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Heavy Quarkonium Hybrids *

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Abstract

We report on a recent investigation on heavy quarkonium hybrids that goes beyond the usual Born-Oppenheimer approximation by including not only the mixing between nearby hybrid states but also the mixing with quarkonium states. We use a systematic effective field theory framework based on NRQCD together with lattice QCD inputs. Short and long distance constraints from weak coupling pNRQCD and the QCD effective string theory are also employed. We calculate the quarkonium and hybrid spectrum for charmonium and bottomonium, and estimate a number of decay widths. Most of the isospin zero XYZ resonances fit in our spectrum either as quarkonia or as hybrid states. The mixing of hybrid states with quarkonia produces enhanced spin symmetry violations, which are instrumental to understand certain decays. We also present new results on the hyperfine splittings.

Keywords: Heavy Quarkonium, Hybrids, NRQCD, pNRQCD, QCD string

1. Introduction

Exotic hadrons, namely those beyond quark-antiquark or three quark (antiquark) states, have been contemplated as a theoretical possibility since the early days of QCD [1]. The interest on exotic hadrons has recently experienced a revival due to the pletora of charmonium, and some bottomonium, resonances, the so called XYZ states, discovered in the last decade that do not easily fit in the quark model spectrum (see [2, 3] for recent reviews). The fact that the charm and bottom quark masses (m_c and m_b) are much larger than the typical hadronic scale, $\Lambda_{\rm QCD}$, has allowed to unambiguously identify tetraquark [4–6] and pentaquark [7] states. We shall focus here on heavy charmonium and bottomonium hybrids, namely $c\bar{c}$ and $b\bar{b}$ states with a non-trivial gluon content, with the aim to understand at least part of the spectrum of isospin zero XYZ states.

Since $m_c, m_b \gg \Lambda_{\rm QCD}$, the heavy quarks move slowly so that they see an instantaneous potential as the effective interaction. In early string models the hybrid potential was associated with the excitations of the string [8, 9]. The use of the Born-Oppenheimer (BO) approximation to obtain the hybrid potential was initiated in [10] within the bag model approach. The first lattice QCD calculation calculation of the hybrid potential was carried out in [11]. More recently, the BO approximation has been revisited in relation with the XYZ states [12]. It has also been incorporated into an effective field theory framework in [13] (see also [14]) elaborating on the weak coupling regime of pNRQCD [15, 16], and in [17, 18] elaborating on the strong coupling regime of pNRQCD [16, 19, 20]. The following sections are based on ref. [18], except for the section on the hyperfine splitting, which presents new material [21].

2. Quarkonium

In order to set the scale of the hybrid spectrum it is important to have the quarkonium spectrum calculated in the same framework. Since lattice calculations exist for both the quarkonium and the hybrid potentials (see Fig. 1), we shall fix the single arbitrary constant

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of all these potentials by fitting to the charmonium and bottomonium spectrum. The shape of the quarkonium potential evaluated on the lattice (Σ_g^+ in Fig. 1) is well described by the Cornell potential, which has the short and long distance behavior expected from QCD perturbation theory and the QCD effective string theory (EST) [22, 23] respectively,

$$V_{\Sigma_g^+}(r) \approx -\frac{\sigma_g}{r} + \kappa_g r + E_g^{Q\bar{Q}}\,, \qquad (1)$$

where we take,

$$\sigma_g = 0.489\,, \qquad \kappa_g = 0.187\,\mathrm{GeV}^2\,, \qquad (2)$$

and we obtain from the comparison with the experimental spectrum,

$$E_g^{c\bar{c}} = -0.242\,\mathrm{GeV} \qquad E_g^{b\bar{b}} = -0.228\,\mathrm{GeV}\,. \qquad (3)$$

Note that $E_g^{Q\bar{Q}}$ should be flavor independent, and indeed $E_g^{c\bar{c}}$ and $E_g^{b\bar{b}}$ agree within a 6%. The spectrum obtained with the potential above is displayed in Tables 1 and 2.

3. Hybrids

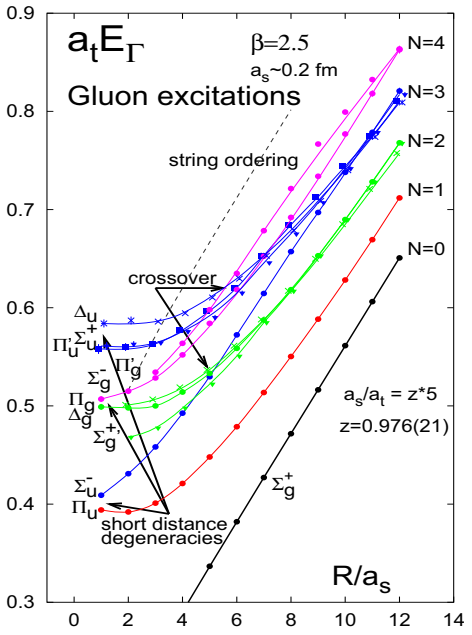

Figure 1: Energy spectrum in the static limit for $n_f = 0$ [24].

The hybrid potentials together with the quarkonium potential (Σ_g^+) are displayed in Fig. 1. The labels correspond to the representations of the $D_{\infty h}$ group, the group of a diatomic molecule. At short distances all the hybrid potentials must approach the repulsive Coulomb potential of the color octet configuration, as perturbation theory dictates. Furthermore, the states should gather in short distance multiplets according to the rotational group [16]. At long distance they must approach the behavior dictated by the QCD EST, namely to the same linear potential as the quarkonium case (Σ_g^+) with a subleading $1/r$ behavior that depends on the string state [23]. Notice that the hybrid potentials, unlike the quarkonium one, have a classical minimum, which must sit at $r \sim 1/\Lambda_{\rm QCD}$ (there is no other scale available). Hence the small energy fluctuations about this minimum have a size $\sqrt{\Lambda_{\rm QCD}^3/m_Q}$, which is parametrically smaller than $\Lambda_{\rm QCD}$. Hence, if we are only interested in the lower lying states for each potential, we are in a situation similar to the strong coupling regime of pNRQCD [16], in which the scale $\Lambda_{\rm QCD}$ can be integrated out. This means that, in a leading approximation, we can ignore the interaction with any other hybrid or quarkonium state with an energy $\gtrsim \Lambda_{\rm QCD}$ above or below the low lying states. In Fig. 1, we observe that the short distance degeneracies are already noticeable close to the minima. Hence, it is natural to chose the degrees of freedom of the effective theory as a wave function field that describe the corresponding short distance multiplet. We are going to restrict ourselves to the lower lying hybrid multiplet, namely that formed by (Σ_u^-, Π_u). At short distances, this wave function field corresponds to a quark-antiquark in a color octet state together with a chromomagnetic field that makes the whole operator color singlet [16]. Hence, we choose a vectorial wave function matrix $\mathbf{H}(\mathbf{0}, \mathbf{r}, t)$ with the same symmetry transformations as that operator. Namely, it transforms as $\mathbf{H} \to h_1 \mathbf{H} h_2^\dagger$, $h_1, h_2 \in SU(2)$ under spin symmetry and as follows under parity, time reversal and charge conjugation,

$$\begin{aligned} P:&\quad \mathbf{H}(\mathbf{R},\mathbf{r},t) \to -\mathbf{H}(-\mathbf{R},-\mathbf{r},t) \\ T:&\quad \mathbf{H}(\mathbf{R},\mathbf{r},t) \to -\sigma^2\mathbf{H}(\mathbf{R},\mathbf{r},-t)\sigma^2 \\ C:&\quad \mathbf{H}(\mathbf{R},\mathbf{r},t) \to -\sigma^2\mathbf{H}^T(\mathbf{R},-\mathbf{r},t)\sigma^2\,, \end{aligned} \qquad (4)$$

As a consequence, the P and C quantum numbers of a hybrid state with quark-antiquark orbital angular momentum L and quark-antiquark spin S are,

$$P = (-1)^{L+1}\,,\; C = (-1)^{L+S+1}\,. \qquad (5)$$

The Hamiltonian at leading order (BO approximation) is chosen such that the projection of $\mathbf{H}$ to $\mathbf{r}$ evolves with $V_{\Sigma_u^-}$ and the projection orthogonal to $\mathbf{r}$ with V_{Π_u} in the