Final State Interactions in Decays of the Exotic $\pi_1$ Meson

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We analyze the role of final state interactions (FSI) in decay of the lightest exotic meson, with $J^{PC} = 1^{-+} \pi_1$. We use the relativistic Lippmann-Schwinger equation for two coupled $\pi b_1$ and $\pi \rho$ channels. The first one is the predicted dominant decay mode of the $\pi_1$, whereas in the other a narrow $\pi_1(1600)$ exotic signal has been reported by the E852 collaboration. The FSI potential is constructed, based on the $\omega$ meson exchange between the two channels. We find that this process introduces corrections to the $\pi_1$ widths of the order of only a few MeV. Therefore, we conclude that a substantial $\pi \rho$ mode cannot be generated through level mixing.

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I. INTRODUCTION

Exotic mesons may provide a unique experimental handle on gluonic excitations. By definition, spin exotic mesons have quantum numbers, spin ($J$), parity ($P$) and charge conjugation ($C$), which cannot be obtained by combining quantum numbers of the valence quark and antiquark alone. Thus other degrees of freedom, in addition to the valence components, must be present in the leading Fock space component of an exotic meson to generate exotic combinations of $J^{PC}$, e.g. $J^{PC} = 0^{-}, 0^{-}, 1^{-}, 1^{-}, 2^{-}, \cdots$. Since multiquark states are expected to fall apart, it is expected that the leading components of exotic mesons will contain excitation of the gluonic field in the presence of just the valence $Q\bar{Q}$ pair. This is confirmed by lattice simulations \cite{1, 2} which find that the lightest exotic meson has $J^{PC} = 1^{-+}$ and overlaps with the state created from the vacuum by application of operators with $Q, \bar{Q}$ and gluon field (link) operators. The mass of this lightest exotic is expected between 1.6 – 2 GeV \cite{3, 4, 5, 6} with the uncertainty in the lattice estimates coming from chiral extrapolations \cite{7}. There are also estimates of exotic meson properties based on the Born-Oppenheimer approximation which indicate that the gluon, if treated as a constituent particle in the ground state exotic meson, has to carry one unit of orbital angular momentum with respect to the $Q\bar{Q}$ \cite{8}.

Whether exotic mesons have been seen experimentally is an open question. Some tantalizing candidates exist in particular in the $\eta'\pi$ and $\rho\pi$ decay channels \cite{9, 10}. The $J^{PC} = 1^{-+}$ partial wave in the $\eta'\pi$ system measured in $\pi^- p \rightarrow \eta'\pi^- p$ is as strong as the $J^{PC} = 2^{++}$ wave from the $a_2$ resonance decay. Unfortunately, the exotic wave peaks at $M_{\eta'\pi} \sim 1.6$ GeV where there are no other well established resonances decaying to $\eta'\pi$. Therefore extraction of the phase of the exotic wave, needed to establish its dynamical character, depends on model assumptions regarding the behavior of other waves. With the GlueX/Hall D experiment one should be able to clarify this issue.

II. FINAL STATE INTERACTIONS IN $\pi_1$ DECAY

There exist several models of exotic meson decays, and all of them predict a small branching ratio to the $\pi \rho$ channel \cite{11, 12, 13}. It is possible, however, that this is modified by final state rescattering between a meson produced in one channel study together with the analysis of production mechanisms. If current experimental results do indeed correspond to elementary QCD exotic mesons, then there might be a potential discrepancy between model predictions and the experimental data. Most models predict that exotic mesons should primarily decay to "exotic" channels dominated by two meson final states with one of them being an $S$-wave and the other $P$-wave meson \cite{11, 12, 13, 14} e.g. $b_1\pi$, $f_1\pi$, $K_1\pi$, $\pi\rho$, while the experimental sightings come for analysis of "normal" decay channels \textit{i.e.} with two $S$-wave mesons, like the $\eta'\pi$ and the $\rho\pi$ modes. In this paper we examine the effects of final state rescattering between mesons as a potential mechanism for shifting strength from one channel to another.

In a previous paper \cite{12} we constructed normal and hybrid meson states and studied kinematical relativistic effects at the quark and gluon level. In this work we will estimate the size of corrections to the $\pi_1$ decays originating from the meson exchange forces between mesons. Since particle number is not conserved and the particle momenta are of the same order as their masses, this problem should be treated in a relativistic formalism. In Sec. II we apply the Lippmann-Schwinger equation for the two coupled $\pi \rho$ and $\pi b_1$ channels. The first is the channel in which a narrow $\pi_1(1600)$ exotic has been observed \cite{10}, and the latter is the predicted dominant decay mode of the $\pi_1$. Then in Sec. III we describe the computational procedure of solving the resulting integral equations and we present a discussion of the numerical results.
can subsequently decay into $\pi$ and $\omega$, and the $\omega$ can absorb the other $\pi$ and produce a $\rho$, as shown in Fig. 1 The $\omega$ exchange is expected to dominate the mixing potential; this is the dominant decay mode of the $b_1$ meson. A similar mixing effect (mediated by $\rho$ exchange) in connection with the $\eta\pi$ $P$-wave enhancement in the 1400 MeV mass region was considered by Donnachie and Page \[16\].

We will denote the single particle exotic state $\pi b_1$ by $|\alpha\rangle$, and the two-particle states $\pi\rho$ and $\pi b_1$ by Roman letters. We first introduce the matrix elements for the potential $V$,

$$V_{\alpha i} = \langle \alpha |V|i\rangle, \quad V_{ij} = \langle i|V|j\rangle,$$

and similarly for $T$. The elements $V_{\alpha i}$ are just the amplitudes of the corresponding "bare" decays of the $\pi_1$ describing transitions at the quark level, whereas $V_{ij}$ are related to the final state interaction potential. In our state space the Lippmann-Schwinger equation can be written as,

$$T = V + VGT,$$

which sums up the ladder of $\omega$ exchanges and the bare exotic meson exchanges.

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where we again used \( [1 - GV]^{-1} = 1 + Gt \). A resonance production process has the form of \( i \to \pi_1 X \to M_1 M_2 X \), where \( i \) denotes an initial state and \( X \) denotes possible other particles in the final state, assumed not to interact with the resonance \( \pi_1 \), and \( M_1 (M_2) \) is the final state meson (isobar) state. The full amplitude has the form

\[
A(i \to \pi_1 X \to M_1 M_2 X) = P_a G_a T_{ai} = \frac{1}{E - m_{ex} - \Sigma_a} (V + V G t)_{ai}, \tag{10}
\]

where \( P_a \) describes the coupling between the bare exotic and the initial state. Thus in the full amplitude the bare exotic propagator \( (E - m_{ex})^{-1} \) is replaced by the dressed one, \( (E - m_{ex} - \Sigma_a)^{-1} \). Furthermore if \( \Sigma_a \), which is energy dependent, is expanded around the physical exotic mass, one obtains the Breit-Wigner propagator. Here, we are interested in the effective coupling of the exotic to the two-meson channels. This is determined by the last term in Eq. (10). In the absence of FSI, the coupling is determined by the matrix \( V \) which changes to \( V + V G t \) due to re-scattering. We will therefore be studying the reduced \( T \) matrix element,

\[
T_{ai} = V_{ai} + V_{aj} G_{ji} t_{ji}, \tag{11}
\]

with \( t_{ji} \) defined by (6). The diagrammatic representation of Eq. (11) is given in Fig. 6.

Eq. (11) written in a full notation in the rest frame of the \( \pi_1 \) is given by

\[
t(p, q, \lambda, \lambda') = V(p, q, \lambda, \lambda') + \sum_{\lambda''} \int \frac{d^3k}{(2\pi)^3 4\omega_1(k)\omega_2(k)} V(p, k, \lambda, \lambda'') G(k) t(k, q, \lambda'', \lambda'), \tag{12}
\]

where

\[
\omega_i(k) = E(m_i, k), \quad H_0(k) = \omega_1(k) + \omega_2(k),
\]

\[
G(k) = [E - H_0(k) + \iota c]^{-1}, \tag{13}
\]

and \( m_i, i = 1, 2 \) are the masses of mesons in the intermediate two-meson state (related to \( G \)). In the above \( p \) and \( q \) are the relativistic relative three-momenta between two mesons, and the dependence on the center-of-mass momentum has been already factored out. Spins \( \lambda \) refer to either \( b_1 \) or \( \rho \).

We introduce the partial wave potentials by

\[
V_{LL'}(p, q) = \sum_{M, M', \lambda, \lambda', j} \int d\Omega_p d\Omega_q \langle L, M; 1, \lambda | j, j \rangle \langle L', M'; 1, \lambda' | j, j \rangle V(p, q, \lambda, \lambda') Y_{LM}(p) Y_{L'M'}^*(q), \tag{14}
\]

where \( d\Omega_q \) is the element of the solid angle in the direction of the vector \( k \) and \( k = |k| \). Similarly we define

\[
t_{LL'}(p, q). \text{ Substituting } V_{LL'}(p, q) \text{ into (12) gives,}
\]

\[
t_{LL'}(p, q) = V_{LL'}(p, q) + \sum_{L'} \int \frac{k^2 dk}{(2\pi)^3 4\omega_{L', 1}(k)\omega_{L', 2}(k)} V_{LL'}(p, k) G_{L'}(k) t_{L' L'}(k, q), \tag{15}
\]
In our state space we can have $L = 0, 2$ (the relative angular momentum between $\pi$ and $b_1$) or $L = 1$ (between $\pi$ and $\rho$). For a $\pi_1$ we also have $J = 1$. Thus in Eq. (16) we must substitute:

$$m_{L,1} = m_\pi, \ m_{0,2} = m_{2,2} = m_{b_1}, \ m_{1,2} = m_\rho. \quad (17)$$

Parity conservation reduced the number of non-vanishing matrix elements of the final state interaction potential to $V_{10}, V_{12}, V_{21}$ and $V_2$. Moreover, from CP invariance we have $V_{10} = V_{10}^*$ and $V_{12} = V_{21}^*$. The integral equation (15) cannot be solved analytically and one needs to replace it by a set of matrix equations. The details will be given in the next section. When this is done and $t_{LL'}$ are found, we may go back to Eq. (11) which becomes

$$\tilde{a}_L(P) = a_L(P) + \sum_{L'} \int \frac{k^2 dk}{(2\pi)^3(2J + 1)\omega_1(k)\omega_2(k)} \times a_{L'}(k)G(k)_{LL'}t_{LL'}(k,P), \quad (18)$$

where $a_L(P)$ are the partial decay amplitudes defined in Ref. [17]. In order to obtain the corrected widths, they must be replaced by the corrected amplitudes $\tilde{a}_L(P)$.

Finally, we proceed to the form of the final state interaction potential. The dominant effective interactions between the $b_1\pi$ and the $\rho\pi$ channels are expected to originate from $\omega$ exchange. The $b_1 \to \pi\omega$ is the dominant decay channel of the $b_1$ meson and the $\rho\pi$ coupling is also known to be large. The effective Lagrangian for the $\rho\pi\omega$ vertex is

$$L_{\rho\pi\omega} = g_{\rho\pi\omega}\epsilon_\mu\nu\lambda\sigma \partial_\mu\rho_\nu\pi^\lambda P^\sigma, \quad (19)$$

and for the $b_1\pi\omega$ vertex is

$$L_{b_1\pi\omega} = g_{b_1\pi\omega} \partial_\mu b_1^{\mu\omega}\pi^\nu. \quad (20)$$

Here the index $i$ corresponds to isospin. The coupling constant, $g_{\rho\pi\omega}$, lies between 0.01 MeV$^{-1}$ and 0.02 MeV$^{-1}$ and we will take a value 0.014 MeV$^{-1}$ [17], whereas $g_{b_1\pi\omega}$, can be obtained by calculating the width of the decay $b_1 \to \pi\omega$ and comparing with its experimental value of 142 MeV [18]. Consequently, the amplitude for $b_1 \to \pi\omega$ is equal to

$$A = -g_{b_1\pi\omega}\epsilon^i(\lambda_1)\epsilon^i(\lambda_\omega, P), \quad (21)$$

where $P$ is given by

$$E(m_\pi, P) + E(m_\omega, P) = m_{b_1}, \quad (22)$$

and $\epsilon^i(\lambda)$ are the spin-1 polarization four-vectors. The absolute value of the square of this amplitude, summed over $\lambda_\omega$ and averaged with respect to $\lambda_{b_1}$, gives

$$|A|^2 = g_{b_1\pi\omega}^2 (1 + \frac{P^2}{3m_\omega^2}), \quad (23)$$

and using $\Gamma_{s-\omega\pi} = 142 \text{ MeV} \cdot 0.92 = 131 \text{ MeV}$ (the factor 0.92 comes from the partial wave D/S ratio) we get $g_{b_1\pi\omega} = 3650 \text{ MeV}$. Therefore $g_{FS1} = g_{\rho\pi\omega} \cdot g_{b_1\pi\omega}$ is approximately 50.

The final state interaction potential can be obtained from Lagrangians of Eqs. (19) and (20) expressed in momentum space and dressed with the instantaneous $\omega$ propagator,

$$V(p, q, \lambda_1, \lambda_\rho) = g_{FS1}\epsilon_{\mu\nu}\sigma\rho^\mu q^\nu \frac{1}{(p - q)^2 + m_\omega^2} \times \epsilon^\sigma(\lambda_1, -p)\epsilon^\sigma*(\lambda_\rho, q), \quad (24)$$

where $p$ is the momentum of the $\pi$ in the $|\pi b_1\rangle$ state and $q$ is the momentum of the $\rho$. At short distances, (for large values of $p$ and $q$) this potential is singular which is a consequence of treating mesons as elementary particles. Therefore we must regulate this potential with an extra factor that tends to zero for large momenta. We will choose an exponential function

$$e^{-|p|/\Lambda}, \quad (25)$$

where $\Lambda$ is a scale parameter expected to be on the order of the inverse of the meson radius $\Lambda \sim 0.5 - 1 \text{ GeV}$.

**III. NUMERICAL RESULTS**

The Lippmann-Schwinger integral equation of Eq. (15) corresponds to outgoing wave boundary conditions. This means that the singularity of the term $G(k)$ is handled by giving the energy $E$ a small positive imaginary part $ie$. An integral of this form may be solved using the Cauchy principal-value prescription,
spin-1 polarization vectors we obtain Clebsch-Gordan coefficients in Eq. (14) in terms of the only need the formula for \( V \) and the quantities \( k \) where \( t \) is obtained from the integration over \( \eta \). 

\[ H_{0,L}(k) = \omega_{L,1}(k) + \omega_{L,2}(k) \]

and the quantities \( k_{0,L} \) are defined by

\[ H_{0,L}(k_{0,L}) = E. \]

In the decay \( \pi_1 \rightarrow \pi b_1 \), the D-wave amplitude is negligible compared to the S-wave, thus we may reduce our angular momentum space to \( L = 0, 1 \). Therefore, we only need the formula for \( V_{01}(p, q) \), and after expressing Clebsch-Gordan coefficients in Eq. (14) in terms of the spin-1 polarization vectors we obtain

\[ V_{01}(p, q) = -\frac{i\sqrt{3}}{4\pi} \sum_{\lambda, \lambda'} \int d\Omega_p d\Omega_q V(p, q, \lambda, \lambda') \times \epsilon^{jk} \epsilon^{i\lambda}(\lambda) \epsilon^{j\lambda'}(\lambda') q^k / q, \]

where \( V(p, q, \lambda, \lambda') \) was given in Eq. (24). The S-matrix is obtained from the t matrix via

\[ S_{LL'} = \delta_{LL'} - i \frac{k_{0,L} k_{0,L'}}{16\pi^2 E} t_{LL'}(k_{0,L}, k_{0,L'}). \]

As a 2 \times 2 unitary matrix it can be parametrized by two real scattering phase shifts \( \delta_0 \) and \( \delta_1 \), and one inelasticity parameter, \( \eta \), \( 0 \leq \eta \leq 1 \),

\[ S_{LL'} = \left( \begin{array}{cc} \eta^{2i\delta_0} & \eta \sqrt{1 - \eta^2} e^{i(\delta_0 + \delta_1)} \\ i \sqrt{1 - \eta^2} e^{i(\delta_0 + \delta_1)} & \eta e^{2i\delta_1} \end{array} \right). \]

Eq. (28) may be solved, as we already mentioned, by converting the integration over \( k \) into a sum over \( N \) integration points \( k_n, n = 1, 2, \ldots N \) (determined by Gaussian quadrature) with weights \( w_n \).

Let us assume that the coupling constant \( g_{FSI} \) is a variable, and we introduce the potential strength \( I = g_{FSI} / \rho_{FSI}^{(0)} \), where \( \rho_{FSI}^{(0)} = 50 \). In Figs. 7 and 8 we present the phase shifts as functions of \( I \), for the energy \( E = 1.6 \) GeV and the scale parameter \( \Lambda = 2 \) GeV. Their dependence on \( E \) for \( I = 1 \) and \( \Lambda = 2 \) GeV is shown in Fig. 9.

Within an experimental range of the expected hybrid mass \( m_{ex} \) is large, thus the Born approximation is accurate. Furthermore, smaller values of \( m_{ex} \) are closer to the \( \pi b_1 \) threshold, which (together with the second power of the FSI potential in momenta) should make the phase shifts small. Therefore, one expects that the final state interaction should not dramatically change the widths of \( \pi_1 \rightarrow \pi b_1 \) and \( \pi_1 \rightarrow \pi \rho \) obtained without final state interactions, such as was discussed in Ref. [15].

In Tables I and II we compare the original widths (in MeV) with the FSI-corrected ones for various values of \( m_{ex} \) and \( \Lambda \), for \( I = 2 \). The strength was doubled in order to compensate for the effect of the regulating factor in the FSI potential. For the values of \( \Lambda \) other than 2 GeV, the coupling constant \( g_{FSI} \) was renormalized such that

\[ m_{ex} \]
Phase shifts $\delta_0$ and $\delta_1$ in degrees, for the interaction $\pi b_1 - \pi \rho$ as functions of the energy $E$, for $I = 1$ and $\Lambda = 2$ GeV.

$q_0$ are the relative momenta between mesons for the $\pi b_1$ and $\pi \rho$ channels, respectively. In Figs. 10 and 11 we show their dependence on $I$ for $m_{\text{ex}} = 1.6$ GeV and $\Lambda = 2$ GeV. The behavior of $\delta_1$ indicates the existence of $\pi \rho$ resonances for the free parameters given.

From Figs. 10 and 11 it follows that the $\pi b_1$ channel can change significantly, if we increase the potential strength enough. However, even for large potentials the $\pi \rho$ channel width remains on the order of a few MeV. These results, which are shown in Tables I and II, are only weakly dependent on the value of the parameter $\Lambda$. Therefore, final state interactions cannot produce a large width for the $\pi_1 \rightarrow \pi \rho$ mode.

$$V(p, q) = -V_0 e^{-(p^2 + q^2)/\Lambda^2},$$

where $\Lambda$ is a free parameter on the order of 1 GeV. The strength $V_0$ was chosen such that the squares of the potentials in Eqs. (31) and (34) were equal, when integrated over $p$ and $q$. We see that the results obtained using this toy potential with a comparable strength are quite similar to those for our full model.

| $m_{\text{ex}}$ [GeV] | $\pi b_1$ | $\pi \rho$ |
|-----------------------|-----------|-----------|
| 1.4                   | 83(85)    | 9(9)      |
| 1.5                   | 145(149)  | 7(7)      |
| 1.6                   | 142(144)  | 5(5)      |
| 1.7                   | 118(117)  | 4(4)      |
| 1.8                   | 91(89)    | 2(2)      |
| 1.9                   | 66(64)    | 1(1)      |
| 2.0                   | 46(44)    | 1(1)      |

TABLE III: FSI-corrected widths in MeV for the $\pi b_1$ and $\pi \rho$ modes of the $\pi_1$ decay, for various values of $m_{\text{ex}}$ and for $\Lambda = 1$ GeV (2 GeV).
IV. SUMMARY

We have found that only unnaturally large potentials in the final state may change the widths for exotic meson decays calculated in microscopic models. In our model, the partial widths tend to change by only a few MeV. However, the $\pi\rho$ channel is subjected to a rather large relative correction. It may be caused by a large value of the original width for the $\pi b_1$ mode. The net width for the $\pi_1 \rightarrow \pi\rho$ channel is still of order of few MeV which agrees with the results of Refs. [13] or [14], but disagrees with results of Ref. [19].

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