Charmed baryons as soliton–D meson bound states

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We show that charmed baryons are remarkably well described as bound states of the D mesons with a soliton. The scheme exploits the isospin–spin transmutation that takes place when a doublet pseudoscalar meson is trapped in the “monopole” field of the SU(2) skyrmion. The model agrees well with the experimental data and uncannily resemble those predicted by quark models. The same model appears to apply equally well for the description of b-flavor baryons. This striking feature is explained in terms of an underlying geometric principle associated with induced gauge fields.

The key ingredients of the bound kaon–soliton model for strange baryons [1] are the SU(2) hedgehog of the u and d flavors, the kaon field in the doublet representation of SU(2) and the Wess–Zumino term. The Wess–Zumino term serves to bind the negative-strangeness (S = −1) kaons to the hedgehog. When the kaon is bound its “isospin” is transmuted to a spin due to the hedgehog, which plays the role of a “magnetic monopole” in a manner first pointed out by Jackiw and Rebbi [2] and independently by Hasenfratz and ’t Hooft [3]. The fine structure effects depend on the other part of the effective interaction, the short-range behavior of which is specific to the model lagrangian density. When the short-distance behavior is described by a hidden gauge symmetric lagrangian, which incorporates the strong vector mesons, the model describes the strange baryons surprisingly well [4].

In this paper, we argue that similar conditions are met with charmed and bottom flavor baryons. Here we will focus on the charmed baryons, making a brief remark on the bottom flavor baryons at the end.

As in ref. [4], we will assume that the soliton is formed of the pion coupled to U(2) vector mesons (i.e., the ρ, ω) introduced as hidden gauge bosons in the Higgs mode. This implies that the O(Nc) term will be the same as in the strange baryons. The D mesons are introduced in the doublet form in analogy with the kaons as

\[
\Phi_D = \begin{pmatrix} D^+ \\ D^0 \end{pmatrix}.
\]

This doublet has charm quantum number C = +1, whereas the C = −1 doublet is \(\Phi_D' = (D^-, D^0)\). The doublet structure is of course inherited from the u and d quark doublet which forms the pseudoscalar boson. One may
incorporate this together with the s quark in the SU(4) flavor group, in which case the D mesons will belong to the generators $\lambda_{9,10,11,12}$. For the present purpose it is, however, simpler to treat $q_1D$ and $K$ as independent doublets, because the coupling between the $K$ and $D$ can be safely ignored in the quadratic expansion of the meson field that we believe to be reliable. In this approximation the equation of motion of the D meson in presence of the Wess–Zumino term and the effective potential $V_{\text{eff}}$ are formally identical to those for the $K$, with the D-meson mass replacing the kaon mass. This turns out to be the only quantity which depends crucially on the length scale relevant in the problem, heavier mesons probing shorter wavelength degrees of freedom. We expect this $O(N_c^0)$ contribution to the mass of the charmed (and bottom) baryons to be poorly described by the hidden gauge symmetry lagrangian involving only low-mass vector mesons ($\rho$, $\omega$, $A_1$, etc.).

The $O(N_c^{-1})$ hyperfine splitting is, on the other hand, not expected to be very sensitive to the short-distance degrees of freedom. The reason for this is a deep issue, the details of which will be discussed in a separate paper [5]. The point, briefly, is as follows. As in the case of the kaon–soliton system, the isospin–spin transmutation can be viewed as an effect of a Berry phase intrinsic to the system that has two scales: one “slow” degree of freedom (i.e., rotation of the hedgehog) and one “fast” degree of freedom (i.e., vibrational excitation of the mesons) which corresponds to a Born–Oppenheimer situation [6,7]. The symmetry group $SU(2)$ then implies that there will also be induced nonabelian gauge fields, which in turn generate a generic hyperfine splitting. To the extent that this is largely of geometric nature [6,7], we expect that the hyperfine splitting should be correctly described for light as well as massive systems.

Following this general argument, it is clear that the mass formula for charmed baryons with no strangeness quantum numbers is identical to that for the strange baryons. It suffices to substitute the subscript $c$ for $s$ everywhere with $c_m = +1$ for $S = -1$. The situation is slightly more complicated when both the $c$ and $s$ quarks come into play. Suppose that we have $n_1$ mesons of species 1 with energy $\omega_1$ and hyperfine splitting parameter $c_1$ (as defined in ref. [4]) bound in the orbital with $j_1 = l_1 \pm \frac{1}{2}$ and $n_2$ mesons of species 2 with $\omega_2$, $c_2$ in $j_2$. The generalized mass formula is then

$$M(I, J, n_1, n_2, J_1, J_2, J_m) = M_{\text{sol}} + n_1 \omega_1 + n_2 \omega_2 + M_{\text{rot}},$$

$$M_{\text{rot}} = \frac{1}{2} \Theta^{-1} \left[ I(I+1) + (c_1 - c_2) [c_1 J_1(J_1 + 1) - c_2 J_2(J_2 + 1)] + c_1 c_2 J_m(J_m + 1) - J_m(J_m + 1) \right] \left( J_1(J_1 + 1) - J_2(J_2 + 1) \right) \left( J_m(J_m + 1) \right),$$

where $I$ is the isospin, $J$ the total angular momentum (sum of the “spin” carried by the mesons and the soliton angular momentum equal to $I$), $J_i = l_i + j_i$, $i = 1, 2$ and $J_m$ is given by $J_m = J_1 + J_2$, ..., $|J_1 - J_2|$. Here $M_{\text{sol}}$ is the soliton mass and $\Theta$ the moment of inertia, both of which are $O(N_c)$ and determined solely by the SU(2) (light-quark) sector. The appearance of $J_m$ in labeling the states is due to the fact that in the case of different flavors, Bose symmetry does not restrict the total angular momentum of the mesons to its maximal value. Note that in a formal $N_c$ counting, the meson energy term $n_1 \omega_1 + n_2 \omega_2$ is $O(N_c^0)$ and the “rotational” mass $M_{\text{rot}}$ is $O(N_c^{-1})$. It is easy to verify that in the case of one heavy flavor, (2) reduces to the formula given in ref. [4].

Given a specific lagrangian with parameters determined in the meson sector, the quantities appearing in (2), namely, $M_{\text{sol}}$, $\omega_1$, $\omega_2$, $c_1$, $c_2$ and $\Theta$, are all calculable. The mass formula is, however, of a lot more general validity than what a specific lagrangian might imply. Therefore we first treat these quantities as parameters to be fixed by the experimental masses of $N(939)$, $\Delta(1232)$, $\Lambda(1116)$, $\Sigma(1193)$, $\Lambda_c(2285)$ and $\Sigma_c(2453)$. The parameter values determined in this way are:

$$M_{\text{sol}} = 866 \text{ MeV}, \quad \Theta = 1.01 \text{ fm}, \quad \omega_1 = 223 \text{ MeV}, \quad \omega_2 = 1418 \text{ MeV}, \quad c_1 = 0.604, \quad c_2 = 0.140.$$ (3)

Here $\omega_1 \equiv \omega_K$ and $\omega_2 \equiv \omega_D$ represent the energies of the $K$ and $D$ mesons in the ground state. These may be considered as “empirical” values. The predictions for the lowest even-parity states built with p-wave mesons

"
are given in table 1 together with the experimental values and quark model [8,9] and lattice calculation [10] results. The agreement with quark model predictions (and available experimental values [11]) is surprisingly good. It is remarkable that there is a one-to-one correspondence in all levels and that they agree to within a few percent. This agreement essentially reflects the validity of the key assumptions that enter in the structure of the model as stated at the beginning.

Let us now see how well one can understand those parameters with a given effective lagrangian. For this, we take the simplified hidden gauge symmetry lagrangian studied in ref. [4]. This lagrangian is a highly truncated one in that even for three flavors, no axial vector fields are included: Only the nonet vector fields are included with, however, the massive (strange) vectors (i.e. the K*’s) integrated out. To apply this lagrangian to charm (and bottom) quark systems, we will simply add terms identical in form except the c (and b) flavor replacing the s flavor. As noted before, we ignore the coupling between the heavy flavors and hence the contributions to the mass will be additive to O(N c). With the meson parameters F = 186 MeV and g = 5.85 (where g is the gauge coupling), the calculated values are

\[ M_{\text{sol}} = 1475 \text{ MeV}, \hspace{1em} \Theta = 0.82 \text{ fm}, \hspace{1em} \omega_1 = 165 \text{ MeV}, \hspace{1em} \omega_2 = 1005 \text{ MeV}, \hspace{1em} c_1 = 0.71, \hspace{1em} c_2 = 0.21. \]  

(4)

The soliton mass is, as always in the Skyrme model with the empirical value for F, too high. No fully satisfactory remedy to this difficulty is known. For the moment, we will, as was done in ref. [4], simply shift the ground state mass down by 627 MeV in order to have the nucleon mass at its empirical value. The hyperfine coefficients c_1 and the moment of inertia \( \Theta \) are in fair agreement with the “empirical” values above. The major disagreements are in the fine structure splitting, at the O(N c) level. As noted in ref. [4], \( \omega_1 \) is too low by about 50 MeV;

\[ \text{The Callan–Klebanov mass formula [1], } \Delta - N = \Sigma^* + \frac{1}{2} \lambda_c, \text{ predicts the as-yet-unmeasured mass } \Sigma^* = 2500 \text{ MeV, consistent with the calculated value 2494 MeV and the quark-model results.} \]

Table 1

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<td>?</td>
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\( ^1 \) The Callan–Klebanov mass formula [1], \( \Delta - N = \Sigma^* + \frac{1}{2} \lambda_c \), predicts the as-yet-unmeasured mass \( \Sigma^* = 2500 \text{ MeV, consistent with the calculated value 2494 MeV and the quark-model results.} \)
this of course is not significant at the level of accuracy aimed at here. The most striking deviation – which we in 
a way expected – is in \( \omega_2 \). The D meson is overbound by 412 MeV. The situation is even more dramatic in b-
quark systems where B mesons are found to be overbound by more than 1 GeV. There is a simple – and plausible –
explanation for this phenomenon which we will sketch below. For the moment, we will simply add the differ-
ences and take

\[
\omega_1 = 223 \text{ MeV}, \quad \omega_2 = 1418 \text{ MeV}.
\] (5)

With this shift in the centroid energy, one obtains practically the same results as the “empirical” ones in table 1.

It is intriguing that the mass formula (2) of the simple model gives results that are quite close to those ob-
tained with quark models. One can easily see, in comparing with the quark model mass formulae [12], that eq. (2) 
reproduces the quark-model mass formulae obtained from a general two-body quark–quark interaction. On 
the other hand, it does not strictly satisfy those which require additional assumptions. For instance, such rela-
tions as the Gell-Mann-Okubo mass formula

\[
2(N+E) = 3A + Z, \quad (6)
\]

\[
2(N+\Xi) = 3\Lambda + \Sigma, \quad (7)
\]

and the equal spacing rule

\[
\Xi^* - \Sigma^* = \Omega - E^* = \Xi^* - \Sigma^* = \Sigma^* - \Delta \quad (8)
\]

are not identically satisfied by eq. (2). With the parameter values (3) the deviation is however quite small: we 
have, for eq. (6), \( 4556 = 4541 \) in MeV, for eq. (7), \( 9382 = 9308 \) and for eq. (8), \( 155 = 153 = 146 = 138 \). This 
provides a compelling support to the thesis that the bound soliton–meson picture closely mimics the quark 
model.

An immediate consequence of our specific model with the prescribed shift of the centroids is the unambiguous 
and nontrivial prediction of excited states containing one charmed quark as there are more than one bound D-
meson states in each channel \(^{61} \). In table 2, the predicted energies of some \( I=0, l \) low-lying excitations relative 
to the lowest \( \Lambda_c \) are listed and compared with the quark-model predictions of refs. [9,13]. The agreement with 
the quark models is again remarkable.

We now discuss the short-range effects associated with the binding energy of the doublet meson occuring at 
\( O(N_c^0) \). While the hyperfine structure, which is seen to be correctly described in all heavy-quark systems, in-
volves a relatively long-wavelength-induced gauge (Berry) potential as discussed in ref. [5], the binding effect 
in question involves the Wess–Zumino term that results from integrating out quark degrees of freedom. For 
light-quark baryons, the resulting Wess–Zumino term contains, in addition to Goldstone bosons, the usual strong 
vector mesons \( (\rho, \omega, \Lambda_c, \ldots) \). For heavy-quark baryons, corresponding massive vector mesons \( (e.g., \ D^* \) in 
the charm-quark case) should, however, appear in the Wess–Zumino term. In the binding of a D meson to the 
soliton, the meson wavefunction, which has short Compton wavelength, probes the deep interior of the hedge-
hog and therefore we expect that the massive vectors will play a more crucial role in the process. There are two 
possible consequences of this. The first one is that the massive vectors might develop mean-field components in 
the hedgehog background field and hence contribute to the classical energy of the hedgehog. We do not know 
whether this would modify significantly the soliton mass (which is typically 600 MeV too high in the usual 
treatment) and if so, in which direction. The second point on which we have some understanding and which we 
can actually check by calculation is that the massive pseudoscalar field is more strongly influenced by the mas-
sive vector fields than the light vectors.

In order to see how the massive vectors modify the binding, we have simply replaced the light vector meson

\(^{62} \) Our model lagrangian predicts, in the notation \( (j, l, n) \), bound \( C=+1 \) D states in the channels: \( (1/2, 1, 1), (1/2, 1, 2), (1/2, 0, 1), \ (1/2, 0, 2), (3/2, 2, 1), (3/2, 2, 2), (3/2, 1, 1), (3/2, 1, 2), (5/2, 3, 1), (5/2, 2, 1). No bound \( C= -1 \) states are predicted. It is possible that some of the \( n=2 \) states and \( j=5/2 \) are not really bound for the reason given in the text.
Table 2
Low-lying excited states of \( \Lambda_c \) and \( \Sigma_c \) (containing one charm quark and no strange quark) calculated with the model of ref. [4]. The energies (in MeV) are given relative to the "ground state" \( \Lambda_c (2285) \). The states labelled by a, b, c, ..., g correspond to the same (intrinsic) bound states, those within the same intrinsic state being rotational excitations.

<table>
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<th>( I )</th>
<th>( J )</th>
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in the equation of motion for the doublet meson by the corresponding massive ones, namely \( D^*(2010) \) and \( B^*(5330) \). The doublet energies so obtained are

\[
\omega_{D^*} = 1265 \text{ MeV}, \quad \omega_{B^*} = 4509 \text{ MeV.}
\] (9)

These are quite close to the "empirical" values. It is clearly tempting to claim that we have correctly isolated the key degrees of freedom, which operate at the respective length scales.

We conclude with the following three tentative but intriguing observations. In the b-quark spectrum, the details of which will be reported elsewhere, we obtain a large number of bound intrinsic states that define the quantum numbers of the b-quark baryons in a way paralleling the charm-quark baryon structure. With the centroid of the b flavor identified according to the argument presented above, the excitation spectrum for \( \Lambda_b \) and \( \Sigma_b \) up to about 1 GeV above \( \Lambda_b (5425) \) agrees, with only four exceptions, with that of the quark model [13]. This is perhaps not so surprising if one believes, as we do, that the hyperfine effect is governed by induced gauge field effects. Furthermore with the account of massive vectors such as \( B^* \), it is even possible to predict correctly the centroid. The simple calculation indicates that this is highly likely. Naturally a fully self-consistent calculation that includes the massive vectors on the same footing as the light vectors would be called for.

One fascinating aspect about our results which also presents an apparent paradox is the uncanny resemblance to the quark-model predictions. Consider the most striking case, the \( \Omega_{ccc} \). To describe this in quark models, one only needs massive charm quarks, a confining potential and color-magnetic interactions: No light-quark degrees of freedom are required. In other words, the \( \Omega_{ccc} \) could exist even if there were no u and d quarks in the universe. In contrast, the soliton description requires the presence of the hedgehog of the light-quark flavors for getting correct quantum numbers as well as hyperfine splittings. We do not understand how the two pictures can be logically related. But we conjecture that there is an intricate duality.

\[ \text{\#3 There are only preliminary data on B*. We will take the best-fit value given in ref. [11], } m_{B^*} \approx 5330 \text{ MeV. In fact, the result for } \omega_{B^*} \text{ is not sensitive to the } B^* \text{ as long as it is higher than 5 GeV.} \]

\[ \text{\#4 With this } \omega_{B^*}, \text{ we get } \Lambda_b^0 \approx 5358 \text{ MeV, close to the experimental value } \approx 5425 \text{ and the quark model prediction 5585 [13]. Without the massive vector, this is overbound by about 1.3 GeV. The details of the B-soliton spectrum will be reported in a separate paper.} \]
In our opinion, the real surprise is that such a simple picture works at all. One explanation which we find to be plausible and exciting is that there is a hierarchy of induced gauge structures operative at various length scales: at high energy – or short distance – it is encoded in the Wess–Zumino term, which at $O(N^0)$, reflects those quark degrees of freedom which are “integrated” out. At $O(N^{-1})$ or hyperfine level, it emerges in a Born–Oppenheimer approximation in the form of the isospin–spin transformation and induced (Berry) potentials. This hierarchy is reminiscent of Nambu’s idea [14] of “tumbling” interactions expressed in terms of Landau–Ginzburg-type (or $\sigma$) models, the domain of application ranging from the Higgs sector down to the IBM interaction in heavy nuclei.

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References

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