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F MESON PRODUCTION IN e^+e^- ANNIHILATION

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August 12, 1980

ABSTRACT

The F meson production rate is calculated on the $\psi(4414)$ resonance using non-relativistic multipole expansion techniques. The results are compared with the available data.

The charm model which successfully describes the properties of the established ψ and D mesons, requires the existence of mesons having both charm and strangeness. The only evidence which supports the presence of these F-mesons in e^+e^- annihilation has been presented by DASP[1]. Data from Mark II places an upper limit for F-production which is close to the positive result of DASP[2]. It is now important to determine the production levels expected from theoretical considerations.

By limiting ourselves to a kinematic regime only slightly above $F\bar{F}$ threshold, we are able to utilize a non-relativistic multipole expansion to couple the gluonic degrees of freedom to the quark systems. As illustrated in fig. 1, the virtual photon from e^+e^- annihilation will produce a color singlet $c\bar{c}$ pair which ultimately couples to an effectively local gauge invariant gluon operator, which is the source of the strange quark content of the final state F-mesons. Although the intermediate gluons are highly virtual and timelike, due to the limited phase space available they have a small 3-momentum component. This long wavelength probe coupled to the small $c\bar{c}$ system naturally lends itself to a multipole expansion. The expansion parameter is of order $(\alpha_s \vec{r} \cdot \vec{p})$, where $\vec{r}$ is the radius of the intermediate $c\bar{c}$ system and $\vec{p}$ is the 3-momentum of a final state quark. As we will show, the relatively small value of this quantity affords good justification for our lowest order calculation.

To evaluate fig. 1, which is the 4-quark production amplitude which must later be projected onto the desired $F\bar{F}$ state, we will

first focus exclusively upon the charmed quark production and its coupling to the intermediate gluon modes. Only one time-ordering is relevant for this part of the process, and the amplitude is given by

$$\begin{aligned}
 A("Y" \rightarrow Q_c \bar{Q}_c + X) &= -ieQ_c(2\pi) \delta(E_f - E_i) \\
 &\times \langle X Q_c \bar{Q}_c | \int_0^\infty dt H_I(t) e^{-iHt} | Q_c' \bar{Q}_c' \rangle \\
 &\times \langle Q_c' \bar{Q}_c' | H_{EM}(0) | "Y" \rangle \frac{d^3 p_c'}{(2\pi)^3} \frac{d^3 \bar{p}_c'}{(2\pi)^3} \quad (1)
 \end{aligned}$$

where $H_I(t)$ is the operator that couples the charmed quarks to the intermediate gluon state, which is denoted by X . The evaluation of the first matrix element in equation 1 uses the now familiar techniques for coupling gluons of small 3-momentum to heavy non-relativistic quark boundstates [3]. To obtain the complete gauge invariant local gluon operator, it is necessary to sum all the diagrams of fig. 2. In order to get the explicit form and normalization of this operator, one must temporarily model the quark-antiquark interaction as being mediated by "color-coulomb" ladder exchange. Once the gauge invariant operator is obtained, it can be used to couple external gluons to a quark bound state whose internal binding mechanism deviates from a simple coulomb-like potential, as is the case for the charmonium system. The reduction of the matrix element yields

$$\begin{aligned}
 &\langle X Q_c \bar{Q}_c | \int_0^\infty dt H_I(t) e^{-iHt} | Q_c' \bar{Q}_c' \rangle \\
 &= (2\pi)^3 \delta^3(p_f - p_i) (igT_b)_{\alpha\bar{\alpha}} \delta_{\lambda_c \lambda_c'} \delta_{\bar{\lambda}_c \bar{\lambda}_c'} \\
 &\times \langle X Q_c \bar{Q}_c | r_i \frac{1}{\frac{p_c'^2}{m_c} + \epsilon' - \epsilon - (p_S^0 + \bar{p}_S^0)} | Q_c' \bar{Q}_c' \rangle \quad (2) \\
 &\times (\partial^i A_b^0 - \partial^0 A_b^i - gf_{bac} A_a^0 A_c^i)
 \end{aligned}$$

where $\epsilon'(\epsilon)$ is the interaction potential energy between the $c\bar{c}$ pair in the color singlet (octet) state, r_i is the relative separation of the $c\bar{c}$ pair, and $(p_S^0 + \bar{p}_S^0)$ is the energy of the emitted gluonic system. Note that this is simply the non-abelian analogue of the dipole $(\underline{r} \cdot \underline{E})$ transition familiar from electromagnetism.

To calculate the contribution from intermediate $c\bar{c}$ resonances we note that the energy denominator that appears in eq. 2 is exactly $(m_{res} - Q)$, where Q is the center-of-mass energy, and m_{res} is the mass of the relevant resonance. To complete the evaluation of eq. 2, we note that the expectation value of r_i for the $\ell = 0$ bound state to "free quark" transition can be estimated without knowing the exact boundstate wavefunction:

$$r_i \cong v_i^{rel} \times \left(\text{boundstate coherence time} \right) \cong \left(\frac{2p_c'}{m_c} \right) \times \frac{1}{E} \quad (3)$$

where $E = |m_{\text{res}} - 2m_c|$. This relationship holds exactly for a coulombic boundstate, and is an excellent approximation for a harmonic oscillator potential.

We can now trivially couple the gluon fields of eq. 2 to an $s\bar{s}$ pair as in fig. 1 and obtain the amplitude to lowest order in the 3-momentum:

$$A("Y" \rightarrow Q_c \bar{Q}_c Q_s \bar{Q}_s) = \frac{(ieQ_c)(igT_b)_{\alpha\bar{\alpha}}(igT_b)_{\beta\bar{\beta}}}{2m_s} (\chi_{\lambda_s}^+ \sigma_i \chi_{\lambda_{\bar{s}}}^-) (\chi_{\lambda_c}^+ \sigma_\mu \chi_{\lambda_{\bar{c}}}^-) \\ \times \frac{2p_c^i}{m_c E} \left(\frac{1 - \frac{3\alpha_s}{4}}{Q - m_{\text{res}} - i\Gamma_{\text{res}}/2} \right) (2\pi)^4 \delta^4(p_f - p_i) \quad (4)$$

where μ is the virtual photon's Lorentz index, and the χ 's are non-relativistic quark 2-spinors. The appearance of the factor $(1 - \frac{3\alpha_s}{4})$ is from the sum of the tree diagram of fig. 1a and the 1-loop diagram of fig. 1b. To obtain the $F\bar{F}$ final states from this 4-quark production amplitude, we must do color, spin, and momentum wavefunction projections. We define the center-of-mass momentum and internal momentum of the F-meson to be $\underline{p} \equiv \underline{p}_c + \underline{p}_s$ and $\underline{q} = \underline{p}_c - \underline{p}_s$, and of the $\bar{F}$ by $\bar{\underline{p}} = \bar{\underline{p}}_c + \bar{\underline{p}}_s$ and $\bar{\underline{q}} = \bar{\underline{p}}_c - \bar{\underline{p}}_s$. The momentum wavefunction projection becomes

$$A("Y" \rightarrow (Q_c \bar{Q}_s)(\bar{Q}_c Q_s)) = A("Y" \rightarrow Q_c \bar{Q}_c Q_s \bar{Q}_s) \times \frac{d^3 q}{(2\pi)^3} \phi_F(q) \times \frac{d^3 \bar{q}}{(2\pi)^3} \phi_{\bar{F}}(\bar{q}) \\ = A("Y" \rightarrow Q_c \bar{Q}_c Q_s \bar{Q}_s) \times \psi_F(0) \times \psi_{\bar{F}}(0) \quad (5)$$

where $\phi_F(\psi_F)$ is the momentum (spatial) wavefunction of the F-meson, and all terms linear in $\underline{q}$ and $\bar{\underline{q}}$ integrate to zero due to the $\ell = 0$ nature of the F. The color projection operator to be inserted in eq. 4 is simply $\delta_{\alpha\beta}\delta_{\bar{\alpha}\bar{\beta}}/3$. The spin projections are equally simple:

$$R_{i\mu} \equiv (\chi_{\lambda_s}^+ \sigma_i \chi_{\lambda_{\bar{s}}}^-) (\chi_{\lambda_c}^+ \sigma_\mu \chi_{\lambda_{\bar{c}}}^-) \quad (6a)$$

where for $F\bar{F}$ production

$$R_{i\mu}^{F\bar{F}} = \delta_{i\mu} \quad (6b)$$

for $F^*\bar{F}$ production

$$R_{i\mu}^{F^*\bar{F}} = -i \sqrt{2} \epsilon_{i\mu j} \epsilon_\lambda^j \quad (6c)$$

with ϵ_λ being the F^* polarization vector, and for $F^*\bar{F}^*$ production

$$R_{i\mu}^{F^*\bar{F}^*} = 2(\epsilon_i^\lambda \epsilon_\mu^\lambda - \epsilon_\lambda^\lambda \epsilon_{i\mu}^\lambda + \epsilon_i^\lambda \epsilon_\mu^\lambda) \quad (6d)$$

with ϵ_λ^λ and ϵ_λ^λ the polarization vectors for the F^* and $\bar{F}^*$ respectively. The projected amplitude becomes

$$A("Y" \rightarrow (Q_c \bar{Q}_s) + (\bar{Q}_c Q_s)) = \frac{-2ig^2 e Q_c \psi_F(0) \psi_{\bar{F}}(0) \cdot R_{iu} (1 - \frac{3\alpha_s}{4}) P^i (2\pi)^4 \delta^4(P_F - P_i)}{3m_c m_s E(Q - m_{res} - i\Gamma_{res}/2)} \quad (7)$$

Adding on the $(e^+ e^- \rightarrow "Y")$ part of the amplitude, squaring, and summing over final states yields

$$\sigma(F\bar{F}\text{-production}) = \frac{(4\pi\alpha_s)^2 (4\pi\alpha Q_c)^2 |\psi_F(0)|^2 |\psi_{\bar{F}}(0)|^2 (1 - \frac{3\alpha_s}{4})^2 m_F^{5/2} (Q - m_F - m_{\bar{F}})^{3/2} R}{27\pi m_c^2 m_s^2 E^2 ((Q - m_{res})^2 + \Gamma^2/4) Q^4} \quad (8)$$

where R is a statistical spin factor, and is equal to $(1, 4, 7)$ for $(F\bar{F}, F^*\bar{F} + F\bar{F}^*, F^*\bar{F}^*)$ production. An interesting feature of this calculation is that the quark binding interaction has not entered in any complicated way. We used a model coulomb potential to derive the form and normalization of the gauge invariant gluon operator, and after that the binding interaction only entered through the F -wavefunction at the origin, which can be estimated from an experimentally determined equation [4] using only the constituent quark masses.

Before numerically evaluating eq. 8, we will show to what extent this non-relativistic perturbative treatment is valid for F -production at the resonance of interest, the $\psi(4414)$. The multipole expansion parameter is of order $(\alpha_s r \cdot p_c)$. Using α_s defined on the scale of the resonance mass, and eq. 3 for r , this parameter

is of order 0.14 for the production of an $F\bar{F}$ pair at a center-of-mass energy of 4.4 GeV. Secondly, we must check the validity of the non-relativistic reduction of the 4-component quark spinors. The weakest possible link is the production vertex of the $s\bar{s}$ pair, and one finds in the reduction of $\bar{u}(p_s) \gamma_i v(\bar{p}_s) \rightarrow \chi_s^\dagger \sigma_i \chi_{\bar{s}}$, the neglected terms are of order 13% for $F\bar{F}$ production at a center-of-mass energy of 4.4 GeV. Lastly, we investigate possible perturbative final state interactions as in fig. 3, Figure 3a is already implicitly included in the use of the effective coupling constant of the lowest order diagram. The color trace of fig. 3b is down by $1/N_c$ from the lowest order diagram, indicating that the expansion parameter here is in fact α_s/N_c . (This is as expected from meson-meson scattering in the $1/N_c$ expansion [5].) So even for an $\alpha_s \sim .3 - .4$, neglect of these terms seems reasonable.

The production cross-section for $F\bar{F}$ states can now be evaluated as a function of energy using eq. 8. We use the DASP values $m_F = 2.03$ GeV and $m_{F^*} = 2.14$ GeV for the F -meson masses. The constituent charm and strange quark masses are 1.5 GeV and 0.5 GeV respectively. Using the quark masses, one estimates [4] $|\psi_F(0)|^2 = 1.1 \times 10^{-2} \text{ GeV}^3$, and α_s is defined at the mass scale of the bound state to which it is associated. Finally, we can now evaluate the $F\bar{F}$ production cross-section contribution from the narrow $(c\bar{c})$ -resonance at 4.414 GeV, whose measured width is 33 ± 10 MeV [6]. We find that for all possible F -meson spin states

$$\sigma_{F\bar{F}}(Q = 4.4 \text{ GeV}) = 1.4 \text{ nb}^{+1.4}_{-.6 \text{ nb}} \quad (9)$$

where the limits are from the experimental uncertainty in the resonance width. This contribution to $F\bar{F}$ production is very strongly peaked in energy and drops to less than 5% of its peak value as one tunes the energy more than 100 MeV off resonance. The spin state content of this cross-section is in the multiple ratio

$$F\bar{F}: F^*\bar{F} + F\bar{F}^*: F^*\bar{F}^* = 1: 2.2: 1.4 \quad (10)$$

The experimental cross-sections, after taking estimated branching ratios into account are, from DASP's inclusive η -production data[1]

$$\sigma_{F\bar{F}}(Q = 4.4 \text{ GeV}) \cong 5.3 \text{ nb}^{+1.3 \text{ nb}}_{-1.3 \text{ nb}}$$

and for Mark II, from $\eta\pi$ production the obtained upper limit[2]

$$\sigma_{F\bar{F}}(Q = 4.4 \text{ GeV}) < 3 \text{ nb}$$

Thus, our calculated F-meson production mechanism could account for all or most of the observed rate.

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FIGURE CAPTIONS

Figure 1: " γ " $\rightarrow Q_c \bar{Q}_c Q_s \bar{Q}_s$ diagrams. λ is the quark spin index, α and β are color indices.

Figure 2: Gauge invariant set of diagrams that gives the local gluon operator which couples to the $c\bar{c}$ resonance.

Figure 3: Possible final state corrections to calculated production mechanism.

—Figure 1—

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Figure 2

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Figure 3

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