta S = 2$ interaction in an artificially confined bag-model problem will induce mixing between the H and the 27-plet member with NN quantum numbers, of the form

$$M = \begin{pmatrix} m_H & \langle H | H_w | 27 \rangle \\ \langle 27 | H_w | H \rangle & m_{27} \end{pmatrix}. \quad (17)$$

This mixing will induce a residue for the H in the amount

$$\bar{r}_H = |\langle 27 | H_w | H \rangle / (m_H - m_{27})|^2 \bar{r}_{27}. \quad (18)$$

This plus eq. (16) leads to our final formula for the decay rate

$$\Gamma_H = |\langle 27 | H_w | H \rangle / (m_H - m_{27})|^2 \bar{r}_{27} / m_H. \quad (19)$$

In applying this result, m_{27} and $\bar{r}_{27}$ are obtained from Mulder's fit to the NN scattering data, with the exception that $\bar{r}_{27}$ is scaled down by S wave phase space (i.e. k_H/k_{27}) in order to reflect the kinematic variation of the residue (cf. eq. (11)). Thus we employ $\bar{r}_{27} = 0.89 \text{ GeV}^2$ at a mass $m_H \approx m_N + m_\Lambda$. Combining then our calculated overlap with the P -matrix relation — eq. (19) — we find

$$\Gamma_H = 1.4 \times 10^{-30} \text{ GeV} \quad (20)$$

and

$$\tau_H = 0.5 \times 10^6 \text{ s}. \quad (21)$$

While in any calculation there is some uncertainty, we do not feel that the lifetime can be much longer than that given in eq. (21). Experience has taught us that, of the various quark models, the bag model generally gives the smallest matrix element and hence longest lifetime when applied to weak decays such as $\Delta S = 1$ nonleptonic decay or proton decay. Other models would then be expected to produce a shorter lifetime. Bag model methods applied to hyperon decay come within a factor of two of the data. Within the framework of the present calculation use of Mulder's fit with $b = 7.0 \text{ GeV}^{-1}$ would yield a lifetime of $14 \times 10^5 \text{ s}$, while the fit with $b = 6.0 \text{ GeV}^{-1}$ we would have to scale down the H radius to about

5 GeV^{-1} , yielding $\Gamma_H = 3.1 \times 10^6 \text{ s}$. Imposing Jaffe and Low's suggested relation for the matching radius, $b = 1.4R$, would further reduce the lifetime. The spin sums in eq. (6) are model independent, and the wavefunction overlap in eq. (7) is not overly large. Thus we feel that, if anything, we have overestimated the lifetime. A recent estimate [11] of the dispersive component of $H \rightarrow NN$ yields $\tau \sim 10^4 \text{ s}$. Given the rough nature of their estimate we take this to be compatible with ours.

Our calculation shows then that the lifetime of a deeply bound H decaying through $\Delta S = 2$ processes is of the order of 10^6 s . This is too small by two orders of magnitude to allow for the H interpretation of the Cygnus X-3 events as described in ref. [2].

We thank L. McLerran for encouraging us to do this calculation and R. Jaffe for helpful conversations. This work was supported in part by the US National Science Foundation.

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