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Final state interaction effects in near threshold enhancement of the $p\bar{p}$ mass spectrum in B and J/ψ decays

D. R. Entem and F. Fernández

Grupo de Física Nuclear, IUFFyM, Universidad de Salamanca, E-37008 Salamanca, Spain
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A near threshold enhancement in the $p\bar{p}$ invariant mass spectrum has been reported by the BELLE, BES, and *BABAR* collaborations for several B and J/ψ decays. This enhancement has been interpreted as a narrow baryonium state $X(1835)$. We investigate its nature using a $p\bar{p}$ interaction derived from a constituent quark model. This interaction does not show any $p\bar{p}$ bound state but a ${}^3P_0$ resonance. We show that $p\bar{p}$ final state interaction can reproduce the mass dependence of the $p\bar{p}$ mass spectrum close to the threshold observed in different B and J/ψ decays.

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I. INTRODUCTION

The observation of a near threshold narrow enhancement in the $p\bar{p}$ invariant mass spectrum from radiative $J/\psi \rightarrow \gamma p\bar{p}$ decay by the BES Collaboration [1] has renewed the interest for the $N\bar{N}$ interaction and its possible baryonium bound states. A fit of the experimental data with an S-wave Breit-Wigner resonance function found a peak mass at $M = 1859 \pm 6$ MeV, which is below the $p\bar{p}$ threshold, whereas a P-wave fit gives a peak mass very close to the threshold at $M = 1876 \pm 0.9$ MeV. The photon polar angle distribution is consistent with $1 + \cos^2\theta$ which suggests that the total angular momentum is very likely to be $J = 0$. So this structure may have quantum numbers $J^{PC} = 0^{-+}$ or $J^{PC} = 0^{++}$ which do not correspond to any known meson resonance. More refined analysis gives a mass of $M = 1835$ MeV and therefore the state is sometimes denoted as $X(1835)$ [2].

The first observation of similar enhancements was reported by the BELLE Collaboration in the decays $B^{\pm} \rightarrow K^{\pm}p\bar{p}$ [3] and $B^0 \rightarrow D^0 p\bar{p}$ [4], although only very recently the *BABAR* measurements of $B^0 \rightarrow \bar{D}^0 p\bar{p}$ and $B^0 \rightarrow \bar{D}^{*0}p\bar{p}$ [5] have shown that the phase space corrected $p\bar{p}$ invariant mass distribution has the same behavior for the two decay channels and very similar to the one found by BES and BELLE collaborations for the $J/\psi \rightarrow \gamma p\bar{p}$, $B^+ \rightarrow K^+ p\bar{p}$, and $B^0 \rightarrow D^0 p\bar{p}$, respectively. Out of this systematic is the $J/\psi \rightarrow \pi^0 p\bar{p}$ which does not show any enhancement at threshold. In addition to threshold enhancements, the pK^- correlation in $B^- \rightarrow K^- p\bar{p}$ decay has been measured by BELLE [6].

These results suggest that the enhancement of the $p\bar{p}$ mass spectrum is related with the dynamics of the $p\bar{p}$ pair and is weakly constrained by the decay vertex.

The simplest interpretation of the experimental $J/\psi \rightarrow \gamma p\bar{p}$ is a baryonium bound state [7–10]. A rich spectrum of $N\bar{N}$ bound states were predicted in the past based on the idea that G-parity transformation of the standard NN interaction models turns the repulsive short range part of the NN potential into an attractive short range $N\bar{N}$ one. Although, as showed by Myhrer and Thomas [11], part of the bound state spectrum produced by the real part may be washed out when annihilation is taken into account, most of the meson-exchange based potentials, like the Paris or Bonn potentials, predicts more than one quasibound state, mainly in S waves. However, no clear experimental indication of these states has been found until now. Using a recent version of the Paris potential, Loiseau and Wycech [10] explain the near threshold enhancement in $J/\psi \rightarrow \gamma p\bar{p}$ as due to the only near threshold quasibound state in the ${}^{11}S_0$ $p\bar{p}$ partial wave that this potential has. However, as pointed out by the authors, the state is strongly dependent of the model parameters and should be confirmed by other experiments. Moreover, this explanation is not compatible with the recent enhancement found in the $B^0 \rightarrow \bar{D}^0 p\bar{p}$ and $B^0 \rightarrow \bar{D}^{*0}p\bar{p}$ decays which would required bound states in the ${}^3S_1$ $p\bar{p}$ partial waves.

A simple short distance argument can explain the threshold baryon-antibaryon enhancements in three body decays and the suppression of two body decays. For the $B^- \rightarrow K^- p\bar{p}$, Suzuki [12] argued that the gluon emitted in the strong penguin decay $b \rightarrow sg^*$ can be almost on shell and produces the enhancement at low $p\bar{p}$ invariant mass. However, this picture predicts that the p and K^- should be emitted in opposite directions when the measured pK^- correlations turns out to be the opposite [6].

Leaving aside more exotic, although possible, explanations, like glueballs or quark fragmentation effects, the other natural way to understand this phenomena is as final state interaction (FSI) effects. The $J/\psi \rightarrow \gamma p\bar{p}$ decay has been the study for several authors with different approaches. Kerbikov *et al.* [13] and Bugg [14] using a complex S-wave $p\bar{p}$ scattering length are able to reproduce the shape of the $p\bar{p}$ invariant mass distribution, although it is not clear the reliability of the scattering length approach in this energy range. Zou and Chiang [15] uses a K matrix formalism but its interaction includes only one pion exchange avoiding other important parts of the $N\bar{N}$ interaction, namely, annihilation. More realistic interactions are used in the calculations of Ref. [16]. These authors uses the Julich $N\bar{N}$ model and show that the mass dependence of the $p\bar{p}$ spectrum close to the threshold can be reproduced by the S-wave $p\bar{p}$ final state interaction in the isospin $I = 1$ state. Also B decays have been studied in this framework [17]. The dominance of the $I = 1$ channel together with the unobservation of the enhancement in the $J/\psi \to \pi^0 p\bar{p}$ channel cast some doubts on the interpretation of the data.

In the present paper we analyze the near threshold enhancement in the $p\bar{p}$ invariant mass distribution using a $N\bar{N}$ interaction derived from a constituent quark model. The use of quark degrees of freedom to describe $N\bar{N}$ interactions presents several advantages over the conventional meson-exchange potentials. First of all the short range part of the interaction can be derived directly from the corresponding NN one without using new parameters. Second, quark annihilation diagrams provides the real part of the annihilation potential, also without additional parameters. Moreover, taken into account that the way to generate the short range and tensor interactions is completely different in quark based models and in meson-exchange models, one may wonder if a quark based $N\bar{N}$ potential can provide different results for the J/ψ and B decays. The potential we use for the $N\bar{N}$ interaction has been described in Ref. [18]. This model is based in the constituent quark model developed in Refs. [19,20] and is able to describe the $N\bar{N}$ scattering data and the energy shifts in protonium, in particular, the enhancement observed in the $^3P_0$ state energy shift. The model presents no quasibound states but a near threshold resonance in the $^3P_0$ $I = 0$ partial wave which is responsible of the enhancement in the protonium energy shift in this particular wave [18].

The paper is organized in the following way. In Sec. II we briefly describe the $N\bar{N}$ potential model we use. In Sec. III we analyze the method to calculate the FSI effects. We compare our results with the different experimental data in Sec IV. The paper ends with some concluding remarks.

II. THE MODEL

Constituent quark models are based on the assumption that the constituent quark mass is generated by the spontaneous breaking of the original $SU(3)_L \otimes SU(3)_R$ symmetry of the QCD Lagrangian for (almost) massless quarks at some momentum scale [21].

The picture of the QCD vacuum as a dilute medium of instantons [22] explains nicely such symmetry breaking, which is the most important nonperturbative phenomenon for hadron structure at low energies. Quarks interact with fermionic zero modes of the individual instantons in the medium modifying the light quark propagator which acquires a momentum dependent mass which drops to zero for momentum higher than the inverse of the average instanton size $\bar{\rho}$.

The momentum dependent mass acts as a natural cutoff of the theory. Moreover, the Goldstone boson modes appearing as a consequence of the chiral symmetry breaking provides an interaction between the constituent quarks.

Beyond the chiral symmetry breaking scale one expects the dynamics being governed by QCD perturbative effects. There are consequences of the one-gluon fluctuations around the instanton vacuum and we take it into account through the qqg coupling. Since it is not allowed any quark-antiquark exchange between N and $\bar{N}$, the $N\bar{N}$ interaction from one-gluon exchange is very different from the NN one. As the gluon carries color, it cannot be exchanged between colorless states and only contributes through annihilation diagrams. A similar contribution appears when a quark-antiquark pair annihilate into a pion which give rise to a new quark-antiquark pair after propagation. These diagrams contribute to the real part of the so-called $N\bar{N}$ annihilation potential. Besides these processes, $N\bar{N}$ can annihilate into a huge number of different channels, mainly meson channels. These annihilation processes are very complicated and very difficult to describe completely in terms of quark degrees of freedom. That is the reason why they are usually described using phenomenological optical potentials like

$$V_{q\bar{q}}^{\mathrm{Anh}}(\vec{q}) = iW_i e^{-(q^2 b'^2/3)}, \tag{1}$$

where W_i gives the strength of the interaction and b' its range. Finally, constituent quarks are confined into hadrons. The only well-established indications we have about the nature of this interaction is provided by lattice studies which show that $q\bar{q}$ systems interact at short distances by a linear potential which mimic the effects of a one-dimensional color flux tube. Spontaneous creation of virtual $q\bar{q}$ pairs may give rise to a breakup of the color flux tube screening the linear potential. This potential, which is necessary to describe hadron structure, does not contribute to the $N\bar{N}$ interaction as far as we describe baryons and antibaryons as singlet color clusters, but avoids that N and $\bar{N}$ collapse under the interaction described above. Explicit expressions for all these potentials and the value of the parameters are given in Refs. [19,20]. All but the last two parameters are fixed from the NN sector, whereas W_i and b' are fitted to the total annihilation cross section for the $p\bar{p}$ system in Ref. [18].

Once the microscopic model is fixed we use the resonating group method to derive the $N\bar{N}$ interaction in the same way as we did in the NN case. The wave functions for the baryon (antibaryon) states are

$$\psi_B = \phi_B(\vec{p}_{\xi_1}, \vec{p}_{\xi_2})\chi_B \xi_c[1^3], \tag{2}$$

where χ_B is the spin-isospin wave function, ξ_c is the color wave function,

$$\phi_B(\vec{p}_{\xi_1}, \vec{p}_{\xi_2}) = \left[\frac{2b^2}{\pi}\right]^{3/4} e^{-b^2 p_{\xi_1}^2} \left[\frac{3b^2}{2\pi}\right]^{3/4} e^{-(3b^2/4)p_{\xi_2}^2} \tag{3}$$

is the orbital wave function, b is the parameter related with

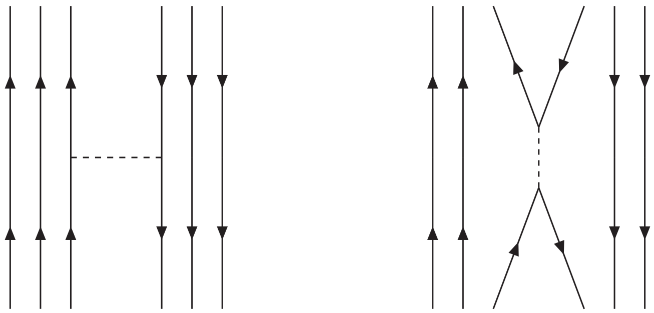

FIG. 1. Diagrams that contribute to the $N\bar{N}$ potential.

the size of the baryon (antibaryon), and $\vec{p}_{\xi_1}$, $\vec{p}_{\xi_2}$ is the Jacobi momenta of the baryons.

The N and $\bar{N}$ wave functions are the same provided that we use for the spin-isospin of the $\bar{N}$ for the G-parity transformed.

As mentioned before, only direct diagrams contribute to the $N\bar{N}$ potential. as seen in Fig. 1. For the first diagram we calculate the interaction as in the NN case but without exchange diagrams. It is interesting to note that since no exchange diagram is present, the interaction is local and from the orbital part we only get a form factor.

III. MESON DECAY RATES AND THE FSI EFFECTS

Owing the evidence that no baryonium bound states appear in our calculation [18], we will now study possible final state interaction effects in the decays of B and J/ψ mesons as an explanation of the near threshold enhancement observed in the data. After averaging over the spin states and integrating over the angles, the differential decay rate of the process $X \rightarrow y p\bar{p}$ can be written as

$$\frac{d\Gamma}{dM} = \frac{\lambda^{1/2}(m_X^2, M^2, m_y^2)\sqrt{M^2 - 4m_p^2}}{2^6\pi^3 m_X^2}|A|^2, \tag{4}$$

where $\lambda(x, y, z)$ is the Kallén function, M is the invariant mass of the $p\bar{p}$ system, and A is the total $X \rightarrow y p\bar{p}$ reaction amplitude. Assuming for A a constant value, one obtains the so-called phase space distribution. An elementary calculation shows [16] that close to the $p\bar{p}$ threshold the data deviate from the phase space distribution indicating the possible influence of final state interaction effects in the $p\bar{p}$ system. A simple way to single out the energy dependence of this FSI effect comes from Watson and Migdal [23]. These authors suggested that the energy dependence of the decay amplitude is given by the energy dependence of the on shell $p\bar{p}$ T matrix. Therefore, the decay amplitude can be factorized in terms of an elementary production amplitude and the T matrix corresponding to the $p\bar{p}$ scattering. Let us outline the derivation of this formula in order to establish its validity. We start from the distorted wave Born approximation for the amplitude of the process

$$A = M + MG^{p\bar{p}}T, \tag{5}$$

where M is the production amplitude, $G^{p\bar{p}}$ is the propagator of the $p\bar{p}$ system and T its scattering amplitude. As pointed out by Watson, if the interaction is short range, M depends weakly on the energy, and we can write $M = \tilde{M}k^L$ where $\tilde{M}$ is a constant, k is the relative $N\bar{N}$ momentum and L its angular momentum. With this assumption we take the correct momentum dependence for the production amplitude at threshold.

For an uncoupled partial wave this equation is written as

$$\begin{aligned} A &= \tilde{M}\left(p^L + \int k^{2+L}dk\frac{1}{E - \frac{k^2}{2\mu}}T(E; k, p)\right) \\ &= \tilde{M}\Big(p^L(1 - i\pi\mu pT(E; p, p)) \\ &\quad + \mathcal{P}\int k^{2+L}dk\frac{2\mu}{p^2 - k^2}T(E; k, p)\Big), \end{aligned} \tag{6}$$

where $\mu = m_p/2$ is the $N\bar{N}$ reduced mass, p the on shell relative $N\bar{N}$ momentum, $\mathcal{P}$ denotes a principal value, and all amplitudes are those of the selected partial wave.

Using the expression of the K matrix in terms of scattering amplitude $T(E; p, p)$

$$K = \frac{T(E; p, p)}{1 - i\pi\mu pT(E; p, p)}, \tag{7}$$

where K is the on shell $p\bar{p}$ K matrix element, one can write the production amplitude as

$$\begin{aligned} A &= \tilde{M}\left(p^L\frac{T(E; p, p)}{K} + \mathcal{P}\int k^{2+L}dk\frac{2\mu}{p^2 - k^2}T(E; k, p)\right) \\ &= \tilde{M}T(E; p, p)\left(\frac{p^L}{K} + \mathcal{P}\int k^{2+L}dk\frac{2\mu}{p^2 - k^2}\frac{T(E; k, p)}{T(E; p, p)}\right), \end{aligned} \tag{8}$$

Now using the effective range expansion

$$\frac{p^{2L}}{K} \sim -\frac{1}{a_L} + \frac{1}{2}r_L p^2 + \ldots, \tag{9}$$

we find

$$\begin{aligned} A &= \tilde{M}\frac{T(E; p, p)}{p^L}\Big(-\frac{1}{a_L} + \frac{1}{2}r_L p^2 + \ldots \\ &\quad + 2\mu p\mathcal{P}\int \bar{k}^{2+L}d\bar{k}\frac{1}{1 - \bar{k}^2}\frac{T(E; k, p)}{T(E; p, p)}\Big), \end{aligned} \tag{10}$$

with $\bar{k} = \frac{k}{p}$.

In this expression all the off shell effects are contained in the principal value of the integral. Assuming that these effects are small and the effective range expansion is a good approximation, one recovers the Watson-Migdal prescription, giving a FSI factor

$$F = \left| \frac{T(E; p, p)}{p^L} \right|^2. \tag{11}$$

The above formula shows that the energy dependence of the decay amplitude is given by F although its absolute value depends on the neglected off shell effects and the unknown constant $\tilde{M}$. Indeed, it is only required that the principal value integral $\mathcal{P}$ has a weak energy dependence compared to the T matrix.

For S waves the Watson-Migdal prescription is a good approximation at moderate energies. However for P waves it may be adequate to use the complete expression

$$F = \left| p^L - i\pi\mu p^{L+1} T(E; p, p)) + \mathcal{P} \int k^{2+L} dk \frac{2\mu}{p^2 - k^2} T(E; k, p) \right|^2. \tag{12}$$

Now F depends on the half off shell $N\bar{N}$ T matrix which is experimentally unknown and theoretically very model dependent. That is the reason why some authors neglect the principal value, using the FSI factor

$$F = \left| p^L \frac{T(E; p, p)}{K} \right|^2. \tag{13}$$

This approximation avoids the difficulties of the Watson-Migdal approach and still only depends on on shell quantities which are free from the flaws described above.

An important consequence of the Watson-Migdal approach is some kind of "universality" on the energy dependence of the $X \to y p\bar{p}$ amplitude which only depends on the $p\bar{p}$ pair and is weakly constrained neither by the decay particle nor by the emitted particle. We will see in the next section that this is in fact the case for the B and J/Ψ decays which suggest the use of the Watson-Migdal approach.

IV. RESULTS

Let us first compare the three different approximations explained above. In Figs. 2–6 we show the F factor for $I = 0$ and $I = 1$ and for the partial waves ${}^1S_0$, ${}^3S_1$, ${}^3P_0$, ${}^3P_1$, and ${}^1P_1$, respectively. The solid line represents the Watson-Migdal (WM) approach, the dashed line the result for Eq. (12) (approximation A), and the dashed-dotted line from Eq. (13) (approximation B). One finds enhancement at threshold in almost all partial waves for the WM approximation. The only exception is the $I = 1$ ${}^1P_1$ partial wave in which neither WM nor A or B approximations show any enhancement. This partial wave will be relevant for the following discussion. As expected, the WM results are similar to the other two approaches in S waves but are different in P waves. The results for approximations A and B are similar except in the ${}^3P_0$ due to the resonance in this partial wave.

We have analyzed the $J/\Psi \to \gamma p\bar{p}$ and $J/\Psi \to \pi^0 p\bar{p}$ decays together with $B^0 \to \bar{D}^0 p\bar{p}$, $B^0 \to \bar{D}^{*0} p\bar{p}$, and $B^+ \to K^+ p\bar{p}$ decays. As shown in Ref. [5], the phase space corrected $p\bar{p}$ invariant mass distributions are the same for all decays but in the $J/\Psi \to \pi^0 p\bar{p}$ one.

Let us first analyze the $J/\Psi \to \gamma p\bar{p}$ reaction. The J^{PC} conservation limits the number of $p\bar{p}$ final states. The allowed states are listed in Table I and only a few possibilities exist for each reaction channel. In this decay isospin is not defined for $p\bar{p}$; therefore, one can expect that the physical result may be a mixture of $I = 0$ and $I = 1$. Coming back to Fig. 2 one can compare the results of our

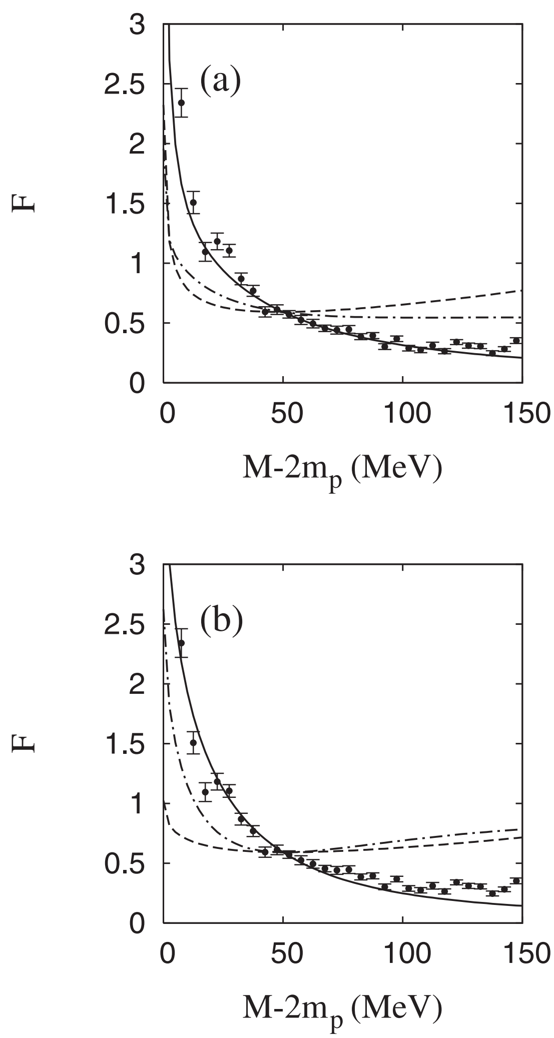

FIG. 2. Final state interaction factor in arbitrary units for the ${}^1S_0$ (0^{-+}) partial wave in the $I = 0$ channel (a) and $I = 1$ channel (b). Data points are for the reaction $J/\Psi \to \gamma p\bar{p}$ from [1,16,24]. The solid line shows the result coming from the Watson-Migdal approach Eq. (11); dashed line from Eq. (12); and dashed-dotted from Eq. (13). All lines are normalized to the experimental value at $M - 2m_p = 50$ MeV.

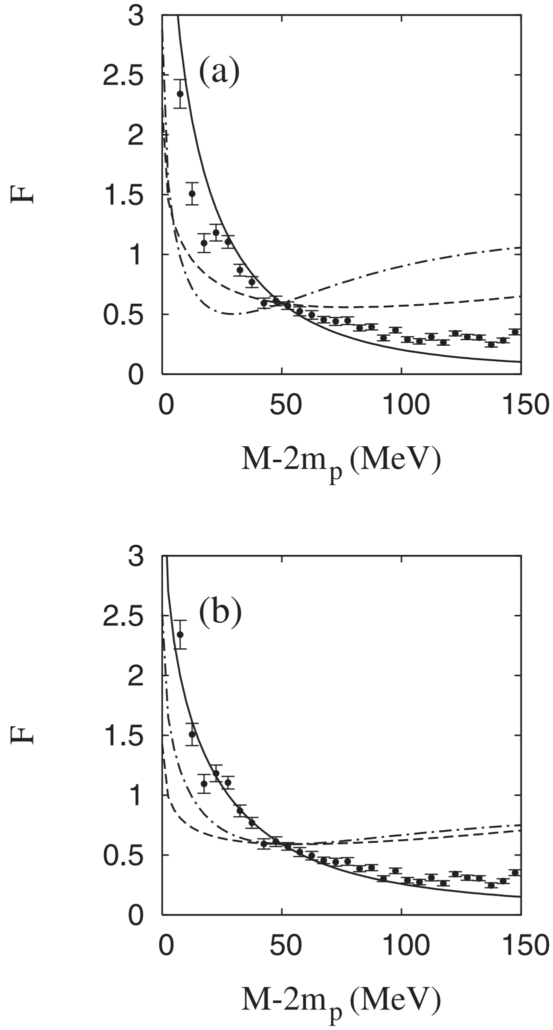

FIG. 3. Final state interaction factor for the $^3S_1$ (1^{--}) partial wave in the $I = 0$ channel (a) and $I = 1$ channel (b). Labels are the same as in Fig. 2. Notice that, although data points are shown, this partial wave is not allowed in the $J/\Psi \to \gamma p\bar{p}$ reaction.

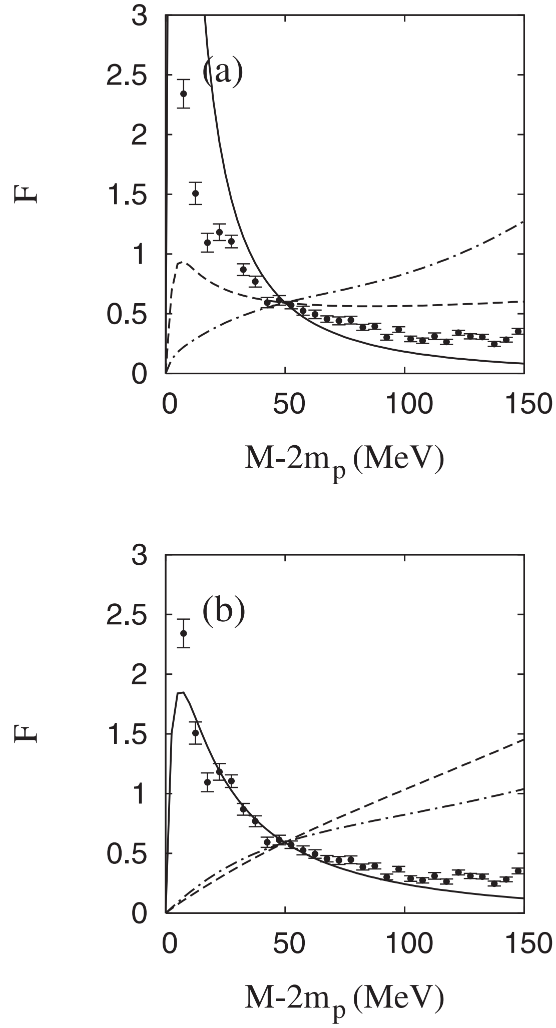

FIG. 4. Final state interaction factor for the $^3P_0$ (0^{++}) partial wave in the $I = 0$ channel (a) and $I = 1$ channel (b). Labels are the same as in Fig. 2.

calculation for the $^1S_0$ partial wave with the $J/\Psi \to \gamma p\bar{p}$ experimental data. Points are the phase space corrected data [16,24] from the BES collaboration [1]. All lines are normalized to the experimental value at $M - 2m_p =$ 50 MeV assuming that only this partial wave contributes. Both isospin channels show a similar threshold enhancement that explains the data assuming that S waves dominate the reaction. In Figs. 4 and 5 one also can compare theory and experiment for the $^3P_0$ and $^3P_1$ partial waves. The Watson-Migdal approach shows similar threshold enhancements in agreement with the data.

Recent data from *BABAR* [5,25] show similar threshold enhancement for the reactions $B^0 \to \bar{D}^0 p\bar{p}$, $B^0 \to \bar{D}^{*0} p\bar{p}$, and $B^+ \to K^+ p\bar{p}$. In Figs. 7 and 8 we show the data compared to the FSI factor for S waves. For the first and third reaction only the $^3S_1$ contributes and in the second one also the $^1S_0$ contributes. The fact that in all these reactions the phase space corrected $p\bar{p}$ invariant mass distribution agrees up to a constant indicates that all S waves should have a similar threshold enhancement. This discards the explanation of the enhancement due to possible baryonium states since they should exists in the $^3S_1$ and in at least the $^1S_0$ or $^3P_0$ waves giving similar contributions, which is very unlikely. In order to support this universality we also show data from the reactions $B^0 \to D^- p\bar{p}\pi^+$, $B^0 \to D^{*-} p\bar{p}\pi^+$, and $e^+e^- \to \gamma p\bar{p}$ in which $p\bar{p}$ states are not constrained by symmetry requirements. All data shows a good agreement up to a constant and could be explained using the Watson-Migdal approximation in S waves.

It is interesting to study the $J/\Psi \to \pi^0 p\bar{p}$ decay since it is the only one in which no threshold enhancement is observed. The isospin of the $p\bar{p}$ system is constrained to $I = 1$ and only $^3S_1$ and $^1P_1$ waves contribute up to $J = 2$. Of course the $^3S_1$ will give an enhancement but the

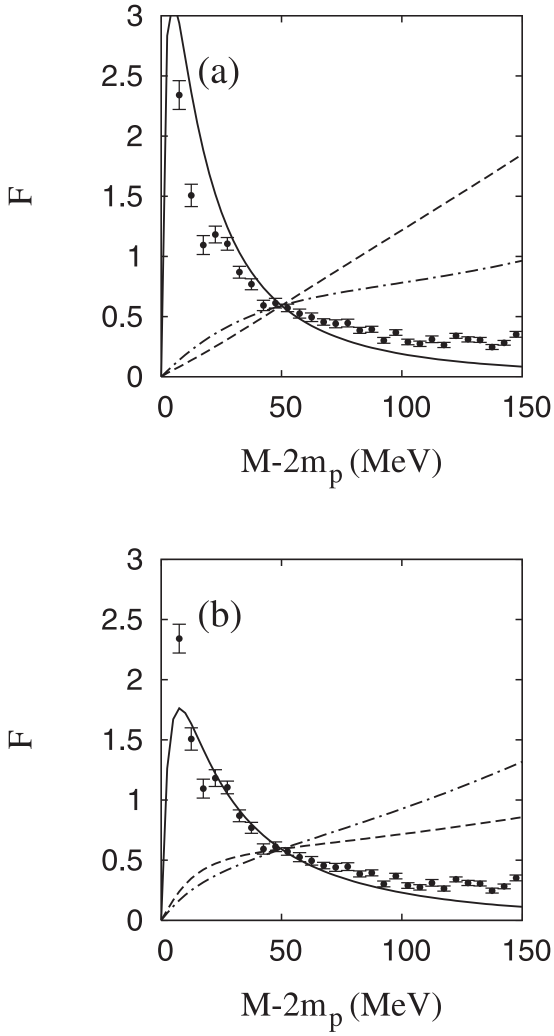

FIG. 5. Final state interaction factor for the $^3P_1$ (1^{++}) partial wave in the $I = 0$ channel (a) and $I = 1$ channel (b). Labels are the same as in Fig. 2.

$^1P_1$ $I = 1$ is in agreement with the data and does not show the enhancement. One explanation why the $^3S_1$ is suppressed was given by Loiseau and Wycech [10]. They assume that $c\bar{c}$ quarks in the J/Ψ state annihilate into an $N\bar{N}$ pair with $I = 0$ and $J^{PC} = 1^{--}$ so it is a $^3S_1$ state. Then a π^0 is emitted via the standard $\pi N\bar{N}$ coupling $\vec{\sigma} \cdot \vec{q}\tau_3$ which requires a final angular momentum $L = 1$ and isospin $I = 1$, so the final state is $^{31}P_1$.

Although we compare the FSI factor with experimental data assuming only one partial wave, the full result should be the sum of the contributions for those allowed by symmetry requirements, the weight of each partial wave given by the decay vertex. Since the $^1S_0$, $^3S_1$, $^3P_0$, and $^3P_1$ waves show the same behavior, the full result will show an enhancement at threshold, and as explained before this also allows an explanation of the different reactions.

The relative sizes of different contributions are essential for other observables. As shown by Suzuki [12], the domi-

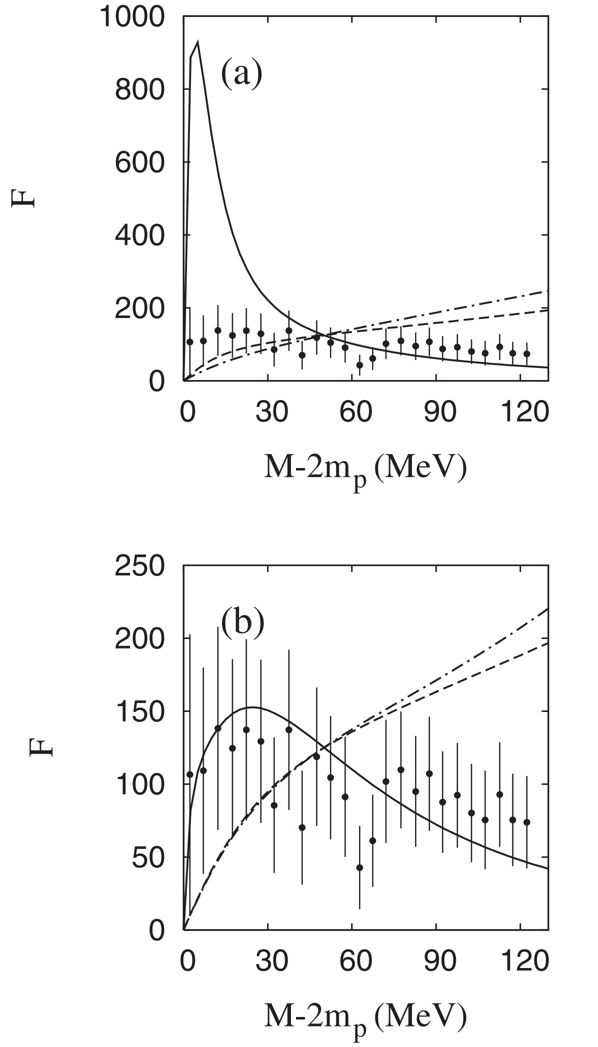

FIG. 6. Final state interaction factor for the $^1P_1$ (1^{+-}) partial wave in $I = 0$ (a) and $I = 1$ (b). Data points are for the reaction $J/\Psi \to \pi^0 p\bar{p}$ from [1,16,24]. The solid line shows the result coming from the Watson-Migdal approach Eq. (11); dashed line from Eq. (12); and dashed-dotted line from Eq. (13).

TABLE I. Lowest $p\bar{p}$ partial waves ($J \leq 1$) that contribute to the different reactions.

Reaction	Partial waves	I
$J/\Psi \to \gamma p\bar{p}$	$^1S_0$, $^3P_0$, $^3P_1$	0, 1
$J/\Psi \to \pi^0 p\bar{p}$	$^3S_1$, $^1P_1$	1
$B^0 \to \bar{D}^0 p\bar{p}$	$^3P_0$, $^3S_1$	0, 1
$B^0 \to \bar{D}^{*0} p\bar{p}$	$^1S_0$, $^1P_1$, $^3P_1$, $^3S_1$	0,1
$B^0 \to D^- p\bar{p}\pi^+$	All	0,1
$B^0 \to D^{*-} p\bar{p}\pi^+$	All	0,1
$B^+ \to K^+ p\bar{p}$	$^3P_0$, $^3S_1$	0,1
$e^+e^- \to \gamma p\bar{p}$	All	0,1

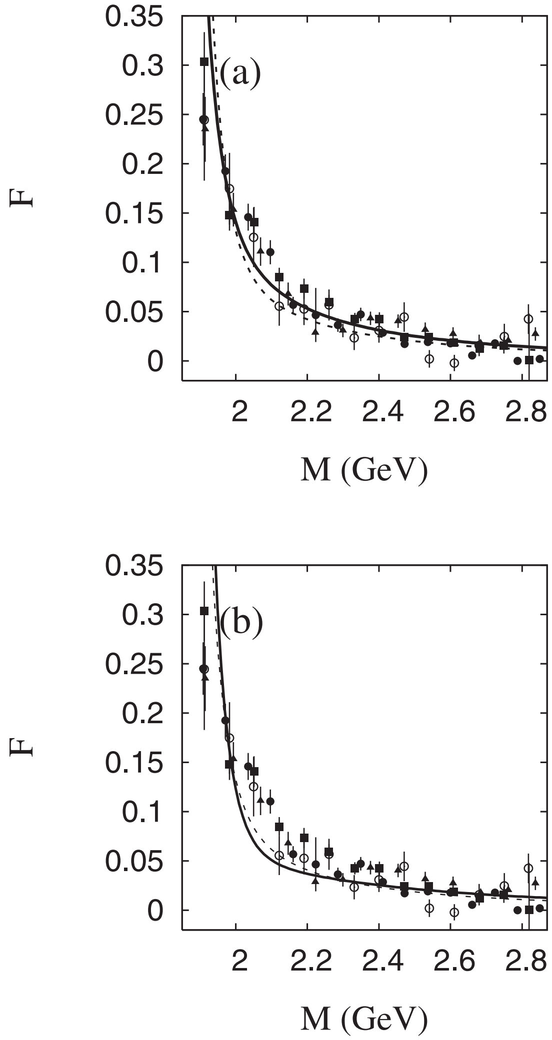

FIG. 7. The phase space corrected $p\bar{p}$ invariant mass distribution for the decay modes: $B^0 \to \bar{D}^0 p\bar{p}$ (triangles), $B^0 \to \bar{D}^{*0} p\bar{p}$ (open circles), $B^0 \to D^- p\bar{p}\pi^+$ (squares), and $B^0 \to D^{*-} p\bar{p}\pi^+$ (filled circles) from *BABAR* [5] compared with the final state interaction factor for the $^1S_0$ (a) and $^3S_1$ (b) channels for $I = 0$ (solid line) and $I = 1$ (dashed line).

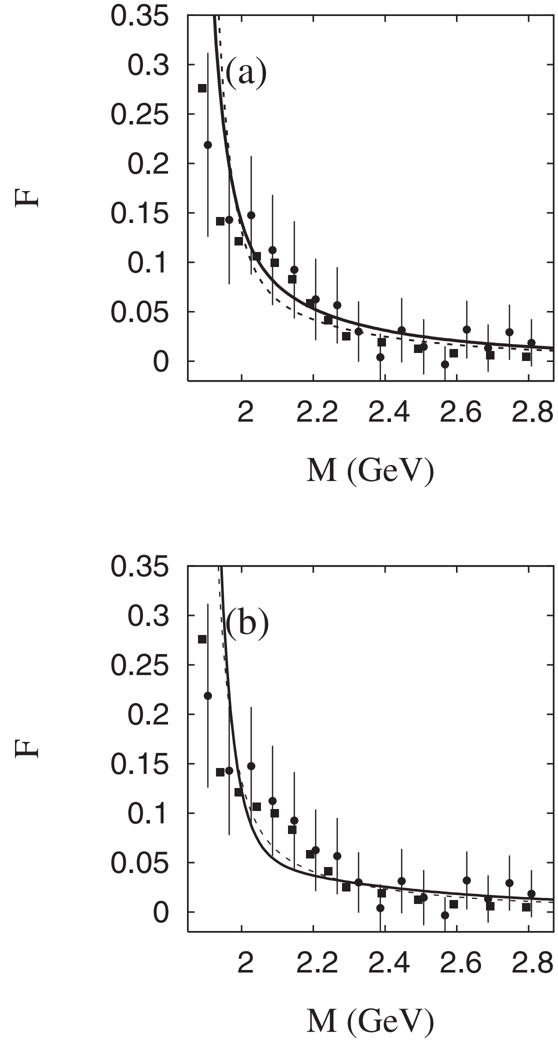

FIG. 8. The phase space corrected $p\bar{p}$ invariant mass distribution for the decay modes: $e^+e^- \to \gamma p\bar{p}$ (squares) from [27] and $B^+ \to K^+ p\bar{p}$ (circles) from [25] compared with the final state interaction factor for the $^1S_0$ (a) and $^3S_1$ (b) channels for $I = 0$ (solid line) and $I = 1$ (dashed line).

nance of only one wave will not explain the pK^- correlation measured in $B^- \to K^- p\bar{p}$. In our calculation we cannot do it since we do not have a model for the decay vertex. Usually the size of the contributions decrease with the corresponding angular momentum, unless a resonance is present. In our calculation the $^3P_0$ wave is enhanced by FSI over other P waves due to the near threshold $p\bar{p}$ resonance in this channel. Therefore, we should expect a stronger interference of the $^3P_0$ and $^3S_1$ which could explain the $K^- p$ correlation [12].

V. CONCLUSIONS

Using a quark model based $N\bar{N}$ interaction which describes the $N\bar{N}$ cross sections and protonium energy shifts, we have investigated the near threshold enhancement found in the $p\bar{p}$ invariant mass spectrum of the $J/\Psi \to y p\bar{p}$ and $B \to z p\bar{p}$ decays, where y states for γ and π^0 and z for $\bar{D}^0$, $\bar{D}^{*0}$, and K^+, recently reported by BES and *BABAR* collaborations. The first clear indication of such enhancement found in the decay $J/\Psi \to \gamma p\bar{p}$ triggered all kind of speculations about the discovery of the long awaited baryonium, bound or quasibound states, or more exotic states like glueballs. However, the discovery of similar enhancements in different reactions, all but the $J/\Psi \to \pi^0 p\bar{p}$ suggests that the enhancement is more related with the $p\bar{p}$ system than with a particular vertex. We have studied several ways to take into account the final state interactions in the decay amplitude. We showed that the near threshold enhancement is understood qualitatively in terms of this state's interactions in the outgoing $p\bar{p}$

system, which in some sense explains the universality of the decay energy dependence.

Surprisingly, the Watson-Migdal prescription for FSI effects gives the best quantitative description of the energy dependence of the decay amplitude even in the case of the $J/\Psi \to \pi^0 p\bar{p}$, which does not show any enhancement but its data can be described by the FSI effects in the $^1P_1$ partial wave, being the $^3S_1$ suppressed by vertex constraints.

Our study shows that the enhancement seen in the $B \to zp\bar{p}$ and $J/\Psi \to yp\bar{p}$ decays can be explained without resort to any $p\bar{p}$ bound state. We conclude that the data are not a signal of baryonium, since they should appear in several partial waves giving similar contributions, and can be explained by final state interaction in the $p\bar{p}$ channel.

Other enhancements have been observed in many baryonic B decay modes having a dibaryon system other than $p\bar{p}$ [26]. As the pion exchange between quarks is the main one responsible for the threshold enhancement, a similar explanation could be valid for all channel decays in which final $B_1\bar{B}_2$ states, involving light quarks, are present. However, many theoretical and experimental efforts are necessary to shed more light on the dynamics of B and J/Ψ decays.

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[1] J. Z. Bai *et al.* (BES Collaboration), Phys. Rev. Lett. **91**, 022001 (2003).
[2] S. Jin (BES Collaboration), Plenary talk at the International Conference on QCD and Hadronic Physics, 2005 Beijing, China.
[3] K. Abe *et al.* (Belle Collaboration), Phys. Rev. Lett. **88**, 181803 (2002).
[4] K. Abe *et al.* (Belle Collaboration), Phys. Rev. Lett. **89**, 151802 (2002).
[5] B. Aubert *et al.* (*BABAR* Collaboration), hep-ex/0607039.
[6] M. Z. Wang *et al.* (Belle Collaboration), Phys. Lett. B **617**, 141 (2005).
[7] A. Datta and P. J. O'Donnell, Phys. Lett. B **567**, 273 (2003).
[8] Chao-Hsi Chang and Hou-Rong Pang, Commun. Theor. Phys. **43**, 275 (2005).
[9] Gui-Jun Ding and Mu-Lin Yan, Phys. Rev. C **72**, 015208 (2005).
[10] B. Loiseau and S. Wycech, Phys. Rev. C **72**, 011001(R) (2005).
[11] F. Myhrer and A. W. Thomas, Phys. Lett. **64B**, 59 (1976).
[12] M. Suzuki, hep-ph/0609133.
[13] B. Kerbikov, A. Stavinsky, and V. Fedotov, Phys. Rev. C **69**, 055205 (2004).
[14] D. V. Bugg, Phys. Lett. B **598**, 8 (2004).
[15] B. S. Zou and H. C. Chiang, Phys. Rev. D **69**, 034004 (2004).
[16] A. Sibirtsev, J. Haidenbauer, S. Krewald, Ulf-G. Meissner, and A. W. Thomas, Phys. Rev. D **71**, 054010 (2005).
[17] J. Haidenbauer, Ulf-G. Meissner, and A. Sibirtsev, Phys. Rev. D **74**, 017501 (2006).
[18] D. R. Entem and F. Fernández, Phys. Rev. C **73**, 045214 (2006).
[19] D. R. Entem, F. Fernández, and A. Valcarce, Phys. Rev. C **62**, 034002 (2000).
[20] A. Valcarce, H. Garcilazo, F. Fernández, and P. González, Rep. Prog. Phys. **68**, 965 (2005); J. Vijande, F. Fernández, and A. Valcarce, J. Phys. G **31**, 481 (2005); J. Vijande, A. Valcarce, F. Fernández, and B. Silvestre-Brac, Phys. Rev. D **72**, 034025 (2005).
[21] A. Manohar and H. Georgi, Nucl. Phys. **B234**, 189 (1984).
[22] D. Diakonov, Prog. Part. Nucl. Phys. **51**, 173 (2003).
[23] K. Watson, Phys. Rev. **88**, 1163 (1952).
[24] J. Haidenbauer *et al.* (private communication).
[25] B. Aubert *et al.* (*BABAR* Collaboration), Phys. Rev. D **72**, 051101 (2005).
[26] Hai-Yang Cheng, Int. J. Mod. Phys. A, **21**, 4209 (2006).
[27] B. Aubert *et al.* (*BABAR* Collaboration), Phys. Rev. D **73**, 012005 (2006).