DECOUPLING OF HIGHER VECTOR MESONS AND CHARMED MESON PRODUCTION IN e^+e^- ANNIHILATION

J.G. KÖRNER

II Institut für Theoretische Physik der Universitat, Hamburg, Germany

M. KURODA [1] and G. SCHIERHOLZ

Deutsches Elektronen-Synchrotron DESY, Hamburg, Germany

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The striking feature that $\bar{D}^*D^*$ production is greatly enhanced at energies just above threshold is explained. We argue that higher excited vector mesons successively decouple from lower spin final states which manifests itself in an explicit decoupling scheme. This can be verified on quite general grounds. A quantitative analysis of charmed meson production near threshold is given. Further evidence of the decoupling scheme coming from old data is discussed

The recently discovered charmed mesons [1] D and D* add further evidence to the simple picture that hadrons are bound states of quarks (and antiquarks) which come in (at least) four different flavours.

In the $c\bar{c}$ channel three $J^P = 1^-$ states, $\psi(3.1)$, $\psi'(3.7)$ and $\psi''(4.028)$, are established [e.g. 2], and three is further structure around 4.4 GeV. In the standard charmonium picture $\psi'(3.7)$ and $\psi''(4.028)$ are interpreted as radial excitations, and we expect many more of them as we go up in energy [e.g. 3]. However, they might be hard to detect in low multiplicity hadronic exclusive channels as we shall argue later on.

The analysis of charmed meson production in e^+e^- annihilation furthermore has shown the striking feature that $\bar{D}^*D^*$ production is greatly enhanced near threshhold. At $\sqrt{q^2} = 4.028$ GeV (i.e., at the peak of the resonance where the charmed meson yield is high) the measured cross sections $\sigma(\bar{D}^oD^o)$, $\sigma(\bar{D}^oD^{*o} + \bar{D}^{*o}D^o)$ and $\sigma(\bar{D}^{*o}D^{*o})$ are in the ratios [4] 1 : 22 . 20. If phase space is taken into account ‡1 the true enhancement factor is about 1 . 40 : 1400. This is to be compared with the naive quark model ratios [5] 1 . 4 . 7.

This fact has led several authors [6,7] to the conjecture that the peak at 4.028 GeV corresponds to a molecular $\bar{D}^*D^*$ state. We believe that this picture is very ill-founded. Such a state would strongly mix with the $c\bar{c}$ bound states (apart when in an exotic configuration). The actual calculation of the mass matrix [8], including loop corrections, show that $\bar{D}D$, $\bar{D}D^* + \bar{D}^*D$ and $\bar{D}^*D^*$ intermediate states have relatively little effect ‡2 on the mass of the $J^P = 1^-$ resonances. Moreover, $\bar{D}D$ and $\bar{D}^*D^*$ intermediate states contribute about equal so that the $\bar{D}^*D^*$ channel seems not to be distinguished from the $\bar{D}D$ channel. Note also that molecular $\bar{D}^*D^*$ states would predict a substantial ψ inclusive cross section which PLUTO [9] has failed to observe.

We shall argue that, more likely, a *decoupling scheme of excited states* (higher vector mesons in this case) *from low lying resonances* lies at the root of the puzzle. Such scheme can be verified in various models as we shall see. This is also a feature of a Bethe–Salpeter model of confined quarks [10]. The decoupling scheme could already have been discussed in connection with old data, e g., $\rho''(1600) \not\rightarrow 2\pi$. But there the situation is not so spectacular.

The decoupling scheme says ‡3, e.g., that only the first $J_1 + J_2 + 1$ vector particles (ground state plus $J_1 + J_2$ excitations/daughters) may couple to a pair of leading spin-J_1 and spin-J_2 particles. This is understood

[1] Alexander von Humboldt Foundation Fellow.

‡1 The available decay energies are 295 MeV, 155 MeV and 16 MeV respectively.

‡2 Though the effect increases with increasing energy

‡3 The supposition on which the decoupling rules hold will be stated later on

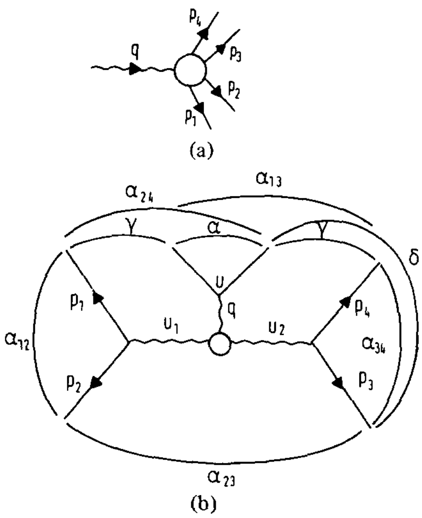

Fig 1. (a) Single current amplitude with four hadrons. (b) Choice of variables for 6-point Veneziano amplitude.

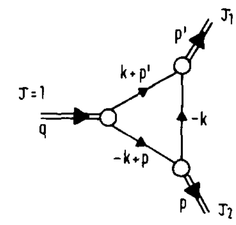

Fig. 2. Quark loop diagram

to hold for all different Regge trajectories, e.g., ρ, ω, ϕ, ψ, etc. separately. This predicts, e.g., $\rho'(1250) \not\to 2\pi$, $\rho''(1600) \not\to 2\pi$ and $\rho''(1600) \not\to \pi\omega$ which seems to be in excellent agreement with experiment [e.g. 11] ‡4. As a further consequence $\psi''(4.028)$ may not couple to $\bar{D}D$ and $\bar{D}D^* + \bar{D}^*D$ which we will discuss in connection with the charm production data later on. The decoupling scheme can in principle be extended to arbitrary spin particles (i.e., beyond vector mesons). This, however, will be hard to test for the time being.

There are further indications for the validity of the decoupling scheme, though indirect. The pion form factor shows a monopole behaviour and can be well described by the ρ pole [e.g. 13]. The nucleon form factor, on the other hand, behaves like a dipole and its isovector part seems to be well accounted for [14] by the ρ and ρ'. This is exactly what we expect from the decoupling scheme. As a corollary of this the universality of the ρ coupling constant should be broken which indeed can be inferred from ref. [14] for the ρNN coupling (using $g^2_{\rho\pi\pi}/4\pi = 2.84 \pm 0.06$). Further evidence comes from the $\omega \to \pi\gamma$ and $\pi^0 \to \gamma\gamma$ decays. The $\omega \to \pi\gamma$ decay can be well described by ρ and ρ' dominance according to our rule ‡5 [16]. The $\pi^0 \to \gamma\gamma$ decay, on the other hand, seems to require an infinity of vector mesons [16] which is predicted by the decoupling scheme ‡6.

We shall now give several heuristic derivations of the decoupling rules. Whenever necessary we assume that hadrons are made out of spin-zero quarks. We believe, however, that this assumption is not crucial. We shall also restrict ourselves to bosons. Apart from the nucleon and, perhaps, the Δ very little can be said about baryons.

(i) We find that the most decreasing electromagnetic form factor coupling a spin-J_1 boson to a spin-J_2 boson behaves asymptotically like (within logarithms) $(q^2)^{-J_1-J_2-1}$. This follows from a slight extension of ref. [17] to arbitrary spin particles and canonical form factor behaviour [18], i.e., ϕ^4 for scalar constituents corresponding to γ_μ coupling for spin-1/2 constituents. *We believe this behaviour to be symptomatic of asymptotic freedom* ‡7. For this form factor let us write down the dispersion relation

$$F(q^2) = \frac{1}{\pi} \int_{q_0^2}^{\infty} \mathrm{d}q'^2 \, \frac{\mathrm{Im}\, F(q'^2)}{q'^2 - q^2}, \tag{1}$$

which, considering only a single two-body or quasi two-body intermediate state for the moment, has the solution ‡8 (i.e., the inverse D matrix)

‡4 The dominant decay mode of the $\rho''(1600)$ seems to be $\rho(\pi\pi)_{I=0}$ Note that the $(\pi\pi)_{I=0}$ system cannot be in a $J = 0$ resonant state [12] so that there is no contradiction to the decoupling scheme

‡5 Single ρ dominance does not agree with experiment [15].

‡6 For the coupling of two excited vector mesons see later on

‡7 For ϕ^3, which differs from the canonical form factor behaviour by one power of q^2, see also ref. [19]

‡8 Our attitude is to adopt the maximal analyticity principle and so we assume that there are no CDD poles. For leading spin-J_1 and spin-J_2 particles this is the standard assumption and seems not to contradict any sacred principle. For degenerate daughter states (e.g., $\gamma_V \to \rho'\pi$) this might, however, not be tenable in general by arguments of selfconsistency of the decoupling scheme.

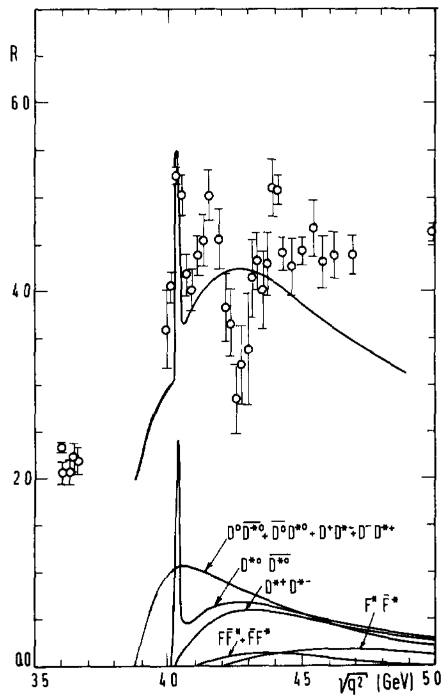

Fig. 3. Fit to the total cross section. The non-charmed contribution is assumed to be $R = 2$. Also shown are the single channel contributions. The $\bar{D}D$ ($F\bar{F}$) cross section contributes about 0.05 (0.004) to R around $\sqrt{q^2} = 4.3$ GeV. Not included is the heavy lepton which contributes about 0.6 at $\sqrt{q^2} = 4.3$ GeV. Data are taken from ref. [31].

$$F(q^2) = \exp\left\{\frac{1}{\pi}\int_{q_0^2}^{\infty} dq'^2 \frac{\delta(q'^2)}{q'^2 - q^2}\right\}, \quad \delta(q_0^2) = 0\,, \quad (2)$$

except for a normalization constant. In principle eq. (2) could be multiplied by a polynomial [‡9] in q^2. All dispersion calculations are afflicted with similar ambiguities. Such polynomial can, however, be excluded for reasons of analyticity in our case. Analyticity of F in J_1 and J_2 says that the polynomial must be of the same order for all J_1, J_2 which means that it can only be of order zero given the asymptotic behaviour of F. Levinson's theorem then tells us that the asymptotic behaviour of eq. (2) is

$$F(q^2) \approx (q^2)^{-N}\,, \quad (3)$$

where N is the number of resonances that the phase shift δ is passing through, i.e., the number of resonances coupling to F. this can be generalized as to include bound states (again assuming that there are no CDD poles) which gives the same result but N being now the number of bound states plus resonances coupling to F. If we combine this with the asymptotic form factor behaviour obtained from an underlying quark structure (asymptotic freedom), we find $N = J_1 + J_2 + 1$. This result can be extended to the case of a coupled channel dispersion relation [20] allowing for absorption. Our result $N = J_1 + J_2 + 1$ is exactly the decoupling scheme. In principle we could employ the same method for the baryons if the asymptotic behaviour of the form factors would be known. This, however, involves a good deal of model dependence due to the three-quark nature of the baryons. The trouble starts already with the nucleon. Brodsky and Farrar [21] obtain $F_2 \sim (q^2)^{-3}$. Other models [‡10] give $F_2 \sim (q^2)^{-2}$. It seems to be certain that two resonances couple to the nucleon (e.g., the ρ and ρ'). But it cannot be excluded that even a third (e.g., the ρ'') will couple [24]. Note that this derivation, in general, does not allow any conclusions on the nature of these resonances, e.g., if they are the lowest lying excitations, etc.

(ii) The dual model for currents constitutes our next example. Following Ademollo and Del Giudice [25] the amplitude for one vector current and four scalar particles (fig. 1a) reads

$$\begin{aligned} T_\mu = &\int_0^1 du\, u^{-\delta-1}(1-u)^{-\alpha-1} \int_0^1\!\!\int_0^1 du_1 du_2 u_1^{-\alpha_{12}-1} \\ &\times (1-u_1)^{-\gamma-1} u_2^{-\alpha_{34}-1}(1-u_2)^{-\gamma-1}(1-uv_1)^{-\alpha_{24}+\alpha-\gamma} \\ &\times (1-uu_2)^{-\alpha_{13}+\alpha+\gamma}(1-uu_1u_2)^{-\alpha_{23}+\alpha_{13}+\alpha_{24}-\alpha} \\ &\times \Big[(1-u)(1-u_1)u_2(p_2+p_3+p_4-p_1)_\mu \\ &+ \frac{(1-u)(1-uu_2)(1-u_2)}{u(1+u_2-u_1u_2-uu_2)}(p_3+p_4-p_1-p_2)_\mu\Big]. \end{aligned} \quad (4)$$

The notation is explained in fig. 1b. The matrix element of the vector current coupling to a spin-J_1 and spin-J_2 particle is obtained by factorizing eq. (4) at the poles $\alpha_{12} = J_1$, $\alpha_{34} = J_2$. Note that the first factor in the squared bracket in eq. (4) vanishes at these poles.

‡9 We distinguish between "dynamical" and "kinematical" ambiguities.

‡10 E.g., the gauge invariant vector model [22] or if the spectator quarks form a mesonic bound system [e.g. 23].

For the most decreasing form factor we find [26]

$$F(q^2) = \int_0^1 \mathrm{d}u\, u^{-\alpha(q^2)}(1-u)^{-\delta-2+J_1+J_2}$$

$$= B_4(-1-\delta+J_1+J_2,\, 1-\alpha(q^2)) . \qquad (5)$$

The constant δ has the meaning of a trajectory in an exotic channel (fig. 1b) so that $\delta = \lim_{t\to-\infty}\alpha(t)$. Dimensional counting [18,21] then gives[‡11] $\delta = -2$ for which eq. (5) becomes asymptotically

$$F(q^2) \sim (q^2)^{-J_1+J_2-1} , \qquad (6)$$

i.e., the familiar result[‡12]. A simple exercise now shows that F (eq. (5)) has exactly J_1+J_2+1 poles corresponding to the leading vector meson and the first J_1+J_2 spin-1 daughters in the photon channel, whereas the (J_1+J_2+1)st, (J_1+J_2+2)nd, etc. daughters decouple. This verifies our decoupling rules once again. The mechanism which makes the higher daughters decouple is the same as that causing the Odorico zeroes [27] in the strong Veneziano amplitudes. In the strong Veneziano amplitudes low multiplicity states must also decouple from high mass resonances as has been remarked by Gliozzi [28]. There, the decoupling scheme has, however, not been developed that far. If one or both of the spin-J_1, spin-J_2 particles are nonleading, i.e. daughter states themselves, the general analysis [26] tells us that daughters count like their leading ancestors in the decoupling scheme.

(iii) Our final example rests on a class of bound state models of confined quarks. Here the vector meson to spin-J_1 and spin-J_2 coupling proceeds via the quark loop diagram as shown in fig. 2. We will base our discussion on the model presented in ref. [10], but believe it to be more general. The model has been solved in Euclidian space where it exhibits O(4) symmetry as a matter of quark confinement. The mesonic bound states lie on linear Regge trajectories which follow the familiar O(4) classification. At Euclidian (external) momenta fig. 2 can be expressed by an integral over Euclidian internal momenta. This leaves us with an integral over O(4) spherical harmonics and reduced wave functions. We define the coupling constant at the (symmetric) point where all external momenta are zero. It is clear that this will lead to decoupling rules due to O(4) Clebsch–Gordan coefficients entering in fig. 2. In order to calculate the coupling constant (in general there are many) from the helicity amplitudes we have to factor out the invariants [e.g. 19] before going to the limit of zero momenta since many of them vanish in this limit. To give an example let us consider $J_1 = J_2$. The (leading) spin-J_1 particles contribute a spherical harmonic $\mathcal{Y}_{J_1 J_1 \lambda_1}$ each. The vector meson contributes[‡13] $\mathcal{Y}_{n1\lambda}$ where $n = 1, 2, \ldots$ Here $n = 1$ denotes the parent particle, $n = 2$ the first daughter, etc. After factorization[‡14], the vertex with the highest number of daughters possible is proportional to the O(4) Clebsch–Gordan coefficient

$$\left(\begin{matrix} \frac{J_1}{2}\,\frac{J_1}{2} & \frac{J_1}{2}\,\frac{J_1}{2} \\ & , \\ J_1\,\lambda_1 & J_2\,\lambda_2 \end{matrix}\left|\begin{matrix} \frac{n-1}{2}\,\frac{n-1}{2} \\ \\ 0 \quad 0 \end{matrix}\right.\right) , \qquad (7)$$

which leads to the selection rule

$$n = 1, 2, \ldots, \quad 2J_1+1 , \qquad (8)$$

in accordance with the decoupling scheme. This can be generalized to arbitrary J_1, J_2 giving back the general decoupling rules. We have not made any attempt to calculate the coupling constant away from zero momenta. It should, however, not depend on the external momenta since its only dependence could be an entire function in the external momenta due to quark confinement which is not tolerable.

From what has been said above it is clear that the decoupling scheme holds *separately* for any Regge trajectory (e.g., any type of quark) as there are ρ, ω,

‡11 For a different argument see also ref. [25].

‡12 This is not surprising since $\delta = -2$ reflects the same dynamics underlying our earlier results. We could have also left δ a free parameter to be fitted to the canonical pion form factor behaviour which, however, would lead us to the same conclusions.

‡13 Note that the odd daughters have charge conjugation parity plus in contrast to the Veneziano model. This shortcoming is due to scalar quarks and can be cured in the case of spin-1/2 quarks without changing our results. Here, we may have the wave functions $(\gamma_4)^n \mathcal{Y}_{n1\lambda}$, $(\gamma_4)^n\, \gamma\, S_3 \mathcal{Y}_{n-100}$, etc. as can be inferred from ref. [10] extended to spin-1/2 quarks. Note also that the interpretation of ψ', ψ'', etc. is quite different from the charmonium picture here.

‡14 Here, it is sufficient to factor out $(p-p')_\mu$ from $\mathcal{Y}_{n1\lambda}$ (i.e., the relative momentum $k_\mu - \frac{1}{2}(p-p')_\mu$ and keeping $(p-p')_\mu$, for the kinematics see fig. 2). The rest of the helicity amplitude stays finite in the limit of zero momenta

ϕ, ψ, etc. This is to say that, e.g., ψ, ψ' and ψ'' as well as ρ, ρ' and ρ'' may couple tc $\bar{D}^*D^*$. As for approach (i) this means that the D matrix becomes a sum over various Regge trajectories. From approach (iii) we expect further trajectories associated with hyperradially excited states [10] starting at 5 – 6 GeV.

If we turn to the more realistic case of spin-1/2 quarks we notice the possibility of a further degeneracy according to a (dynamical) decoupling of the ${}^3S_1 - {}^3D_1$ system. This seems to be given in the charmonium picture and be realized in nature. The recently reported resonance [‡15] at 3.77 GeV may be understood to be the ${}^3D_1$ $c\bar{c}$ state in agreement with charmonium estimates. Hence, we expect a doubling of the vector meson spectrum, i.e., two Regge trajectories for each family of quarks, at least for the new quarks. In practice this means that on top of ψ, ψ' and ψ'' we also expect $({}^3D_1)$, $({}^3D_1)'$ and $({}^3D_1)''$ $c\bar{c}$ states coupling to $\bar{D}^*D^*$ and similarly for any other final state.

Let us now come back to the charmed meson production. The cross section for $\bar{D}D$, $\bar{D}D^* + \bar{D}^*D$ and $\bar{D}^*D^*$ production is given by ($\eta = q^2/(m+m')^2$)

$$\frac{d\sigma}{d\cos\theta} = \frac{8\pi}{3}\,\frac{\alpha^2 p_{cm}^3}{(q^2)^{5/2}}\,\frac{3}{4}\sin^2\theta\,|G_E|^2\,, \tag{9}$$

$$\frac{d\sigma}{d\cos\theta} = \frac{8\pi}{3}\,\frac{\alpha^2 p_{cm}^3}{(q^2)^{5/2}}\,\frac{3}{8}(1+\cos^2\theta)\,4\eta\,|G_M^*|^2\,, \tag{10}$$

and

$$\frac{d\sigma}{d\cos\theta} = \frac{8\pi}{3}\,\frac{\alpha^2 p_{cm}^3}{(q^2)^{5/2}}\,[\tfrac{3}{8}(1+\cos^2\theta)\,4\eta\,|G_M|^2 + \tfrac{3}{4}\sin^2\theta\,(3|G_E|^2 + \tfrac{8}{3}\eta^2|G_Q|^2)]\,, \tag{11}$$

respectively [‡16]. In terms of invariant form factors the multipole form factors read

$$G_E = F\,, \tag{12}$$

$$G_M^* = \tfrac{1}{2}(m+m')F^*\,, \tag{13}$$

and

$$G_E = F_1 + \tfrac{2}{3}\eta[F_2 + 2(1-\eta)F_3]\,,$$
$$G_M = F_1 + F_2\,, \tag{14}$$
$$G_Q = -F_2 - 2(1-\eta)F_3\,,$$

respectively (in terms of the form factors of ref. [29]) we have $F_1 = F^{(1)}$, $F_2 = -F^{(1)} + F^{(2)}$ and $F_3 = F^{(3)}$).

We now expect one vector meson coupling to F and two vector mesons coupling to F^*. Furthermore, we expect [‡17] three vector mesons coupling to F_3 and two vector mesons coupling to F_1 and F_2 each. The form factors are written as a product of poles $\Pi_i(1 - q^2/m_i^2)$ (and Breit–Wigners for ψ'') as argued in the first part of this paper. The contribution of ρ, ω, ϕ and ψ type vector mesons and their recurrences are considered separately. For the ρ, ω and ϕ recurrences we assume $(\alpha')^{-1} = 1\ \mathrm{GeV}^2$. The D and D* masses are taken to be $m_{D^0} = 1.867$ GeV, $m_{D^+} = 1.871$ GeV and $m_{D^{*0}} = 2.006$ GeV, $m_{D^{*+}} = 2.013$ GeV [4].

In case of $\bar{D}D^* + \bar{D}^*D$ production we need to specify $G_M^*(0)$ and the relative contribution of ρ, ω, ϕ and ψ type vector mesons. For the latter we call upon the charmonium picture which tells us that the ϕ and ψ contributions have to be weighted by factors m_u/m_s and m_u/m_c respectively [‡18]. The normalization is then given by the $\omega \to \pi\gamma$ decay rate ($g_{\omega\pi\gamma} = 2.56\ \mathrm{GeV}^{-1}$) which leads to $\Gamma_{D^{*+}\to D^+\gamma} = 1.64$ keV and $\Gamma_{D^{*0}\to D^0\gamma} = 24.7$ keV.

In case of $\bar{D}^*D^*$ production we totally rely on the quark model since the magnetic moments of even the old vector mesons are unknown. The naive quark model predicts for the $c\bar{c}$ sector $G_M^{\psi}(0) = 2m_{D^*}/3m_c$ and relative weighting by the quark masses as before. Little is known about the electric quadrupole moment except that it should be small since it arises from the d-wave admixture in the D* wave function. The relative weight, however, can be calculated and turns out to be given by the square of the quark mass ratios [30]. We shall normalize $F_3^{\psi}(0)$ to the peak of the DORIS total cross section data [31] corresponding to $m_{\psi''} = 4.034$ GeV. This gives $F_3^{\psi}(0) = -\frac{2}{3}\cdot 0.78$ and a comparatively small ratio $G_Q(0)/G_M(0) \approx 0.24$. For the ψ'' (into $\bar{D}^{*0}D^{*0}$) width we obtain $\Gamma_{\psi''} = 7.5$ MeV which finally leads to $\Gamma_{\psi''}^{e^+e^-} \approx 0.1$ keV.

‡15 G Feldman, L. Galtieri and M. Perl, SLAC telegram (May 26, 1977)

‡16 Using degenerate masses eqs (9)–(11) can be seen to lead to the quark model cross section ratios 1 . 4 7 close to threshold ($\eta \approx 1$) if the quark results $G_E = G_M = G_{M^*}$ and $G_Q = 0$ are used.

‡17 See, e.g., ref. [26] generalized to vector currents (essentially eq. (4)) and ref. [19] scaled to the canonical form factor behaviour.

‡18 We have assumed $m_u = 0\ 32$ GeV and $m_c = 1\ 6$ GeV.

The fit to the total cross section is shown in fig. 3. Due to the p-wave threshold factor entering in the Breit–Wigner width, the resonance curve is quite distorted. Note that the resonance would have been broadened towards higher energies even further had we included $D^{*+}D^{*-}$ final states. Also shown are the contributions of the various final states separately (the cross section for $\bar{D}D$ production is too small to be included in the figure). On the peak the cross sections $\sigma(\bar{D}^0D^0)$, $\sigma(\bar{D}^0D^{*0} + \bar{D}^{*0}D^0)$ and $\sigma(\bar{D}^{*0}D^{*0})$ are found in the ratios [‡19] 1 · 46 . 201. The relative amount of $\bar{D}^*D^*$ changes quite rapidly in the peak area and only stabilizes outside the resonance. If the theoretically predicted cross section is binned in intervals of, say, 20 MeV we obtain about the SPEAR value for the ratio $\sigma(\bar{D}^0D^{*0} + \bar{D}^{*0}D^0) . \sigma(\bar{D}^{*0}D^{*0})$ in the region covering the peak. On the peak the $\bar{D}^{*0}D^{*0}$ cross section arises almost entirely from ψ'' coupling to F_3 which predicts a purely longitudinal (i.e., $1 - \cos^2\theta$) distribution in this region. We have not attempted a precision fit to the total cross section allowing for further energy dependence of the width (as, e.g., suggested by the effective range approximation). So our width is strongly limited from above by the shape of the total cross section.

In fig. 3 we have also drawn our predictions for $F^+F^{*-} + F^{*+}F^-$ and $F^{*+}F^{*-}$ production using $m_F = 1.975$ GeV and $m_{F^*} = 2.061$ GeV [32] (again, the F^+F^- cross section is too small to be included in the figure). The cross sections are small as compared to D and D* production which might explain why the F and F* have not yet been seen.

Our decoupling rules are of great impact for future e^+e^- phenomenology. We expect the ψ''' (mass 4.4 GeV?) *not* to decay into a single $\bar{D}^*D^*$ pair and any lower spin configuration, i.e., $\bar{D}D$, $\bar{D}D^* + \bar{D}^*D$, etc. But the $(^3D_1)'$ and $(^3D_1)''$ (if established) may couple to $\bar{D}^*D^*$ and the $(^3D_1)'$ also to $\bar{D}D^* + \bar{D}^*D$. This might help to sort out the structure of the total cross section in the 4–4.5 GeV region. At higher energies we expect $\sigma(\bar{D}^*D^*)$ (and σ of any lower spin configuration) to be smooth. This should manifest itself also in the K^0 and μ/e inclusive cross section. Remember, however, that there is the possibility that trajectories of hyperradially excited states show up at higher energies.

‡19 The ratio $\sigma(\bar{D}^0D^{*0} + \bar{D}^{*0}D^0) . \sigma(\bar{D}^{*0}D^{*0})$ is quite dependent on the input assumption of the fit. If one, e.g., takes the naive quark model value $F^{*\psi}(0) = 2/3m_c$ we obtain $\sigma(\bar{D}^0D^{*0} + \bar{D}^{*0}D^0) : \sigma(\bar{D}^{*0}D^{*0}) = 1 \cdot 3\ 5$ rather than 1 . 4.4.

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