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Title: Hybrid Baryon Signatures

Author(s): Philip R. Page, T-16,
Theoretical Division, Los Alamos National Lab.

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HYBRID BARYON SIGNATURES

P. R. PAGE[a]
Theoretical Division, MS-B283, Los Alamos National Laboratory,
Los Alamos, NM 87545, USA
E-mail: prp@lanl.gov

We discuss whether a low–lying ***hybrid baryon*** should be defined as a three quark – gluon bound state or as three quarks moving on an excited adiabatic potential. We show that the latter definition becomes exact, ***not*** only for very heavy quarks, but also for specific dynamics. We review the literature on the signatures of hybrid baryons, with specific reference to strong hadronic decays, electromagnetic couplings, diffractive production and production in ψ decay.

1 What is a hybrid baryon?

Historically a low–lying hybrid baryon was defined as a three quark – gluon composite. However, from the viewpoint of the Lagrangian of Quantum Chromodymanics (QCD) this definition is non–sensical. This is because gluons are massless, and hence there is no reason not to define a hybrid baryon, for example, as a three quark – two gluon composite. Neither is one possibility distinguishable from the other, since strong interactions mix the possibilities. The place where this definition of a hybrid baryon is most useful is in large Q^2 deep inelastic scattering, where a Fock state expansion of a state can rigorously be defined, and one can at least talk about the three quark – gluon component of such a state.[1] However, in other situations the definition becomes perilous. A case in point is recent work on large N_c hybrid baryons, where their properties depend critically on the fact that the gluon is in colour octet, and hence the three quarks in colour octet, so that the entire state is colour singlet.[2] The bag model circumvents the objections raised against this definition, since gluons become massive due to their confinement inside the bag.[3]

More recently, a low–lying hybrid baryon was defined as three quarks moving on the low–lying excited adiabatic potential.[4] From the viewpoint of QCD this can be a perfectly sensible definition. One can always evaluate the energy of a system of three fixed quarks as a function of the three quark positions,

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called the adiabatic potential. A calculation along these lines has been performed in flux–tube models[5] and a first attempt has been made in lattice QCD.[6] The three quarks are then allowed to move in a three–body equation, typically a non–relativistic Schrödinger equation. Treating a three quark system via this two step process is called the adiabatic or Born–Oppenheimer approximation. Note that it is not appropriate to talk about case where the quarks are actually infinitely heavy, because of the lack of kinetic energy. The adiabatic approximation is expected to become exact if the quark masses are much greater than the scale of the strong interactions Λ_{QCD}. The criterion for the validity of the adiabatic approximation is that the slow degrees of freedom (the quarks) should move much slower than the fast degrees of freedom (the gluons). It is possible to argue that for conventional baryons the relative velocities of quarks behave like the strong coupling constant α_S as the quarks become heavier.[7] Because of the asymptotic freedom of QCD, the quark relative velocities go to zero, ensuring the validity of the adiabatic approximation. In fact, this is the basis for the NRQCD expansion. However, since α_S goes to zero only logarithmically, one may need quarks heavier than the bottom quark for the adiabatic approximation to be valid. Depending on the shape of the adiabatic potential the possibility of an NRQCD expansion for *hybrid* baryons,[8] and hence the validity of an adiabatic approximation, may be in jeopardy.

Here we point out for the first time that for specific dynamics, the adiabatic approximation can be *exact*, even for light quarks, if one redefines the adiabatic potential suitably. We call this the "redefined adiabatic approximation", which employs a "redefined adiabatic potential". It was noted in a flux–tube model that "For light quarks almost all corrections may be incorporated into a redefinition of the potentials. Mixing between [new] potentials is of the order of 1%".[9] We develop the following technique for obtaining the redefined potential. The quark positions are still fixed relative to each other, but the quark and gluon positions are not defined relative to the quark positions, but relative to the centre of mass of the quarks and the gluons, as recently pioneered.[5] For dynamics where the redefined adiabatic approximation is exact one can rigorously define a low–lying hybrid baryon as three quarks moving in the low–lying redefined excited adiabatic potential.

2 Redefined adiabatic approximation

Consider the non–relativistic hamiltonian for three quarks at positions $\mathbf{r}_1, \mathbf{r}_2, \mathbf{r}_3$ and a junction at position $\mathbf{r}_4$. The junction represents the gluons. The classical hamiltonian with a simple harmonic oscillator potential is

$$H = \frac{1}{2}M(\dot{\mathbf{r}}_1^2 + \dot{\mathbf{r}}_2^2 + \dot{\mathbf{r}}_3^2) + \frac{1}{2}m\dot{\mathbf{r}}_4^2 + \frac{1}{2}k((\mathbf{r}_1 - \mathbf{r}_4)^2 + (\mathbf{r}_2 - \mathbf{r}_4)^2 + (\mathbf{r}_3 - \mathbf{r}_4)^2) \quad (1)$$

where M, m and k are the mass of the quarks, the mass of the junction and coupling constant respectively. The first and second terms are the kinetic and potential energy terms respectively.

2.1 Exact solution

Making the variable transformations

$$\boldsymbol{\rho} = \frac{1}{\sqrt{2}}(\mathbf{r}_1 - \mathbf{r}_2) \qquad \boldsymbol{\lambda} = \frac{1}{\sqrt{6}}(\mathbf{r}_1 + \mathbf{r}_2 - 2\mathbf{r}_3)$$
$$\boldsymbol{\sigma} = \frac{2}{\sqrt{3}}\frac{m}{m+3M}(\mathbf{r}_1 + \mathbf{r}_2 + \mathbf{r}_3 - 3\mathbf{r}_3) \quad (2)$$

one obtains the diagonal hamiltonian

$$H = \frac{1}{2}M(\dot{\boldsymbol{\rho}}^2 + \dot{\boldsymbol{\lambda}}^2 + \frac{1}{4}(1+3\frac{M}{m})\dot{\boldsymbol{\sigma}}^2) + \frac{1}{2}k\,(\boldsymbol{\rho}^2 + \boldsymbol{\lambda}^2 + \frac{1}{4}(1+3\frac{M}{m})^2\boldsymbol{\sigma}^2) \quad (3)$$

when working in the overall centre of mass frame. The hamiltonian describes three independent simple harmonic oscillators so that the energies solving the (quantized) Schrödinger equation, are

$$E = (n_\rho + \frac{3}{2})\omega_\rho + (n_\lambda + \frac{3}{2})\omega_\lambda + (n_\sigma + \frac{3}{2})\omega_\sigma \quad (4)$$

where e.g. $n_\rho = n_\rho^x + n_\rho^y + n_\rho^z$, with n_ρ^x, n_ρ^y and n_ρ^z the degrees of excitation in the three space directions. The vibrational frequencies are

$$\omega_\rho{}^2 = \omega_\lambda{}^2 = \frac{k}{M} \qquad \omega_\sigma{}^2 = (1+3\frac{M}{m})\frac{k}{M} \quad (5)$$

The wave functions solving the Schrödinger equation can be denoted by

$$\psi_{n_\rho}(\boldsymbol{\rho})\,\psi_{n_\lambda}(\boldsymbol{\lambda})\,\psi_{n_\sigma}(\boldsymbol{\sigma}) \quad (6)$$

where e.g. n_ρ denotes the set $\{n_\rho^x, n_\rho^y, n_\rho^z\}$.

One can verify that the quark relative velocities $\sim (\frac{k}{Mm^2})^{\frac{1}{4}}$, so that they go to zero if $M \gg \frac{k}{m^2}$. The criterion for the validity of the adiabatic approximation is hence satisfied for large quark masses.

2.2 Adiabatic solution

Fix the quark positions so that $\dot{\mathbf{r}}_1 = \dot{\mathbf{r}}_2 = \dot{\mathbf{r}}_3 = 0$. The potential energy depends only on the differences between positions, and is hence unaffected by fixing the quark positions. The kinetic energy in Eq. 1 is $\frac{1}{2}m\dot{\mathbf{r}}_4^2$, and using $2\sqrt{3}\dot{\mathbf{r}}_4 = -(1+3\frac{M}{m})\dot{\boldsymbol{\sigma}}$ from Eq. 2, the hamiltonian (1) is

$$H_j = \frac{1}{2}\frac{m}{12}(1+3\frac{M}{m})^2\dot{\boldsymbol{\sigma}}^2 + \frac{1}{2}k\left(\boldsymbol{\rho}^2 + \boldsymbol{\lambda}^2 + \frac{1}{4}(1+3\frac{M}{m})^2\boldsymbol{\sigma}^2\right) \tag{7}$$

Only the variable $\boldsymbol{\sigma}$ is dynamical. The energies (adiabatic potentials) are

$$E_j = (n_\sigma + \frac{3}{2})\omega_\sigma^{'} + \frac{1}{2}k(\boldsymbol{\rho}^2 + \boldsymbol{\lambda}^2) \qquad \omega_\sigma^{'\,2} = 3\frac{k}{m} \tag{8}$$

Now allow the quarks to move in their centre of mass frame, so that the quark motion hamiltonian is

$$\begin{aligned} H_q &= \frac{1}{2}M(\dot{\mathbf{r}}_1^2 + \dot{\mathbf{r}}_2^2 + \dot{\mathbf{r}}_3^2) + E_j \\ &= \frac{1}{2}M(\dot{\boldsymbol{\rho}}^2 + \dot{\boldsymbol{\lambda}}^2) + (n_\sigma + \frac{3}{2})\omega_\sigma^{'} + \frac{1}{2}k(\boldsymbol{\rho}^2 + \boldsymbol{\lambda}^2) \end{aligned} \tag{9}$$

The energies are

$$E_q = (n_\rho + \frac{3}{2})\omega_\rho + (n_\lambda + \frac{3}{2})\omega_\lambda + (n_\sigma + \frac{3}{2})\omega_\sigma^{'} \tag{10}$$

It is easy to see that these adiabatic approximation energies only agree with the exact energies (4) when $\frac{M}{m} \gg 1$, i.e. when the quarks are much heavier than the gluons. This is in accord with one's naïve expectation for the validity of the adiabatic approximation.

2.3 Redefined adiabatic solution

We follow the same procedure as for the adiabatic solution, with one critical change. The quark positions are still fixed will respect to each other, but all positions are now defined with respect to the overall centre of mass, before we allow the quarks to move. Define the overall centre of mass as

$$\mathbf{R} = \frac{M(\mathbf{r}_1 + \mathbf{r}_2 + \mathbf{r}_3) + m\mathbf{r}_4}{m+3M} \quad \Rightarrow \quad \dot{\mathbf{R}} = \frac{m}{m+3M}\dot{\mathbf{r}}_4 = -\frac{1}{2\sqrt{3}}\dot{\boldsymbol{\sigma}} \tag{11}$$

Define new coordinates which are the positions of each particle with respect to the overall centre of mass. The time derivatives of these coordinates are

$$\dot{\mathbf{r}}_1' = \dot{\mathbf{r}}_2' = \dot{\mathbf{r}}_3' = -\dot{\mathbf{R}} = \frac{1}{2\sqrt{3}}\dot{\boldsymbol{\sigma}} \qquad \dot{\mathbf{r}}_4' = \dot{\mathbf{r}}_4 - \dot{\mathbf{R}} = -\frac{\sqrt{3}}{2}\frac{M}{m}\dot{\boldsymbol{\sigma}} \tag{12}$$

The kinetic energy in terms of the new coordinates is

$$\frac{1}{2}M(\dot{\mathbf{r}}_1'^2 + \dot{\mathbf{r}}_2'^2 + \dot{\mathbf{r}}_3'^2) + \frac{1}{2}m\dot{\mathbf{r}}_4'^2 = \frac{1}{2}\frac{M}{4}(1+3\frac{M}{m})\dot{\boldsymbol{\sigma}}^2 \tag{13}$$

This kinetic energy combined with the potential in Eq. 1 is

$$H_j = \frac{1}{2}\frac{M}{4}(1+3\frac{M}{m})\dot{\boldsymbol{\sigma}}^2 + \frac{1}{2}k\,(\boldsymbol{\rho}^2 + \boldsymbol{\lambda}^2 + \frac{1}{4}(1+3\frac{M}{m})^2\boldsymbol{\sigma}^2) \tag{14}$$

We used the fact that the potential energy depends only on the differences between positions. The energies (redefined adiabatic potentials) are

$$E_j = (n_\sigma + \frac{3}{2})\omega_\sigma + \frac{1}{2}k(\boldsymbol{\rho}^2 + \boldsymbol{\lambda}^2) \tag{15}$$

and the junction wave functions $\psi_{n_\sigma}(\boldsymbol{\sigma})$. Allowing the quarks to move, the quark motion hamiltonian is

$$H_q = \frac{1}{2}M(\dot{\boldsymbol{\rho}}^2 + \dot{\boldsymbol{\lambda}}^2) + (n_\sigma + \frac{3}{2})\omega_\sigma + \frac{1}{2}k(\boldsymbol{\rho}^2 + \boldsymbol{\lambda}^2) \tag{16}$$

The energies are

$$E_q = (n_\rho + \frac{3}{2})\omega_\rho + (n_\lambda + \frac{3}{2})\omega_\lambda + (n_\sigma + \frac{3}{2})\omega_\sigma \tag{17}$$

which are identical to the exact solution (4). The quark wave functions are $\psi_{n_\rho}(\boldsymbol{\rho})\ \psi_{n_\lambda}(\boldsymbol{\lambda})$. In order to obtain the full wave functions of the system, we take the direct product of the quark and junction wave functions, i.e. $\psi_{n_\rho}(\boldsymbol{\rho})\ \psi_{n_\lambda}(\boldsymbol{\lambda})\ \psi_{n_\sigma}(\boldsymbol{\sigma})$. These are identical to the wave functions of the exact solution (6).

In conclusion, the solution obtained via the redefinited adiabatic approximation is *exact.*

We shall now show how the specific dynamics (1) provide a toy model for the quark model, in the sense that many of the features needed for the validity of quark model, are *exact* in the toy model. Quark models enable the dynamics of quarks, while freezing out the dynamics of gluons. In the toy model this corresponds to claiming that gluons are in the same wave function for all conventional baryons. But this is manifestly the case in the redefined adiabatic approximation. The redefined adiabatic potential (15) is explicitly dependent on quark masses. This corresponds to quark models: the Coulomb interaction is usually postulated to have a coupling that depends on quark masses. On the other hand, the adiabatic potential (8) is not dependent on quark masses.

The wave functions obtained from the adiabatic approximation differ from those in the redefined adiabatic approximation only for the gluons. When one freezes the dynamics of gluons, as in the quark model, one does not notice the difference between these two approximations. Even the functional forms of the gluon wave functions are the same: the wave functions only differ in the dimensionful scale they contain. Thus the forms can be used interchangebly.

When a process is studied that involves more than one redefined adiabatic potential, e.g. hybrid baryon decay to a conventional baryon and meson, care has to be taken to use the redefined adiabatic gluon wave functions instead of the adiabatic ones. The former gluon wave functions are dependent on ω_σ, which is itself dependent on the quark masses.

3 Where can one search for hybrid baryons?

We first outline the results of a recent flux–tube model calculation.[5] The flavour, non-relativistic spin S and J^P of the seven low–lying hybrid baryons are $(N,\Delta)^{2S+1}J^P = N^2\frac{1}{2}^+,\ N^2\frac{3}{2}^+,\ \Delta^4\frac{1}{2}^+,\ \Delta^4\frac{3}{2}^+,\ \Delta^4\frac{5}{2}^+$, where the first two states double. The bag model has the same number of states. The pair $N^2\frac{1}{2}^+,\ N^2\frac{3}{2}^+$ has the same quantum numbers as in the flux-tube model.[3] The other five states in the bag model are $\Delta^2\frac{1}{2}^+,\ \Delta^2\frac{3}{2}^+,\ N^4\frac{1}{2}^+,\ N^4\frac{3}{2}^+,\ N^4\frac{5}{2}^+$. The state $N\frac{1}{2}^+$ was studied in QCD sum rules.[10] The hyperfine interaction moves the Δ states upwards and the N states downwards.[5] Hence there are four low–lying N hybrid baryons with a mass of 1870 ± 100 MeV. This is somewhat higher, but sometimes within errors, of bag models and QCD sum rules which find a mass around 1.5 GeV.[3,10] The wave function sizes are estimated as $\sqrt{\langle\rho^2\rangle} = \sqrt{\langle\lambda^2\rangle} = 2.1$ (conventional baryons) and 2.5 (hybrid baryons) GeV^{-1}. Hybrid baryons are hence larger than conventional baryons.

The following techniques may enable the detection of hybrid baryons:

• *Overabundance of states*: This approach is troublesome, since not even all conventional baryons in the appropriate mass region have been discovered yet. However, hybrid baryon states are likely to be discovered before all conventional baryons in a relevant mass region, and hence need to be studied alongside conventional baryons.

• *Decays*: Except for a QCD sum rule motivated suggestion that hybrid baryons should decay to $N\sigma$,[11] no decay calculations have been performed. However, decay of hybrid baryons to $N\rho$ and $N\omega$ is *a priori* interesting since it isolates states in the correct mass region, without contamination from lower–lying conventional baryons. Study of these decay modes will also yield information on photo– and electroproduction in Hall B and C at Jefferson Lab, since the ρ and ω couple to the photon via vector meson dominance. The process $N\pi \rightarrow$ hybrid $\rightarrow N\omega$ can be studied at the D–line at Crystal Ball E913.

• *Diffractive γN and πN production:* The detection of the hybrid meson candidate $\pi(1800)$ in diffractive πN collisions by VES[12] may indicate that hybrid mesons are producted abundantly via meson–pomeron fusion. If this is the case, one expects significant production of hybrid baryons via baryon–pomeron fusion, i.e. production in diffractive γN and πN collisions.

• *Production in ψ decays:* Naïve expectations are that the gluon–rich environment of ψ decays should lead to dominant production of glueballs, but also signifant production of hybrid mesons and baryons. The large branching ratios[13] $Br(\psi \rightarrow p\bar{p}\omega,\ p\bar{p}\eta') \sim 10^{-3}$ may contain hybrid baryons decaying to $(p, \bar{p})\ \omega$ especially. Recently a $J^P = \frac{1}{2}^+$ 2σ peak at mass 1834^{+46}_{-55} MeV was seen in $\Psi \rightarrow p\bar{p}\eta$.[14]

• *Electroproduction:* In the flux–tube model, i.e. the adiabatic picture of a hybrid baryon, there is the qualitative conclusion that "$ep \rightarrow eX$ should produce ordinary N^*'s and hybrid baryons with comparable cross–sections".[15] However, the conclusions obtained from the three quark – gluon picture of a hybrid baryon is different. For large Q^2 electroproduction, the Q^2 dependence of the amplitudes is summarized in Table 1. Since the photon has both a transverse and longitudinal component, the amplitude for a conventional baryon is expected to dominate that of the hybrid baryon as Q^2 becomes large.[1] For small Q^2 the conclusion agrees with the large Q^2 result for transverse photons, but is more dramatic for longitudinal photons: the amplitude vanishes.[16] It has accordingly been concluded that the (radially excited) conventional baryon is dominantly electroproduced, with the hybrid baryon subdominant to the resonances $S_{11}(1535), D_{13}(1520)$ and Δ as Q^2 increases.[16] The Q^2 dependence of the electroproduction of a resonance can be measured at Jefferson Lab Hall B and an energy upgraded Jefferson Lab. A hybrid baryon is expected to behave different from nearby conventional baryons as a function of Q^2. One needs to perform partial wave analysis at different Q^2. For large Q^2 cross–sections are small, which would make this way of distinguishing conventional from hybrid baryons challenging.

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Table 1: Q^2 dependence of amplitudes for the electroproduction of conventional or hybrid baryons with transverse or longitudinal photons.[1]

	Conventional	Hybrid
Transverse	$1/Q^3$	$1/Q^5$
Longitudinal	$1/Q^4$	$1/Q^4$

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