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F. Buccella and L. Oliver

May 1979

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LIMITS ON CHARM-CHANGING NEUTRAL CURRENTS*

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ABSTRACT

From data on D^0 , $\bar{D}^0$ decays, we get an upper limit on the mass difference δm of the two CP-eigenstates. Due to a critical dependence of δm on uncertain QCD corrections, this bound cannot exclude large (g_L, g_R) charm-changing neutral couplings. Limits on neutrino-production of charm by neutral currents and of charm decay into e^+e^- + hadrons are then necessary to further constrain the couplings. However, if the (c,u) current is assumed to be pure left handed, δm gives the severe bound $|g_L| < 0.6 \times 10^{-3}$.

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** On leave of absence of LPTHE, Orsay, France.

I. INTRODUCTION

The experimental suppression to orders G and $G\alpha$ of strangeness-changing neutral currents in $K_L \rightarrow \mu^+\mu^-$, $K_L - K_S$, etc., is theoretically understood as a consequence of the "natural" conservation of flavor by neutral currents.¹ Within this framework, a charm-changing neutral current is also unexpected. The present experimental limits on neutrino-production of charm and on charm semileptonic decay by neutral currents, as well as on $D^0 - \bar{D}^0$ mixing (decays into "wrong" sign kaons or leptons) indicate indeed qualitatively the absence of this current. The purpose of this note is to make a quantitative study - which lacks in the literature - of this question.

To be general, let us add an arbitrary $V, A, |\Delta C| = 1$, $\Delta Q = 0$ current

$$g_L \bar{c} \gamma_\mu \left(\frac{1-\gamma_5}{2} \right) u + g_R \bar{c} \gamma_\mu \left(\frac{1+\gamma_5}{2} \right) u + \text{h.c.} \quad (1)$$

to the standard model neutral current²

$$J_\mu^0 = \bar{\nu}_e \gamma_\mu \left(\frac{1-\gamma_5}{2} \right) \nu_e - \bar{e} \gamma_\mu \left(\frac{1-\gamma_5}{2} \right) e + 2 \sin^2 \theta_W \bar{e} \gamma_\mu e + \dots$$

The new term will then give rise to $|\Delta C| = 2$ and $|\Delta C| = 1$ pieces in the current $\times$ current lagrangian $\sqrt{2} G J_\mu^0 J_\mu^0$:

$$\mathcal{L}^{|\Delta C|=2} = \sqrt{2} G \left\{ \left[g_L \bar{c} \gamma_\mu \left(\frac{1-\gamma_5}{2} \right) u + g_R \bar{c} \gamma_\mu \left(\frac{1+\gamma_5}{2} \right) u \right]^2 + \text{h.c.} \right\} \quad (2)$$

$$\mathcal{L}^{|\Delta C|=1} = \sqrt{2} G \left\{ \left[-\bar{e} \gamma_\mu \left(\frac{1-\gamma_5}{2} \right) e + 2 \sin^2 \theta_W \bar{e} \gamma_\mu e + \right. \right. \\ \left. \left. + \bar{\nu}_e \gamma_\mu \left(\frac{1-\gamma_5}{2} \right) \nu_e + \dots \right] \left[g_L \bar{c} \gamma_\mu \left(\frac{1-\gamma_5}{2} \right) u + g_R \bar{c} \gamma_\mu \left(\frac{1+\gamma_5}{2} \right) u \right] + h.c. \right\} \quad (3)$$

This $|\Delta C| = 2$ piece would contribute to $D^0 - \bar{D}^0$ mixing, while the $|\Delta C| = 1$ piece to neutrino-production ($\nu_\mu + u + \nu_\mu + c$) and semileptonic decay ($c + u + e^+ e^-$) of charm by the neutral current.

In Section II we will deduce an upper limit on the mass difference between the two CP eigenstates from D^0 decay into "wrong" sign kaons or leptons. In Section III we shall derive a constraint on g_L and g_R from the bound on δm . We will pay special attention to the QCD corrections to δm . Further constraints on g_L and g_R will be obtained in Section IV from experimental limits on the processes $\nu_\mu + N \rightarrow \nu_\mu + \text{Charm}$ and on $\text{Charm} \rightarrow e^+ e^- + \text{Hadrons}$. Some details of the calculations will be given in the Appendices.

II. $D^0 - \bar{D}^0$ OSCILLATIONS

Let us first consider the pattern of possible decay amplitudes for a D^0 ($c\bar{u}$) in the framework of the standard model,

$$c \rightarrow s + u + \bar{d} \sim \cos^2 \theta_c \quad (4a)$$

$$c \rightarrow s + u + \bar{s} \sim \cos \theta_c \sin \theta_c \quad (4b)$$

$$c \rightarrow d + u + \bar{d} \sim \cos \theta_c \sin \theta_c \quad (4c)$$

$$c \rightarrow d + u + \bar{s} \sim \sin^2 \theta_c \quad (4d)$$

$$c \rightarrow s + \ell^+ + \nu_\ell \sim \cos \theta_c \quad (4e)$$

$$c \rightarrow d + \ell^+ + \nu_\ell \sim \sin \theta_c \quad (4f)$$

where we have omitted the quark spectator, $\bar{u}$. We expect then, from a D^0 , a single negative (positive) strangeness kaon with amplitude $\cos^2 \theta_c$ ($\sin^2 \theta_c$) and, if any, a positive charged lepton. The opposite is true for the $\bar{D}^0$. One has however to take into account the possibility of mixing between the D^0 and the $\bar{D}^0$. Indeed, a state, which at $t = 0$ is a D^0 , becomes, at time t

$$\frac{1}{2} \left[\exp(-im_+ t - \frac{\lambda_+}{2} t) + \exp(-im_- t - \frac{\lambda_-}{2} t) \right] |D^0\rangle + \\ + \frac{1}{2} \left[\exp(-im_+ t - \frac{\lambda_+}{2} t) - \exp(-im_- t - \frac{\lambda_-}{2} t) \right] |\bar{D}^0\rangle + \\ + |\text{Decay products}\rangle. \quad (5)$$

where $m_\pm$ and $\lambda_\pm$ are the masses and the inverse lifetimes of the two CP eigenstates $(D^0 \pm \bar{D}^0)/\sqrt{2}$, and one has neglected CP violation.

From Eq. (5) it is easy to deduce³ the ratio of probabilities of getting a positive or a negative electron starting from a D^0 :

$$\frac{N(D^0 \rightarrow e^- + \dots)}{N(D^0 \rightarrow e^+ + \dots)} = \frac{\delta m^2 + \delta \lambda^2/4}{2\lambda^2 + \delta m^2 - \delta \lambda^2/4} \quad (6)$$

$$\text{where } \lambda = \frac{\lambda_+ + \lambda_-}{2}, \quad \delta \lambda = \lambda_+ - \lambda_-, \quad \delta m = m_+ - m_-.$$

To compute the probability of finding "wrong" sign (positive) kaons one has to keep into account the possibility of direct decay of a D^0 into a state with positive strangeness, with amplitude proportional to $\sin^2 \theta_c$, (4d). Let us consider the decay (4a) of a D^0 into a particular state with a single strange

particle K^- , $|i\rangle = |K^- + \dots\rangle$ with amplitude α_i , proportional to $\cos^2 \theta_c$. From the decay (4d) we see that $\bar{D}^0$ can also decay into the same mode with an amplitude β_i , proportional to $\sin^2 \theta_c$. The C-transformed of this decay is the direct decay of D^0 into a "wrong" sign kaon. It is easy to obtain the probability of getting a K^- or a K^+ as the single strange particle starting from a D^0 . In Appendix I we generalize the formulas of Ref. (3) to take into account the $D^0 - \bar{D}^0$ oscillations as well as the direct decays into "wrong" sign kaons. We get:

$$N(D^0 \rightarrow K^- + \text{pions}) \propto \sum_i \left\{ \frac{\lambda}{2\lambda_+ \lambda_- (\lambda^2 + \delta m^2)} \left[\left(2\lambda^2 + \delta m^2 - \frac{\delta \lambda^2}{4} \right) |\alpha_i|^2 + \left(\frac{\delta \lambda^2}{4} + \delta m^2 \right) |\beta_i|^2 + \frac{\delta \lambda}{2\lambda_+ \lambda_-} \text{Re}(\alpha_i^* \beta_i) - \frac{\delta m}{(\lambda^2 + \delta m^2)} \text{Im}(\alpha_i^* \beta_i) \right] \right\} \quad (7)$$

$$N(D^0 \rightarrow K^+ + \text{pions}) \propto \sum_i \left\{ \frac{\lambda}{2\lambda_+ \lambda_- (\lambda^2 + \delta m^2)} \left[\left(2\lambda^2 + \delta m^2 - \frac{\delta \lambda^2}{4} \right) |\beta_i|^2 + \left(\frac{\delta \lambda^2}{4} + \delta m^2 \right) |\alpha_i|^2 \right] + \frac{\delta \lambda}{2\lambda_+ \lambda_-} \text{Re}(\alpha_i^* \beta_i) + \frac{\delta m}{(\lambda^2 + \delta m^2)} \text{Im}(\alpha_i^* \beta_i) \right\}$$

Let us now discuss the magnitude of the different terms in these formulae. Consider first the standard GIM model. β_i is proportional to $G \sin^2 \theta_c$. The width difference $\delta \lambda \sim \sum_n \langle D^0 | H_w | n \rangle \langle n | H_w | \bar{D}^0 \rangle$ is also Cabibbo suppressed, $G^2 \sin^2 \theta_c$, since the possible intermediate states come from the modes (4b) and (4c). δm is also very small in the standard

model. We give here the expression⁴ just for completeness*

$$(\delta m)_{\text{GIM}} = \frac{2}{3} \frac{G}{\sqrt{2}} f_D^2 m_D \frac{\alpha}{4\pi} \left(\frac{m_s^2 - m_d^2}{M_W^2 \sin^2 \theta_W} \right) \sin^2 \theta_c \cos^2 \theta_c.$$

Therefore, in the GIM model, the probability of getting a negative charged lepton or $S = +1$ state from the decay of a D^0 is expected to be of order $\sin^4 \theta_c$.

If we introduce the new $|\Delta C| = 2$ piece (2), we would have a contribution to δm at order G . Neglecting in (6) and (7) the terms $\delta \lambda^2$, $\delta \lambda \beta_i$, $|\beta_i|^2$ (all of order $\sin^4 \theta_c$), we obtain:

$$\frac{N(D^0 \rightarrow e^- + \dots)}{N(D^0 \rightarrow e^+ + \dots)} \cong \frac{\delta m^2}{2\lambda^2 + \delta m^2} \quad (8)$$

$$\frac{N(D^0 \rightarrow K^+ + \text{pions})}{N(D^0 \rightarrow K^- + \text{pions})} \cong \frac{\delta m^2}{2\lambda^2 + \delta m^2} + \quad (9)$$

$$+ \frac{2\delta m \lambda}{2\lambda^2 + \delta m^2} \sum_i \text{Im}(\alpha_i^* \beta_i).$$

Although we have assumed CP-invariance and Born approximation for the weak interaction, the last term in (9) (which is potentially important since λ is not Cabibbo suppressed) could be non-zero due to final state strong interactions. Therefore, we will keep it.

* The factor $2/3$, often forgotten in the literature, comes from color (see below).

Let us now look for the experimental limits on $D^0 - \bar{D}^0$ mixing. The best limit is given by the reaction⁵ $e^+e^- \rightarrow e^\pm e^\pm + \dots$:

$$\frac{N(e^+e^+) + N(e^-e^-)}{N(e^+e^+) + N(e^+e^-) + N(e^-e^-)} < 5\% \quad (90\% \text{ CL})$$

$$(3.72 \text{ GeV} < E_{\text{cm}} < 4.14 \text{ GeV})$$

This implies, from Eq. (8)

$$\frac{\delta m^2}{2\lambda^2 + \delta m^2} < 0.025. \quad (10)$$

This upper limit is far beyond the prediction of the standard model $\delta m^2/\lambda^2 \sim \sin^4 \theta_c$. The theoretical estimates of the D lifetime range from 0.7×10^{-12} sec (Ref. 6) to 1.2×10^{-12} sec (Ref. 7). The present experimental limits⁸ are 0.50×10^{-12} sec $< \tau(D) < 1. \times 10^{-12}$ sec. From the value 0.75×10^{-12} sec., we get

$$|\delta m| < 0.25 \times 10^{12} \text{ sec}^{-1}. \quad (11)$$

From (10) and (11) we can bound the r.h.s. of Eq. (9), if we make the reasonable assumption $|\beta_1| \approx \tan^2 \theta_c |\alpha_1|$. We predict:

$$\frac{N(D^0 \rightarrow K^+ + \text{pions})}{N(D^0 \rightarrow K^- + \text{pions})} < 3.8\%$$

This limit is consistent with the experimental upper limits:

$$(a) \text{ Ref. 9. } e^+e^- \rightarrow D^{*+} + \dots, D^{*+} \rightarrow K^+ \text{ or } K^-.$$

$$\frac{N(\text{wrong sign kaons})}{N(\text{all } D^{*+} \text{ events})} < 16\% \quad (90\% \text{ CL})$$

$$(E_{\text{cm}} > 5 \text{ GeV})$$

$$(b) \text{ Ref. 10. } e^+e^- \rightarrow D^0 K^+ + \dots$$

$$\frac{N(\text{wrong sign kaons})}{N(\text{all } D^0 \text{ events})} < 18\% \quad (90\% \text{ CL})$$

$$(E_{\text{cm}} = 4.028 \text{ GeV})$$

III. LIMITS ON THE COUPLINGS FROM $D^0 - \bar{D}^0$ MIXING.

Let us now compute δm from the effective interaction (2).

Without considering QCD corrections and with the vacuum insertion as the only intermediate state, we get

$$\delta m = \frac{G}{\sqrt{2}} m_D^2 \left\{ (g_L - g_R)^2 + \frac{1}{3} \left[g_L^2 + g_R^2 + 4g_L g_R \left(\frac{m_D}{m_c + m_u} \right)^2 \right] \right\}. \quad (13)$$

There are two graphs (Fig. 1): annihilation and scattering (Z^0 exchanged respectively in the s- or t-channels), corresponding respectively to the two terms in the bracket. The factor 1/3 comes from color, since to insert the vacuum we must perform a Fierz transformation of the scattering term. Moreover, the Lorentz forms $(V-A) \cdot (V-A)$ and $(V+A) \cdot (V+A)$ are invariant under Fierz, but $(V-A) \cdot (V+A)$ gives $2(P+S) \cdot (P-S)$. This explains the factor 4, the sign in front of it, and the mass ratios in the last term, since to evaluate the matrix elements of the pseudoscalar density¹¹ we use

$$q_u \langle 0 | A_\mu | D^0 \rangle = (m_c + m_u) \langle 0 | P | D^0 \rangle .$$

Note that only the axial piece contributes to the annihilation term. But, due to the Fierz transformation that we must perform, the scattering term gives also a vector contribution. f_D is defined according to:

$$\langle 0 | \bar{c} \gamma_\mu \gamma_5 u | D^0 \rangle = i f_D \frac{p_\mu^D}{\sqrt{2m_D}}$$

Let us now consider the QCD corrections. To evaluate all leading powers of $\alpha_s \log(M_W^2/m_c^2)$ we use the renormalization group method.¹² Since we have assumed a general V, A interaction, the anomalous dimension matrix is not diagonal, unlike the case of non-leptonic decays through the usual (V-A) · (V-A) interaction. We give details of the calculations in the Appendix II. Adopting the obvious notation $O_1 = (V-A)^2$, $O_2 = (V+A)^2$, $O_3 = (V-A) \cdot (V+A)$, $O_4 = (P-S) \cdot (P+S)$, the operators affected by a multiplicative renormalization

$$U_i \rightarrow U_i \left[1 + \frac{25}{3} \frac{\alpha_s(m_c^2)}{4\pi} \log \frac{M_W^2}{m_c^2} \right] \frac{\gamma_{U_i}}{25}$$

are the following combinations:

$$\begin{array}{lll} U_1 = O_1 & \text{with} & \gamma_1 = -2 \\ U_2 = O_2 & \text{with} & \gamma_2 = -2 \\ U_3 = O_3 - \frac{2}{3} O_4 & \text{with} & \gamma_3 = -1 \\ U_4 = O_4 & \text{with} & \gamma_4 = +8 \end{array}$$

Writing the original Lagrangian in terms of U_i and using (14) we get an effective Lagrangian:

$$\begin{aligned} \mathcal{L}_{\text{eff}}^{|\Delta C|=2} = & \frac{G}{2\sqrt{2}} \left\{ g_L^2 f_1 \left[\bar{c} \gamma_\mu (1 - \gamma_5) u \right]^2 + g_R^2 f_1 \left[\bar{c} \gamma_\mu (1 + \gamma_5) u \right]^2 + \right. \\ & + 2g_L g_R \left[f_2 \left(\left[\bar{c} \gamma_\mu (1 - \gamma_5) u \right] \left[\bar{c} \gamma_\mu (1 + \gamma_5) u \right] - \frac{2}{3} \left[\bar{c} (\gamma_5 - 1) u \right] \cdot \right. \right. \\ & \left. \left. \cdot \left[\bar{c} (\gamma_5 + 1) u \right] \right) + f_3 \frac{2}{3} \left[\bar{c} (\gamma_5 - 1) u \right] \left[\bar{c} (\gamma_5 + 1) u \right] \right] + \text{h.c.} \right\} \end{aligned}$$

where f_1, f_2 and f_3 are strong interaction factors, related to the quantity f_+ which appears in non-leptonic hyperon decays^{12,6},

$$f_1 = f_+, \quad f_2 = f_+^{1/2}, \quad f_3 = f_+^{-4} \quad (16)$$

with

$$f_+ = \left[1 + \frac{25}{3} \frac{\alpha_s(m_c^2)}{4\pi} \log \left(\frac{M_W^2}{m_c^2} \right) \right]^{-\frac{6}{25}} \quad (17)$$

From (15) we get

$$\begin{aligned} \delta m = & \frac{G}{\sqrt{2}} m_D f_D^2 \left\{ \frac{4}{3} (g_L^2 + g_R^2) f_1 + 2g_L g_R \cdot \right. \\ & \left. \cdot \left[-\frac{1}{9} (8f_2 + f_3) + \frac{2}{3} f_3 \left(\frac{m_D}{m_c + m_u} \right)^2 \right] \right\} \quad (18) \end{aligned}$$

which of course reduces to (13) in the limit $\alpha_s \rightarrow 0$.

For f_D we adopt the value $f_D = 180$ MeV. We discuss this choice in Appendix III. Taking $m_D = m_c + m_u$, we get the domain from (11),

$$-0.3 \times 10^{-6} < A(g_L^2 + g_R^2) + B g_L g_R < 0.3 \times 10^{-6}$$

where $A = 4 f_1/3$, $B = 2(5f_3 - 8f_2)/9$. This domain is very sensitive to the value of the QCD corrections within the expected theoretical range. For $f_+ < 0.75$ - as is the case $f_+ = 0.68$ from Ref. 6 ($\alpha_s = 0.7$), plotted in Figure 2 - this domain lies between two hyperbolas. For $f_+ = 0.75$, these become straight lines and only the vector part is bounded. When $f_+ > 0.75$, the domain closes to an ellipse, as shown in the figure for the case without QCD corrections ($f_+ = 1$). Strictly speaking, large g_L, g_R couplings are not excluded. Due to this critical dependence on the value of f_+ , $D^0 - \bar{D}^0$ mixing cannot exclude large (dominantly axial-vector) couplings.

However, if we assume a pure left-handed (c,u) current (as was the case for the (s,d) current in the Cabibbo theory before the GIM mechanism), we get the severe limit ($f_+ > 0.6$):

$$|g_L| < 0.6 \times 10^{-3} \quad (19)$$

IV. LIMITS FROM OTHER PROCESSES

Further constraints on the couplings in (1) can be obtained from limits on neutrino production and semileptonic decay of charm by neutral currents. The relevant interaction is given by (3). We get, in the valence quark parton model:

$$R = \frac{\sigma(\nu_\mu N \rightarrow \nu_\mu C)}{\sigma(\nu_\mu N \rightarrow \nu_\mu X)} = \frac{\frac{1}{4}(g_L^2 + \frac{1}{3}g_R^2)}{\left(\frac{1}{2} - \sin^2 \theta_W + \frac{20}{27} \sin^4 \theta_W\right)} \quad (20)$$

$$\bar{R} = \frac{\sigma(\bar{\nu}_\mu N \rightarrow \bar{\nu}_\mu C)}{\sigma(\bar{\nu}_\mu N \rightarrow \bar{\nu}_\mu X)} = \frac{\frac{1}{4}(g_L^2 + 3g_R^2)}{\left(\frac{1}{2} - \sin^2 \theta_W + \frac{20}{9} \sin^4 \theta_W\right)}$$

where N is an isoscalar target, and C(X) denote final states with (without) charm.

From the experimental limit¹³

$$R < 2\%$$

we get, with $\sin^2 \theta_W = 0.25$,

$$g_L^2 + \frac{1}{3}g_R^2 < 0.025. \quad (21)$$

In a similar way we can compute

$$\frac{\Gamma(c \rightarrow u + e^+ e^-)}{\Gamma(c \rightarrow s + e^+ \nu_e) + \Gamma(c \rightarrow d + e^+ \nu_e)} = \frac{1}{4}(g_L^2 + g_R^2) \times (1 - 4 \sin^2 \theta_W + 8 \sin^4 \theta_W), \quad (22)$$

which, within the parton model, is expected to be equal to the ratio¹⁴

$$\frac{\Gamma(\text{Charm} \rightarrow e^+ e^- + \text{hadrons})}{\Gamma(\text{Charm} \rightarrow e^+ \nu_e + \text{hadrons})} < 2\%.$$

We get the much less strict bound:

$$g_L^2 + g_R^2 < 0.16. \quad (23)$$

As an illustration we plot in Fig. 3 domain (21) together with the region allowed by $D^0 - \bar{D}^0$ mixing due to uncertainties in the QCD corrections when $f_+ > 0.6$.

If the (c,u) current is assumed to be pure left-handed, the limit from $D^0 - \bar{D}^0$ mixing (19) predicts:

$$\frac{\sigma(\nu_\mu N \rightarrow \nu_\mu C)}{\sigma(\gamma_\mu N \rightarrow \nu_\mu X)} < 3 \times 10^{-7}$$

$$\frac{\sigma(\bar{\nu}_\mu N \rightarrow \bar{\nu}_\mu C)}{\sigma(\bar{\nu}_\mu N \rightarrow \bar{\nu}_\mu X)} < 2 \times 10^{-7}$$

$$\frac{\Gamma(\text{Charm} \rightarrow e^+ e^- + \text{hadrons})}{\Gamma(\text{Charm} \rightarrow e^+ \nu_e + \text{hadrons})} < 3 \times 10^{-8}$$

V. CONCLUSION

We have obtained bounds on the strengths of a charm-changing neutral current $g_L(c,u)_L + g_L(c,u)_R$. Due to the critical dependence on the QCD corrections, $D^0 - \bar{D}^0$ cannot exclude large couplings, and the experimental limits on neutrino-production of charm by neutral currents are necessary to further constrain the general current. However, assuming a pure left-handed coupling, $D^0 - \bar{D}^0$ mixing gives the severe bound $|g_L| < 0.6 \times 10^{-3}$.

APPENDIX I.

We deduce here the formulae (7), which take into account $D^0 - \bar{D}^0$ oscillations as well as direct $D^0(\bar{D}^0)$ decays into "wrong" sign kaons K^+ (K^-). This is a simple generalization of the formalism of reference 3.

$$\begin{aligned} \text{A } D^0 \text{ produced at } t = 0 \text{ evolves as} \\ a(t)|D^0\rangle + b(t)|\bar{D}^0\rangle + \sum_i c_i(t)|X_i^-\rangle + \sum_i d_i(t)|X_i^+\rangle + \\ + |\text{Other decay products}\rangle \end{aligned}$$

where $|X_i^-\rangle$ is a decay product containing a K^- , and $|X_i^+\rangle$ is a decay containing a K^+ , CP transformed one of another. The amplitudes $a(t)$, $b(t)$ are given by

$$\begin{Bmatrix} a(t) \\ b(t) \end{Bmatrix} = \frac{1}{2} \begin{bmatrix} e^{-im_+t - \lambda_+t/2} + e^{-im_-t - \lambda_-t/2} \\ e^{-im_+t - \lambda_+t/2} - e^{-im_-t - \lambda_-t/2} \end{bmatrix}$$

$m_\pm$ and $\lambda_\pm$ are the masses and inverse lifetimes of the two CP-eigenstates. A state $|X_i^-\rangle$ could come from the decay of a D^0 through the process (4a) with amplitude α_i , or from the decay of a $\bar{D}^0$ through (4d) with amplitude β_i . The reverse is true for the state $|X_i^+\rangle$. We have then the differential equations for the amplitudes of probability:

$$\frac{d}{dt} c_i(t) = \alpha_i a(t) + \beta_i b(t) - i E_i c_i(t)$$

$$\frac{d}{dt} d_i(t) = \alpha_i b(t) + \beta_i a(t) - i E_i d_i(t)$$

The solutions are, for $t \rightarrow \infty$:

$$\begin{pmatrix} c_i \\ d_i \end{pmatrix} \xrightarrow{t \rightarrow \infty} e^{-iE_i t} \left\{ \left(\frac{\alpha_i + \beta_i}{2} \right) \frac{1}{\left[\frac{\lambda_i^+}{2} + i(m_+ - E_i) \right]} + \left(\frac{\alpha_i - \beta_i}{2} \right) \frac{1}{\left[\frac{\lambda_i^-}{2} + i(m_- - E_i) \right]} \right\}$$

The intensity of "wrong" (K^+) or right (K^-) sign kaons is then given by

$$N(D^0 \rightarrow \begin{Bmatrix} K^+ \\ K^- \end{Bmatrix} + \text{pions}) = \sum_i \int \left\{ \frac{|d_i|^2}{|c_i|^2} \right\} dE_i.$$

From these expressions, integrating over the Breit-Wigner's we get the formulae (7).

APPENDIX II.

We sketch here the calculation of the QCD corrections. We start from the Lagrangian (2). We have two types of interactions, $(V \pm A)^2$ and $(V - A) \cdot (V + A)$. Computing the lowest order graphs (Fig. 4) in the $\alpha_s \log$ approximation, we get the variation of these operators. The $(V \mp A)^2$ operators become

$$[\bar{c} \gamma_\mu (1 \mp \gamma_5) u]^2 - \frac{3}{2} \beta [\bar{c} \lambda^\alpha (1 \mp \gamma_5) u]^2$$

and the $(V - A) \cdot (V + A)$

$$[\bar{c} \gamma_\mu (1 - \gamma_5) u] \cdot [\bar{c} \gamma_\mu (1 + \gamma_5) u] + \frac{3}{2} \beta [\bar{c} \lambda^\alpha \gamma_\mu (1 - \gamma_5) u] \cdot [\bar{c} \lambda^\alpha \gamma_\mu (1 + \gamma_5) u]$$

$$\text{where } \beta = \frac{\alpha_s}{4\pi} \log \left(\frac{M_W^2}{m_c^2} \right)$$

Of course, in this case the vertex correction and self energy graph cancel. This result can be expressed without the color matrices by performing a Fierz transformation:

$$(\bar{q}_4 q_3)(\bar{q}_2 q_1) = \frac{1}{3} (\bar{q}_4 q_1)(q_2 q_3) + \frac{1}{2} (\bar{q}_4 \lambda^\alpha q_1)(\bar{q}_2 \lambda^\alpha q_3)$$

$(V \pm A)^2$ are Lorentz forms invariant under Fierz, but $(V - A) \cdot$

$(V + A)$ gives $2 \cdot (P + S) \cdot (P - S)$. We get then for the $(V \pm A)^2$ terms:

$$[\bar{c} \gamma_\mu (1 \mp \gamma_5) u]^2 (1 - 2\beta)$$

and for the $(V - A) \cdot (V + A)$ piece:

$$[\bar{c} \gamma_\mu (1 - \gamma_5) u] \cdot [\bar{c} \gamma_\mu (1 + \gamma_5) u] (1 - \beta) + 6\beta [\bar{c}(\gamma_5 + 1)u] \cdot [\bar{c}(\gamma_5 - 1)u].$$

The anomalous dimension matrix is thus non-diagonal. We need to compute the variation of $[\bar{c}(\gamma_5 + 1)u] \cdot [\bar{c}(\gamma_5 - 1)u]$ under lowest order gluon corrections. For this operator we get that the box diagrams (c) and (d) cancel in the leading log approximation, and the self energy and vertex graphs give the same operator:

$$[\bar{c}(\gamma_5 + 1)u] \cdot [\bar{c}(\gamma_5 - 1)u] (1 + 8\beta).$$

This means that a general interaction of the form $L = \sum_i a_i O_i$, where $O_1 = (V - A)^2$, $O_2 = (V + A)^2$, $O_3 = (V - A) \cdot (V + A)$, $O_4 = (P + S) \cdot (P - S)$, which can be represented as the vector $\mathcal{L} = \begin{pmatrix} a_1 \\ a_2 \\ a_3 \\ a_4 \end{pmatrix}$, is changed by lowest order gluon corrections by an amount

$$\delta \mathcal{L} = \frac{\alpha_s}{4\pi} \log \left(\frac{M_W^2}{m_c^2} \right) A \cdot \mathcal{L}$$

where the matrix A is given by

$$A = \begin{pmatrix} -2 & 0 & 0 & 0 \\ 0 & -2 & 0 & 0 \\ 0 & 0 & -1 & 0 \\ 0 & 0 & 6 & 8 \end{pmatrix}.$$

This matrix has the eigenvalues $\gamma_i = -2, -2, -1, 8$, corresponding to the eigenvectors $U_1 = O_1$, $U_2 = O_2$, $U_3 = O_3 - \frac{2}{3} O_4$, $U_4 = O_4$. Summing up all leading log orders, these operators will be affected by a multiplicative renormalization

$$U_i \rightarrow U_i \left(1 + \frac{25}{3} \frac{\alpha_s(m_c^2)}{4\pi} \log \frac{M_W^2}{m_c^2} \right)^{\frac{3\gamma_i}{25}}.$$

APPENDIX III.

Let us discuss the value of f_D that we have adopted. In the non-relativistic limit, this quantity is given in terms of the wave function at the origin by the Van Royen and Weisskopf formula¹⁵

$$f_D = \frac{2|\psi_D(0)|}{\sqrt{m_D}}.$$

A measurement of f_D - very difficult due to the small BR expected for $\mu\nu$ - would give very interesting information on the wave function of a system made up of a very light quark and a heavy quark. This quantity will be indicative of the type of $SU(4)$ -breaking forces. Here we cannot do better than try to get a value from indications on these forces from light mesons. The generalization to $SU(4)$ of the Weisskopf-Van Royen "paradox"

$$|\psi_\pi(0)|/|\psi_K(0)| \cong \sqrt{m_\pi}/\sqrt{m_K}$$

which ensures $f_\pi \cong f_K$ would give a value for f_D within the same range: $f_D \sim 140$ --- 180 MeV. Another estimation could come from the empirical formula¹⁶ for the vector meson wave function at the origin from $V \rightarrow e^+e^-$,

$$|\psi_V(0)|^2 \propto m_V^{-1.89}.$$

Note that this formula applies to vector mesons in which q and $\bar{q}$ have the same mass, and could be invalid for the pseudoscalar $c\bar{u}$ system. Assuming however the same mass dependence, we have $f_D = 180$ MeV. Within the uncertainties of these estimations, this value - adopted in the text - seems to us reasonable.

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FIGURE CAPTIONS

- Fig. 1: Annihilation and scattering contributions to the $D^0-\bar{D}^0$ transition.
- Fig. 2: Domain from the bound on $|\delta m|$ as a function of the QCD corrections.
- Fig. 3: Domain allowed by $\nu_\mu N \rightarrow \nu_\mu C$ and from δm taking into account the uncertainties on the QCD corrections.
- Fig. 4: Lowest order gluon corrections to the $|\Delta C| = 2$ interaction.

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Figure 1

Figure 2

Figure 3

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Figure 4

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