INCLUSIVE ELECTRON SPECTRUM IN e^+e^- ANNIHILATION NEAR CHARM THRESHOLD

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The electron energy spectrum from charm meson decays is calculated and compared with the recent experimental data at c.m.s. energy 4.0–4.2 GeV.

Single electrons produced with hadrons in e^+e^- annihilation have been observed at DORIS at cms energies of 4.0–4.2 GeV [1]. The electron yield at lower energies (below ψ') was found to be consistent with the background, indicating a threshold for such events somewhere between 3.7 and 4.0 GeV. We note the following aspects of the data:
(1) The hadrons associated with the electrons have an average multiplicity $\geqslant 3$.
(2) The electrons have an energy spectrum concentrated at low energies.

That the events originate from the production and decay of a heavy lepton pair [2] is improbable, since one then expects the average hadron multiplicity to be $\leqslant 3$. The momentum spectrum of the electron further supports this view; as it is inconsistent with that from the heavy lepton decay, unless the associated neutrino has a mass of about 1 GeV or higher [1].

The charm model [4] predicts charm mesons of about 2 GeV which, when produced in e^+e^- annihilation, would also produce the inclusive electron signals. Recent discovery at SLAC of resonances decaying into non-zero strangeness final states of masses ~1.87 GeV and ~2.01 GeV [5] strengthens the charm interpretation of the observed narrow resonances $J(\psi)$, ψ' etc. Because of the available decay modes of both nonleptonic and semi-leptonic decay modes, the hadron multiplicity associated with the electron in this case is expected to be $\geqslant 3$, and for the same reason the electron momentum is expected to be concentrated at low energies. The gross features of the data seem to be consistent with the production and subsequent decays of charm mesons. It is, however, important to explain in detail the electron spectrum as well as various other electron-hadron correlations in the charm theory. The electron energy spectrum from charm meson decays can be straightforwardly calculated using well-known current algebra methods, generalized to SU(4) symmetry. This note reports the result of such a calculation.

From either theoretical [6] or experimental [5] arguments, the production cross section for $D\bar{D}$ is suppressed compared with that of DD^*. Identifying the observed resonances at SLAC with D (m_D = 1.87 GeV) and D^* (2.01 GeV), we conclude they are produced with rather low momentum at cms energies about 4.0–4.2 GeV. Since D^* decays almost immediately to $D\pi$ or $D\gamma$, we therefore concentrate on the electron spectrum from D meson decays. Assuming $\Delta c = \Delta s$ selection rule for charm changing weak currents, the D meson can decay via

$$D \to K e \bar{\nu}_e, \qquad D \to K\pi e \bar{\nu}_e, \qquad K\pi\pi e \bar{\nu}_e, \text{ etc.} \tag{1}$$

Multi-pion decay modes are not only suppressed by the available phase space but also by the low energy theorem which says that the semi-leptonic decay rate vanishes if any one of the pions is soft [7].

In applying the current algebra to known meson and hyperon decays, one assumes certain pole dominance. We assume here that the $K\pi$ and $K\pi\pi$ channel are dominated by K^* and K_A poles, respectively. The form factors for the currents coupling D to K and K* can be written as:

$$\langle \mathrm{K}(k)|V_\nu|\mathrm{D}(p)\rangle = f_+(q^2)(p+k)_\nu + f_-(q^2)q_\nu$$
$$\langle \mathrm{K}^*(k)|(V_\nu+\lambda A_\nu)|\mathrm{D}(p)\rangle = \mathrm{i}F_1^{\mathrm{V}}(q^2)\,\epsilon_{\mu\nu\lambda\sigma}\,\epsilon^\mu k^\lambda q^\sigma + \lambda\epsilon^\mu[F_1^{\mathrm{A}}(q^2)g_{\mu\nu} + F_2^{\mathrm{A}}(q^2)p_\mu k_\nu + F_3^{\mathrm{A}}p_\mu q_\nu] \quad (2)$$

where $q = p - k$, $\lambda = \mp 1$, for V $\mp$ A currents.

(In calculating the electron spectrum, we can neglect the f_- and F_3^{A} terms whose contribution are proportional to $\mathrm{O}(m_{\mathrm{e}})$). We calculate the above form factors by assuming that the currents are dominated by the corresponding vector and axial-vector meson poles. The three point vertices, so introduced, are determined by employing the hard meson technique of Schnitzer and Weinberg [8]. The essential feature of this technique is to expand the vertices in simple polynomials in momenta and determine the coefficients by subjecting them to Ward identities. Somewhat arbitrarily, but in conformity with the usual practice, we assume that the scalar meson contribution is small and will not change our result dramatically. Since the method is well-described in the literature and the application to SU(4) symmetry is straightforward, we give here only the results:

$$f_+(q^2) = \frac{m_{\mathrm{F}^*}^2}{(m_{\mathrm{F}^*}^2 - q^2)}\left[1 - \left(\frac{1+\delta}{4}\right)\left(1 + \frac{m_{\mathrm{K_A}}^2}{m_{\mathrm{D_A}}^2}\right)\frac{q^2}{m_{\mathrm{K_A}}^2}\right]$$

$$F_1^{\mathrm{A}}(q^2) = \sqrt{2}m_{\mathrm{K}^*} + m_{\mathrm{K}^*}\left(1 + \frac{m_{\mathrm{F_A}}^2}{m_{\mathrm{D_A}}^2}\right)\Big/2\sqrt{2}(q^2 - m_{\mathrm{F_A}}^2)\,[(m_{\mathrm{K}^*}^2 - q^2) + \tfrac{1}{2}\delta(m_{\mathrm{D}}^2 + m_{\mathrm{K}^*}^2 - q^2)] \quad (3)$$

$$F_2^{\mathrm{A}}(q^2) = -\frac{\delta m_{\mathrm{K}^*}}{2\sqrt{2}(q^2 - m_{\mathrm{F_A}}^2)}\left(1 + \frac{m_{\mathrm{F_A}}^2}{m_{\mathrm{D_A}}^2}\right).$$

Here $\mathrm{K_A}$, $\mathrm{F_A}$, $\mathrm{D_A}$ are the axial partners of K*, F* and D*, respectively. δ is a parameter not determined by this method (from the analysis of $\mathrm{A}_{1\rho\pi}$ system, one finds $\delta = -0.5$ is consistent with the width of A_1, but because of the uncertainty in the data and in the value of the pion decay constant F_π, one expects that δ lies between 0 and -1) [8]. In deriving the above results, we have incorporated the SU(4) breaking through the use of First Weinberg sum rule [9]. We assume that $F_{\mathrm{D}} = F_{\mathrm{F}} = F_{\mathrm{K}} = F_\pi$. In simplifying the results, we have used the KSRF relation [10], namely

$$g_\rho^2/m_\rho^2 = 2F_\pi^2 .$$

The vector form factor $F_1^{\mathrm{V}}(q^2)$ cannot be determined by the current algebra method. In pole-dominance form, one has

$$F_1^{\mathrm{V}}(q^2) = \sqrt{2}g_{\mathrm{F}^*}\,g_{\mathrm{F^*K^*D}}/(q^2 - m_{\mathrm{F}^*}^2) \quad (4)$$

We give three estimates for the coupling constant F_1^{V}:

(1) Using SU(4) Clebsch–Gordon coefficient one can relate $g_{\mathrm{F^*K^*D}}$ to $g_{\omega\rho\pi}$ thus obtaining

$$g_{\mathrm{F}^*}g_{\mathrm{F^*K^*D}} = \left(\frac{m_{\mathrm{F}^*}}{m_\rho}g_\rho\right)\frac{g_{\omega\rho\pi}}{\sqrt{2}} = \frac{m_{\mathrm{F}^*}m_\rho}{\sqrt{2}}\left[\frac{3\Gamma(\omega\to\pi\gamma)}{\alpha k^3}\right]^{1/2}$$

(2) Incorporating the SU(4) symmetry breaking in the coupling constants by assuming the Gell-Mann-Okubo ansatz. From the known vector meson decays one then finds a 30% decrease in $g_{\mathrm{F^*K^*D}}$ compared to the value from (4) [11].

(3) Noting that $g_{\mathrm{F^*K^*D}}$ has inverse mass dimension, if the mass scale is typical of charm mesons rather than of a universal mass, the result (4) above is suppressed by a factor m_ρ/m_{F^*}.

Similar methods can be applied to $\mathrm{D}\to\mathrm{K}\pi\pi e\nu_{\mathrm{e}}$. Assuming $\mathrm{K_A}$ dominance, we find the decay rate is less than

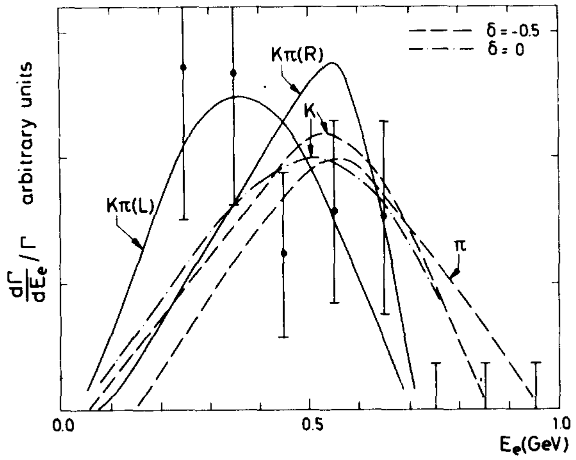

Fig. 1. The energy spectrum for various channels of D decay at rest, i.e. $D \to Ke\bar{\nu}$ ($D \to \pi e\bar{\nu}$) and $D \to K\pi e\bar{\nu}$ for two values of δ. For $D \to K\pi e\bar{\nu}$, the difference between the two cases $\delta = 0$ and $\delta = -0.5$ is negligible.

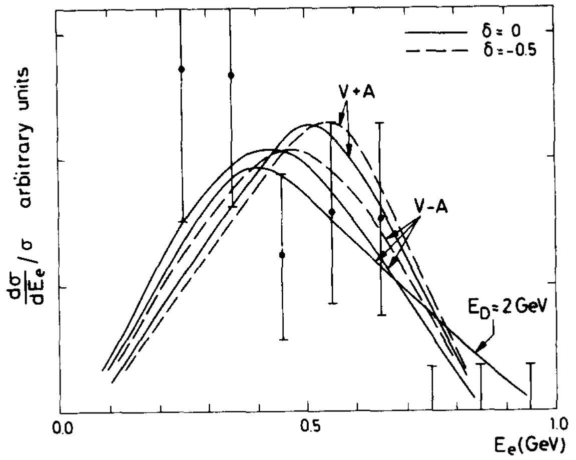

Fig. 2. The energy spectrum of D decay at rest after summing over the $Ke\bar{\nu}$ and $K\pi e\bar{\nu}$ for two values of δ. We also show the spectrum for D decay with $E_D = 2$ GeV.

Table 1
Ratio of the decay rates $R = \Gamma(D \to Ke\nu_e)/\Gamma(D \to K\pi e\nu_e)$ for different values of δ.

δ	0	-0.5	-1.0
R	1.6	2.2	3.1

5% as compared to the rate of $D \to K\pi e\nu_e$ and thus we neglect it in the following discussion.

The electron spectra for $D \to Ke\nu_e$ and $D \to K\pi e\nu_e$ via K^* pole-dominance are calculated numerically. Our results are presented in figs. 1 and 2, together with the available data (we assume $m_{F_A}^2/m_{D_A}^2 \approx m_{F^*}^2/m_{D^*}^2 \approx 1$ for numerical calculation); however the small number of events observed makes the comparison only qualitatively meaningful. Fig. 1 presents the spectrum for individual decay channels of D at rest. Fig. 2 presents the combination of $D \to Ke\nu_e$ and $D \to K\pi e\nu_e$ with the relative normalization determined by eq. (3). (see table 1.) We have assumed (4) in the curves for $D \to K\pi e\nu_e$ to allow the maximum effect of VA interference (see below). We note that the shape of the energy spectrum in each case does not change dramatically for different values of δ, but the decay rate for $D \to Ke\nu_e$ is more dependent on δ (see table 1), which will affect the peak of the spectrum in fig. 2. Several remarks are in order with an eye on the available data in the near future.

(1) $D \to Ke\nu_e$. We note that the spectrum for this mode alone peaks above 500 MeV with about half of the events expected above 500 MeV [12]. Compared with the average value of the data, it seems that this mode alone cannot be the dominant mode responsible for the whole data. The ratio of events above and below 500 MeV with more statistics will shed more information on this question. We also show in fig. 1 the decay spectrum due to $D \to \pi e\nu_e$, which produces more electrons with higher energies.

(2) $D \to K\pi e\nu_e$. We note from table 1 that $D \to K\pi e\nu_e$ via K^* dominance takes up a sizeable fraction of the total events. Our calculation indicates that the off-K^* mass-shell contribution is about 10% of that of the on-mass-shell contribution (as expected from narrow resonance approximation). Thus we believe that the fraction of events for D decaying into a real K^* plus $e\bar{\nu}$ is probably not small and can be experimentally looked for by searching for charged K^* resonance, having an opposite charge versus the electron (via the $\Delta s = \Delta c$ rule), in the $K\pi$ final states. The decay mode $D \to K^* e\nu_e$ provides the possibility of testing the sign of VA interference. According

to the estimate of eqs. (3) and (4), the vector contribution to $D \to K^* e \nu_e$ is about 10% of that of the axial-vector case, but the interference effect is still apparent in the electron energy spectrum. From figs. 1 and 2, one notes that V–A current peaks at lower energy than V+A, supporting the basic structure of charm changing weak currents. But if the vector contribution is suppressed, as argued above, then the interference effect may disappear.

(3) The effect of the momentum of the parent D meson will shift the peak of the spectrum to lower energy but extend the tail to higher energy.

In conclusion, we note that we have presented estimates of the most important decay modes of the D mesons and their ensuing electron momentum spectra. Our calculation shows that the $D \to K e \nu_e$ is still the dominant decay mode with $D \to K\pi e \nu_e \sim 50\% \, (D \to K e \nu_e)$ optimistically, and the higher modes contributing $\leqslant 5\%$. A definite conclusion of the analysis is that the electron spectrum in e^+e^- at $E_{cm} = 4.0$ to 4.2 GeV must decrease below $E_e \sim 350$ MeV. The large experimental errors at low electron momenta in the present data do not allow any definite conclusion to be drawn. But if the present trend of very low momentum enhancement persists, it would seem very difficult to attribute it to the charm D meson decays, at least in the conventional approach.

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