Vol. 29 (1998) ACTA PHYSICA POLONICA B No 11

EXOTIC MESONS AND STRUCTURE OF SOFT GLUE*

Adam P. Szczepaniak

Physics Department
Indiana University, Bloomington IN 47405, USA

(Received July 1, 1998)

Hybrid meson phenomenology is reviewed and models of soft gluons are discussed in the context of heavy quarkonia.

PACS numbers: 12.39.Hg

1. Introduction

The constituent quark model, even though not well understood from first principles remains as one of the main tools for studying hadronic structure. Despite the existence of a number of states for which the naive constituent picture fails, an unambiguous identification of nonvalence, gluonic or higher Fock quark components, has not been established. However, there is a growing interest in theoretical description of phenomena that could require explicit excitations of soft gluonic modes. Recently the BNL E852 collaboration has announced the discovery of an exotic $J^{PC} = 1^{-+}$ state in the $\eta\pi$ system of the reaction $\pi^- p \to \eta\pi^- p$ with 18 GeV pions [1]. If true, this would be a confirmation of the previous VES results which, however seems to be inconsistent with the earlier KEK results. Since the $\eta\pi$ channel resonates in $J^P = 2^+$ wave forming a strong $a_2(1320)$ meson, the exotic signal is buried deeply in the background rendering partial wave analysis complicated. There is also a possibility of spurious exotic signals from "leakage" of strength from the strong nonexotic wave [2]. Future experiments are clearly needed and other production mechanisms should be explored. In fact, there is a growing experimental effort to build a meson facility at the JLab that would focus on spectroscopy in the $1.5 - 2.5$ GeV mass region [3] and search for possible hybrid signals.

---

* Presented at the Meson'98 and Conference on the Structure of Meson, Baryon and Nuclei, Cracow, Poland, May 26–June 2, 1998.

Phenomenology of hybrids has been developed in the bag, QCD sum rule, constituent glue, and flux tube models [4]. The constituent glue model is based on the observation that in a fixed, noncovariant gauge *e.g.* Coulomb gauge, gluon mass may be generated. A heavy gluon could be treated as another constituent, localized color charge interacting with the quarks through a static potential [5,6]. In the Dynamical Quark Model (DQM) [6] the starting point is the Coulomb gauge QCD Hamiltonian regularized using a cutoff on interaction matrix elements which eliminates couplings between states with free energy difference larger than the cutoff. The advantage of this type of regularization over more conventional, momentum cutoff schemes is that it does not reduce the Fock space and eliminates the problem of small energy denominators. Perturbative expansion in the strong coupling constant is then used to calculate effective interactions which bring back the effects eliminated by the cutoff. The expansion is justified for as long as the cutoff is kept larger than the QCD scale, Λ_{QCD}. This renormalization program has been carried out using a similarity transformation of the Hamiltonian which keeps identical the spectra of the original (cutoff independent) and effective Hamiltonians. The cutoff, effective Hamiltonian is then approximated by replacing the nonabelian Coulomb interaction by a Cornell-type potential and diagonalized nonperturbatively. A single gluon dispersion relation is determined self consistently in the mean field approximation using the BCS ansatz for the vacuum. Because of the confining interaction a single gluon becomes infinitely massive, but in a colorless bound state *e.g.* in a presence of an octet $Q\bar{Q}$ pair or another gluon, it behaves as a constituent with a mass of approximately 800 MeV at low momenta. The spectrum of two gluon bound states, glueballs obtained in this approach has been calculated in Ref. [5].

An alternative approach to soft gluon dynamics is to assume that the interaction between quarks is mediated by a flux of chromoelectric field whose quantum excitations correspond to the spectrum of soft gluon modes. This is different from the constituent picture since now the $Q\bar{Q}$ interactions is mediated by a large (infinite) number of intermediate degrees of freedom, beads. In the original approach of Isgur ad Paton [7], the beads were assumed to be colorless scalars. We have recently extended the model to include color and allow for interactions of colored fields with the beads and quark sources [8]. Our ultimate goal is to use hybrid phenomenology to discriminate between various models. In the following I will discuss the adiabatic $Q\bar{Q}$ potentials in the context of the two models described above.

2. Adiabatic $Q\bar{Q}$ potentials

To be able to focus on the structure of gluonic modes, I am going to eliminate quark dynamics by focusing on (infinitely) heavy quarks. Gluonic interactions between heavy quarks typically occur on a time scale much shorter then motion of the quarks and in the adiabatic approximation one describes the interaction between the slow (quark) and fast (gluon) degrees of freedom in terms of static potentials. The static $Q\bar{Q}$ potential in DQM is given by the spectrum of the constituent gluon interacting with the $Q\bar{Q}$ pair through the Cornell-type potential modified to incorporate the appropriate color factors [9],

$$[K+V]\left|\boldsymbol{r},\Lambda_Y^{PC}\right\rangle = E(r)\left|\boldsymbol{r},\Lambda_Y^{PC}\right\rangle\,, \tag{1}$$

Here K is the kinetic energy of a single gluon as discussed above and V is the residual, attractive interaction between gluon and the $Q\bar{Q}$ pair. The states are classified according to their transformation properties under rotation along the $Q\bar{Q}$ axis, Λ, combined gluonic parity and charge conjugation, PC and reflection in the plain containing the $Q\bar{Q}$ pair, Y. The adiabatic energies for $\Lambda = 0$ and $\Lambda = 1$, the Σ and Π surfaces in diatomic molecule notation, are shown in Fig. 1 and are compared to the corresponding surfaces obtained in the flux tube model. In the flux tube model gluonic field between quarks is described in terms of small, nonrelativistic oscillations of a 1-dim string. The corresponding Hamiltonian can be diagonalized in a space of string excitations, phonons coupled to static $Q\bar{Q}$ sources.

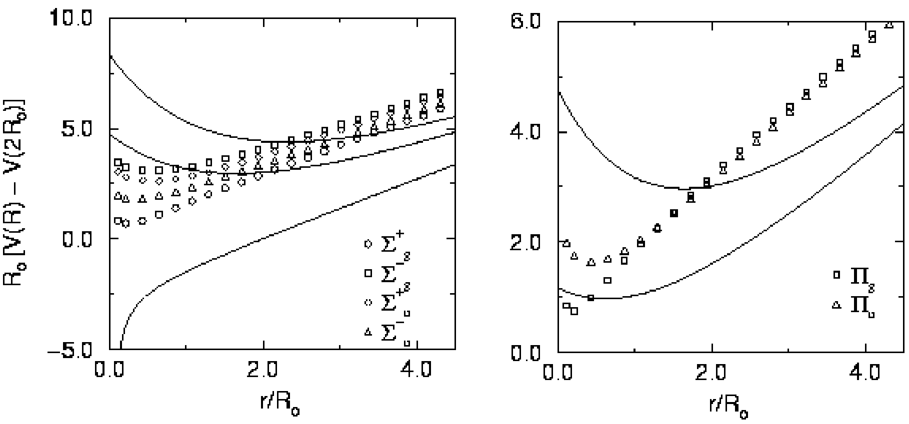

Fig. 1. Σ (left) and Π (right) adiabatic potentials in units of $r_0 \sim 0.5$ fm. Dynamical quark model (symbols) *versus* flux tube model (solid lines) predictions.

The energies are of the form

$$E(r) = br + \frac{N\pi}{r}(1 - \mathrm{e}^{-fb^{1/2}r}) \qquad (2)$$

with $f \sim 1$ being a phenomenological factor introduced to soften the small r behavior and $N = \sum_{m=1} m(n_{m+} + n_{m-})$. The latter represents the total number of right-handed (n_{m+}) and left-handed (n_{m-}) transverse phonon modes weighted by the phonon momentum. The flux tube model predicts the first excited $\Sigma_g'^{+}$ to be split by $N = 2$ from the ground state, and two degenerate Σ_u^+ and Σ_u^- potentials at $N = 3$. The lowest Π state is predicted to be the Π_u in agreement with the lattice results [10] as shown in Fig. 2. Other surfaces are also well described particularly for large interquark separations. Inspecting the graphs in Figs 1,2 it is clear that the constituent model fails to reproduce the adiabatic potential at large interquark separation, r. The success of the flux tube model on the other hand indicates the need for a large number of intervening degrees of freedom.

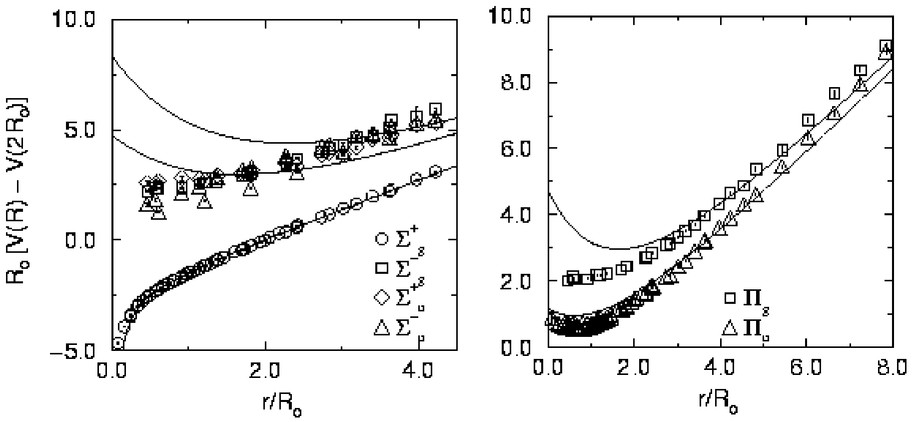

Fig. 2. Σ (left) and Π (right) adiabatic potentials. lattice (symbols) *versus* flux tube model (solid lines) predictions.

A similar conclusion is reach when studying spin-dependent corrections to the adiabatic approximation. Ordering of P-wave heavy quarkonia requires the spin–orbit potential to vanish at large separation between the heavy quarks. Since the spin dependent potentials can be expressed in terms of matrix elements of gluonic $\boldsymbol{E}$ and $\boldsymbol{B}$ fields it is possible to calculate them in the two models discussed here. In Ref. [8] it was shown that an infinite number of intervening degrees of freedom between the quarks is necessary to accommodate this behavior of the spin dependent interactions. In phenomenological models it is often assumed that $Q\bar{Q}$ interactions can be parameterized by an underlying covariant current–current interaction. After nonrelativistic reduction the corresponding heavy quark potentials may be derived. Vanishing of the spin–orbit potential at large quark separation can be obtained if the underlying, covariant interaction is chosen between two Lorentz scalar currents. This is in odds with the notion of the potential as an element of the Hamiltonian, *i.e* fourth component of a Lorentz vector, chiral symmetry and assumption of common confinement for both $Q\bar{Q}$ and QQQ states. As discussed in Ref. [8] the proper way is to start with a vector-like confinement and then to eliminate dynamical gluons. Then the new effective interaction eliminates the spin–orbit potential and makes it look like originating from an underlying scalar kernel. Even if the full $Q\bar{Q}$ kernel could be derived by eliminating an infinite number of Dyson–Schwinger equations there is no reason why it should be local in time, and therefore the whole notion of the Lorentz structure of confinement is not a well defined concept. It makes sense to talk about potentials in the limit of large quark masses but these may have little to do with the relativistic, and time-nonlocal Dyson-Schwinger kernel.

This research was supported by the US Department of Energy.

REFERENCES

[1] D.R. Thompson *et al.*, E852 Collab., *Phys. Rev. Lett.* **79**, 1630 (1997).

[2] A. Dzierba, in NCSU. JLab Workshop on Hybrids and Photoproduction Physics.

[3] For more information go to http:\ dustbunny.physics.indiana.edu\d̃zierba\ Eightplus

[4] T. Barnes, F.E. Close, E.S. Swanson, *Phys. Rev.* **D52**, 5242 (1995) and references therein.

[5] A.P. Szczepaniak, E.S. Swanson, C.-R. Ji, S.R. Cotanch, *Phys. Rev. Lett.* **76**, 2011 (1996).

[6] A.P. Szczepaniak, E.S. Swanson, *Phys. Rev.* **D55**, 1578 (1997).

[7] N. Isgur, J.Paton, *Phys. Rev.* **D31**, 2910 (1985).

[8] A.P. Szczepaniak, E.S. Swanson, *Phys. Rev.* **D55**, 3987 (1997).

[9] E.S. Swanson, A.P. Szczepaniak, hep-ph/9804219.

[10] K. Juge, J. Kuti, C.J. Morningstar, hep-lat/9709131