Flavor-changing neutral currents: Theory and future experiments

Martin Gorn*

Institut für Theorie der Elementarteilchen, Fachbereich Physik, Freie Universität, Arnimallee 3, D-1000 Berlin 33, Germany

(Received 27 November 1978; revised manuscript received 9 April 1979)

An SU(2) consistent introduction of particular flavor changes in the neutral current of a generalized Weinberg-Salam model for hadrons is described and examples are discussed, in particular, two *CP*-violating and flavor-changing five-quark models, which are in agreement with measurements on the K^0-$\bar{K}^0$ and the D^0-$\bar{D}^0$ system and still allow $t \leftrightarrow c$ and $b \leftrightarrow s$ transitions. The evidence for the strong suppression of a neutral $u \leftrightarrow c$ transition from D^0-$\bar{D}^0$ measurements is briefly reviewed. The experimental signatures of the different flavor changes, to be detected or to be ruled out in the future, are discussed.

INTRODUCTION AND SUMMARY

Strangeness conservation in weak neutral currents (NC's) motivated the prediction of a new quantum number, charm, by Glashow, Iliopoulos, and Maiani (GIM),[1] yielding for the Weinberg-Salam model[2,3] a NC completely diagonal in all flavors considered up to that time: u, d, s, c. Charmed particles have been detected[4] and are studied in various experiments, leaving no doubt about the good agreement with the GIM predictions as far as charged-current reactions are concerned. For neutral currents, the GIM model fitted strangeness conservation, as observed in kaons, but also predicted charm conservation. For some time, it was not quite clear whether charm was experimentally conserved[5] because not all the tests used in kaons could be applied to charmed particles and perhaps the interpretation of measurements for charmed particles as obvious as for kaons. Meanwhile strong suppression of charm-changing NC's seems to be accepted[6]; Sec. II of this article reviews very briefly how the absence of D^0-$\bar{D}^0$ mixing is interpreted in this way.

The GIM mechanism and its generalization to more than four quark is not the only way to suppress strangeness change in NC. Many models have been proposed conserving strangeness but changing charm. Now charm is known to be conserved but another flavor has been observed, attributed to the $\Upsilon(9.5)$ resonance in e^+e^- (Ref. 7) and the same question arises again: Are there neutral transitions from this new flavor to the old ones? Section I of this article therefore deals with the more general restrictions on the multiplet structure of an extended Weinberg-Salam model if particular flavors are conserved and others not [in the case of SU(2) doublets and singlets]. As examples, out of the Kobayashi-Maskawa (KM) model[8] two alternative five-quark models with flavor change of $V-A$ type in the NC are obtained. It turns out that even if the neutral flavor-changing (FC) transitions $s \leftrightarrow d$ and $c \leftrightarrow u$ are highly suppressed, a considerable coupling strength for $b \leftrightarrow s$ and $t \leftrightarrow c$, respectively, could be possible. Comparing limits eventually obtainable on these couplings by future experiments (Sec. III), expressed by the KM parameters θ_1, θ_2, θ_3, and δ with other determinations of these parameters from b or t decays and *CP* violation, one could perhaps rule out a five-quark version before experimental signs of a sixth quark have been seen.

Every time detection of a new flavor is reported, a model of weak interactions may be constructed within the lines of Sec. I, having neutral currents changing this new flavor to old ones. The experimental possibilities to support or to contradict such a theory are discussed in Sec. III: First, neutral-mixing experiments can yield limits as low as in the previous cases, provided that the second-order mixing is small enough and that the experimental signature is well defined. Therefore this test might fail for the $\bar{t}$-c system. In deep-inelastic neutrino-nucleon scattering wrong-sign dimuons could indicate neutral FC transition off nucleon quarks. If neutral-mixing signatures are correct, like-sign dimuon rates can give good upper limits. By means of the invariant-mass spectrum of $\mu^+\mu^-$ pairs in trimuons off-diagonal neutral particles could be seen or excluded if their mass is known. Finally the direct coupling of off-diagonal neutral particles to the NC in e^+e^-, as proposed,[9] is briefly mentioned. Monochromatic γ's or π^0's are expected at certain values of s, with signal sizes of (0.25–5)%. Considerations on the width of vector mesons and their mass difference to the corresponding pseudoscalar show that the signal is in principle much higher and in reality depends on the beam resolution. Remarks on bremsstrahlung background and inclusive γ spectroscopy show that this kind of experiment is not unrealistic, particularly when other tests turn out to be less efficient. In the appendices the leptonic widths of D^0 and D^{*0} as well as some estimates on FCNC form factors are given.

I. WEAK-INTERACTION THEORIES WITH AND WITHOUT FLAVOR CHANGE IN THE NC

This discussion will be started by presenting the neutral current for a generalized Weinberg-Salam model,[2,3] using SU(2) × U(1), extended to hadrons,[10] with left-handed and right-handed fermion multiplets of SU(2):

$$j_\mu^N = \bar\psi\gamma_\mu C_L^3(1-\gamma_5)\psi + \bar\psi\gamma_\mu C_R^3(1+\gamma_5)\psi - 2\sin^2\theta_W j_\mu^{em}. \tag{1.1}$$

θ_W is the Weinberg angle, j_μ^{em} the electromagnetic current. C_L^3 and C_R^3 are the third components out of a representation (reducible) of SU(2), which means that they are obtained from the so-called charged currents

$$j_\mu^\pm = \bar\psi\gamma_\mu C_L^\pm(1-\gamma_5)\psi + \bar\psi\gamma_\mu C_R^\pm(1+\gamma_5)\psi \tag{1.2}$$

by

$$[C_R^+, C_R^-] = 2C_R^3, \quad [C_L^+, C_L^-] = 2C_L^3 \tag{1.3}$$

for R and L independently, acting on the right- and left-handed multiplets, respectively. The NC (1.1) couples to the weak neutral intermediate boson Z^0 as

$$\mathcal{L}_{int} = i g_N j_\mu^N Z^{0\mu}, \tag{1.4}$$

where

$$g_N = \frac{e}{2\sin\theta_W\cos\theta_W} = \left(\frac{G}{\sqrt2}\right)^{1/2} M_Z\sqrt2,$$

$$\frac{G}{\sqrt2} \approx 8\times10^{-6}\,\text{GeV}^{-2}.$$

As long as only doublets of the form

$$\begin{pmatrix}\psi_q\\ \psi_{q-1}\end{pmatrix},$$

q standing for the em charge of the field, and singlets are considered, the representation matrices C^i $(i=1,2,3)$ $(C^\pm = C^1 \pm iC^2)$ can be expressed in terms of the well known 2×2 Pauli matrices[11]:

$$C_{L,R}^+ = \tau^+ \otimes D_{L,R} \quad \text{for } R \text{ and } L \text{ separately}, \tag{1.5}$$

meaning that all doublets are written together in one multiplet belonging to a reducible representation. Note that the matrices D_L and D_R can be extracted from charged-current reactions, see (1.2), but are also restricted by the group structure. From the commutation relations for SU(2) one obtains for the D's defined in (1.5) (Refs. 11, 12) the condition

$$DD^\dagger D = D \quad (R \text{ and } L \text{ separately}). \tag{1.6}$$

The neutral current (1.1) is obtained from

$$C_{L,R}^3 = \tfrac12\tau^+\tau^-\otimes D_{L,R}D_{L,R}^\dagger - \tfrac12\tau^-\tau^+\otimes D_{L,R}^\dagger D_{L,R}. \tag{1.7}$$

But the products $\tau^+\tau^-$ and $\tau^-\tau^+$ are just projection operators on each of the two charges appearing in the doublet. This projection operator property holds only in special dimensions (representations) because it depends on what the anticommutators are:

$$T^+T^-T^+T^- = T^+(-\tfrac12[T^+,T^-] + \tfrac12\{T^+,T^-\})T^-.$$

But only the commutators are fixed by the algebra. Therefore in the case of doublets and singlets the neutral current, as indicated in (1.7), decomposes into two parts, one with quarks having all charge q and the neutral-current transitions among which are given by the matrix $DD^\dagger$ and the other with charge $q-1$, whose transitions are given by $D^\dagger D$. Note that in general $DD^\dagger$ and $D^\dagger D$ are independent. They directly indicate the structure of the NC for each charge sector and each helicity sector (L and R). If they are diagonal, there is no flavor change.

From (1.6) follows that if D is regular ($\det D \neq 0$) it has to be unitary and that all flavors are conserved. The same is true if D is of the form $D = \begin{pmatrix}\tilde D & 0\\ 0 & 0\end{pmatrix}$ with $\tilde D$ regular. A nontrivially singular matrix however can yield flavor change, because rearrangements of lines or columns by unitary transformations do not have the same effect on $DD^\dagger$ and $D^\dagger D$; in other words, $D^\dagger D$ and $DD^\dagger$ do not necessarily commute. They do so, for instance, when D is regular, but then there is no flavor change.

The history of the model by Glashow, Iliopoulos, and Maiani[1] can now easily be reviewed as follows: For the original three-quark version with u, d, and s (see Ref. 10) the D's defined in (1.5) are

$$D_L = \begin{pmatrix}\cos\theta_C & \sin\theta_C\\ 0 & 0\end{pmatrix} \text{ and } D_R = 0. \tag{1.8}$$

D_L is a singular matrix, satisfying (1.6). For an extended multiplet $\psi^T = (u^T, d^T, s'^T, s^T)$, D_L avoids coupling the s' $(q=\frac23)$ currents (it therefore physically appears nowhere). The products appearing in (1.7) which describe the NC structure become

$$D_L D_L^\dagger = \begin{pmatrix}1 & 0\\ 0 & 1\end{pmatrix}$$

and

$$D_L^\dagger D_L = \begin{pmatrix}0 & \cos\theta_C\sin\theta_C\\ \cos\theta_C\sin\theta_C & 0\end{pmatrix}, \tag{1.9}$$

immediately showing up the flavor change in the $q-1$ sector, the strangeness-changing NC. Now it is clear how the strangeness-conserving extension of the three-quark model has to be made: Replace (1.8) by

$$D_L = \begin{pmatrix} \cos\theta_C & \sin\theta_C \\ -\sin\theta_C & \cos\theta_C \end{pmatrix}, \quad D_R = 0\,, \tag{1.10}$$

being the only *regular* matrix satisfying (1.6). This is exactly the model by Glashow, Iliopoulos, and Maiani having no strangeness or charm change in the NC. While the first was a fit to the experimental data from K mesons, the charm conservation was a prediction. Bjorken[5] refers to some estimates[13] for

$$\mathcal{H}_{\text{eff}}^{|\Delta S|=2} = \frac{2G}{\sqrt{2}}(G\Lambda^2)\cos^2\theta_C \sin^2\theta_C\, j_\mu^{N\Delta S} j^{\mu,N\Delta S}$$

of $\Lambda \leqslant 4$ GeV, that is for the effective coupling of a strangeness-changing neutral current.

However it is possible within the framework of $SU(2)\times U(1)$ models, discussed up to now, to still have charm change in the neutral current.

Before elaborating on this point let me briefly quote the six-quark extension of the GIM model by Kobayashi and Maskawa[8] which is now in consideration because of the recently observed resonance structure near 9.5 GeV (Ref. 7) and in addition tries to parametrize CP violation. To that extent the KM predictions for CP violations agree with or differ from those of a superweak theory;[14] see Ref. 15. For other experimental consequences see Ref. 16. It has an additional doublet $\binom{t}{b}_L$ and in terms of the definition (1.5) the KM model just becomes

$$D_L = \begin{pmatrix} c_1 & -s_1c_3 & -s_1s_3 \\ s_1c_2 & c_1c_2c_3 - s_2c_3e^{i\delta} & c_1c_2s_3 + s_2c_3e^{i\delta} \\ s_1c_2 & c_1s_2c_3 + c_2s_3e^{i\delta} & c_1s_2s_3 - c_2c_3e^{i\delta} \end{pmatrix}, \tag{1.11}$$

where

$$c_i = \cos\theta_i, \quad s_i = \sin\theta_i, \quad i = 1,2,3\,,$$

the whole multiplet being

$$\psi^T = (u^T, d^T, c^T, s^T, t^T, b^T)\,.$$

As one sees from (1.11) D_L is unitary and therefore all flavors are conserved in the KM model in first order.

However, this need not be the case. Even if strangeness is conserved in the NC, there is a priori no reason why neutral transitions like $u \leftrightarrow c$, $d \leftrightarrow b$, or $c \leftrightarrow t$ should not take place already at the first-order level in $G/\sqrt{2}$. The above considerations allow insertion of such flavor-changing NC's consistently with the $SU(2)\times U(1)$ framework: Take the KM matrix D_L in (1.11) and replace for instance the last row (column) by zeros. Surprisingly, it gives rise to a new model: a five-quark model, that is KM without the t quark (b quark), and the new reduced matrix $D_L^{(5)}$ still satisfies the SU(2) restrictions (1.6). $D_L^{(5)}$ is nontrivially singular and therefore yields flavor change in the NC. The first possibility, the five-quark version with the b quark as the fifth quark, has the FCNC

$$[-s_1c_1c_3 + s_1c_2(c_1c_2c_3 - s_2s_3e^{-i\delta})]\,\bar{s}\gamma_\mu(1-\gamma_5)d + \text{H.c.},$$

$$[-s_1c_1s_3 + s_1c_2(c_1c_2s_3 + s_2c_3e^{-i\delta})]\,\bar{b}\gamma_\mu(1-\gamma_5)d + \text{H.c.}, \tag{1.12a}$$

$$[s_1^2s_3c_3 + (c_1c_2s_3 + s_2c_3e^{-i\delta})(c_1c_2c_3 - s_2s_3e^{i\delta})] \times \bar{b}\gamma_\mu(1-\gamma_5)s + \text{H.c.},$$

and *no flavor change* for the $q=\frac{2}{3}$ quarks, whereas for the second possibility, the five-quark model containing the t quark as the fifth quark, we get

$$[c_1s_1c_2 - s_1c_3(c_1c_2c_3 - s_2s_3e^{i\delta})]\,\bar{c}\gamma_\mu(1-\gamma_5)u + \text{H.c.},$$

$$[c_1s_1s_2 - s_1c_3(c_1s_2c_3 + c_2s_3e^{i\delta})]\,\bar{t}\gamma_\mu(1-\gamma_5)u + \text{H.c.}, \tag{1.12b}$$

$$[s_1^2s_2c_2 + (c_1s_2c_3 + c_2s_3e^{i\delta})(c_1c_2s_3 - s_2s_3e^{-i\delta})] \times \bar{t}\gamma_\mu(1-\gamma_5)c + \text{H.c.},$$

and *no flavor change* for the $q=\frac{1}{3}$ quarks.

Before adding some knowledge on FC parameters, note that Cabibbo universality demands (see, e.g., Refs. 15,16)

$$s_3^2 \leqslant 0.06\,. \tag{1.13}$$

Comparing these restrictions and the predictions for CP violation coming from the complex phase $e^{i\delta}$ with the results from future searches for FCNC (see Sec. II of this work) would give an indication as to whether a five-quark model may still be possible or whether it should be ruled out.

In order to make some predictions we still have to insert the strong suppression of strangeness change (SC) into (1.12a). From Refs. 5 and 13 the coefficient of the term $\bar{s}\gamma_\mu(1-\gamma_5)d$ in (1.12) is bounded by

$$s_1s_2(-c_1c_3c_2 - c_3c_2e^{-i\delta}) \lesssim 10^{-3}\,, \tag{1.14}$$

for which the two extreme cases

$$\text{SC1: } s_2 \lesssim 10^{-2}$$

or

$$\text{SC2: } \theta_2 \approx \theta_3 \text{ and } e^{-i\delta} \approx -1 \tag{1.14$'$}$$

are solutions (among others). This yields from (1.12a) the coupling strengths for $\bar{b}\gamma_\mu(1-\gamma_5)d$,

SC1: ~10^{-3},

SC2: max of 10^{-1} (depending on s_2, s_3);

(1.14″)

for $\bar{b}\gamma_\mu(1-\gamma_5)s$,

SC1: max of 0.6 (depending on s_2, s_3),

SC2: max of 0.4 (depending on s_2, s_3),

where the bound (1.13) and $s_1 \approx 0.23$ has been used.

Similarly, we insert charm conservation (CC) (reviewed in Sec. II) into (1.12b). It sets the bound

$$s_1 s_3 (s_3 c_2 c_1 + s_2 c_3 e^{i\delta}) \lesssim 4\times 10^{-4}. \qquad (1.15)$$

Two possible solutions are

CC1: $s_3 \lesssim 4\times 10^{-4}$,

(1.15′)

CC2: $\theta_2 \approx \theta_3$ and $e^{i\delta} \approx -1$.

They predict for the coupling strengths

$\bar{t}\gamma_\mu(1-\gamma_5)u$

CC1: ~4×10^{-4}

CC2: max of ~6×10^{-2},

(1.15″)

$\bar{t}\gamma_\mu(1-\gamma_5)c$,

CC1: max of 0.6 (depending on s_2, s_3),

CC2: max of 0.4 (depending on s_2, s_3).

We see that neutral transitions among the fifth quark of a five-quark model and the non-nucleonic quarks (s or c) with an appreciable strength are still consistent with present data. They are just the extensions of the original three-quark version (with only $V-A$ currents) of the original Weinberg-Salam model. In part III of this work experimental signatures for FCNC's are discussed. Their suppression in comparison with other determinations of θ_2, θ_3, and δ could rule out the above five-quark models in favor of the flavor-conserving six-quark KM model, before the quantum number or other signs of a sixth quark have really been observed experimentally. Therefore, within this construction, an inconsistency will lead to the prediction of a sixth quark, similar to the situation with charm (GIM mechanism[1]).

The above procedure of introducing neutral flavor changes in first order can obviously be applied to any model with SU(2) doublets and singlets. Other models have been proposed with FCNC's by Achiman and Walsh.[17] They can also be described by the above framework, showing that they are consistent with SU(2) in the sense of condition (1.6). For instance, their extension of the standard model by the right-handed doublet $\binom{u}{b}_R$ and the singlet c_R becomes in terms of definition (1.5)

$$D_R = \begin{pmatrix} \cos\alpha & 0 \\ \sin\alpha & 0 \end{pmatrix} \text{ and } D_L \text{ as in standard GIM}.$$

However, out of these models, those having nucleon quarks in right-handed doublets are in poor agreement with experimental data, if one combines[18] the strong suppression of the neutral $u \leftrightarrow c$ transition (Sec. II) and diagonal-neutral-current data as analyzed by Ecker.[19]

Apart from possibly predicting a new quark there is another reason why FCNC will always be an important subject of measurements. Glashow and Weinberg have formulated the concept of "natural flavor conservation",[20] requiring that, if all flavors are strongly conserved in nature, then they should be conserved independently of all the variable parameters of a theory. Chanowitz, Ellis, and Gaillard[21] have then shown that this principle of naturalness, if fully applied, could possibly (unless some peculiar effects are observed) restrict the group structure for a gauge theory of weak and em interactions to $SU(2)_L\times SU(2)_R\times U(1)$ or $SU(2)_L\times U(1)^m$, $m\leq 2$.

II. EVIDENCE FOR SUPPRESSION OF NEUTRAL CHARM CHANGE FROM D^0-$\bar{D}^0$ MIXING

Some properties of neutral-particle mixing theory and its validity of application are reviewed. Once a state $|D^0\rangle$ or $|\bar{D}^0\rangle$ is produced, there is a mixing of these states due to the off-diagonal matrix elements of the weak Hamiltonian induced by the charm-changing neutral current:

$$j_\mu^{N\Delta C} = \frac{d}{2}[\bar{c}\gamma_\mu(1-f\gamma_5)u + \bar{u}\gamma_\mu(1-f\gamma_5)c], \qquad (2.1)$$

d is the (dimensionless) coupling strength relative to the Fermi constant, depending on the model [see as an example expression (1.12b)]. Mixing theory[22] describes the time development of the states $|D^0\rangle$ and $|\bar{D}^0\rangle$ for times much larger than the decay time, meaning that the effects from relativistic quantization can after that time be neglected and that the D^0-$\bar{D}^0$ system is then to be treated as a classical one. Using the fact that the Klein-Gordon equation is equivalent to the Schrödinger equation,[22] if the commutator between the interaction term and the classical Hamiltonian is small compared to $M_{D^0}{}^2$, which is true because of Γ and $\Delta m \ll m_{D^0}$ (Γ = decay width, Δm = mass difference), one gets the well known mixing formulas, e.g.,

$$|\langle \overset{(-)}{D^0} | D^0(t)\rangle|^2 = \tfrac{1}{4}[e^{-\Gamma_1 t} + e^{-\Gamma_2 t} \underset{(-)}{+} 2e^{-(\Gamma_1+\Gamma_2)t}\cos(\Delta m t)], \qquad (2.2)$$

plus a term proportional to

$$\frac{\sqrt{\Gamma_1+\Gamma_2}\,\exp[-(p^2+m_0{}^2)^{1/2}t]}{(p^2+m_D{}^2)t^{3/2}},$$

which in fact can be neglected.

From (2.1) one can calculate $\Gamma_{1/2}$ and Δm: the result is[18]

$$\Delta m = \frac{G}{\sqrt{2}} \frac{d^2}{2} f_D{}^2 m_D \quad \left[= O\left(\frac{G}{\sqrt{2}}\right)\right] ,$$

$$\Gamma_1 + \Gamma_2 = \Gamma_0 + \Gamma_c \frac{d^2}{2} n \quad \left[= O\left(\left(\frac{G}{\sqrt{2}}\right)^2\right)\right] , \qquad (2.3)$$

$$\Delta\Gamma = \Gamma_c \left(n \tan^2\theta_C + \frac{d^2}{2} n \right) .$$

Γ_0 stands for the total D^0 width with only charged-current contributions ($d=0$):

$$\Gamma_0 = \Gamma_c (1 + \tan^4\theta_C + n \tan^2\theta_C) .$$

This is the sum of Cabibbo allowed, suppressed, and partly suppressed partial widths, the factor n counting the additional channels not present in the first two kinds of decay. But n is, by pure counting of channels and phase space, restricted by $2 \leq n \leq 4$, therefore not much of a change. f_D is the pseudoscalar form factor. The above way of writing Γ_0 has been chosen because a rather accurate estimate by Gaillard *et al.* exists for Γ_c, where all multipion channels are taken into account: $\Gamma_c = 1.65 \times 10^{-10}$ MeV.[23,24]

The fact that (2.1) gives a first-order contribution to the mass difference whereas the widths $\Gamma_{1/2}$ remain of second order in the Fermi constant is well known[25] and leads to an enormous value for the mixing parameter:

$$\frac{\Delta m}{\Gamma_1 + \Gamma_2} = 9 \times 10^5 \frac{d^2}{1 + \frac{1}{4} n d^2} \qquad (2.4)$$

(with $f_D \approx f_\pi$, which seems to be a good guess.[16] Considering the quantity, obtained from (2.2),

$$R_M = \frac{\Gamma(D^0 \to \overline{D}^0 \to \overline{X})}{\Gamma(D^0 \to X)}$$

$$= \frac{\left(\frac{\Delta\Gamma}{2(\Gamma_1+\Gamma_2)}\right)^2 + \left(\frac{\Delta m}{\Gamma_1+\Gamma_2}\right)^2}{\frac{1}{2} - \left(\frac{\Delta\Gamma}{2(\Gamma_1+\Gamma_2)}\right)^2 + \left(\frac{\Delta m}{\Gamma_1+\Gamma_2}\right)^2} , \qquad (2.5)$$

and using (2.3) to get

$$R_M = 5.2 \times 10^{13} \frac{\left(\frac{nd^2}{4+nd^2}\right)^2}{2 + 5.2 \times 10^{13} \left(\frac{nd}{4+nd}\right)^2}$$

(for d not too small) , (2.6)

$$R_M \leq \frac{\left(\frac{n \tan^2\theta_C}{1+\tan^2\theta_C}\right)^2 + \tan^4\theta_C}{2 - \left(\frac{n\tan^2\theta_C}{1 + n\tan^2\theta_C}\right)^2 + \tan^4\theta_C} \approx 5\% \quad (\text{for } d=0) ,$$

(2.6′)

one can derive upper limits for d from experiments. Goldhaber *et al.*[26] have measured the ratio of equally charged to oppositely charged kaons coming from D^0-$\overline{D}^0$ pairs produced in e^+e^-. Another measurement, by Feldman *et al.*,[27] tests the decay products of $D^{*+}D^{*-}$ pairs. The two measurements set the bound

$$\frac{d}{\sqrt{2}} \leq (4-7) \times 10^{-4} \qquad (2.7)$$

under the condition that the above estimates on the total D^0 width and the form factor turn out to be correct. A similar number can be found in Ref. 6. On the Fermi level, (2.7) corresponds to $\sim 10^{-7} G/\sqrt{2}$.

Consider finally the $\bar{u}$-c coupling (2.1) to be mediated by a neutral gauge boson which has been introduced in models for weak interactions with higher groups [e.g., SU(3), Ref. 28]. If there were neither a Clebsch-Gordan coefficient nor a mixing angle, the mass of this boson had to be larger than 3×10^3 times the Z^0 mass.

III. DETECTION OF FLAVOR CHANGE IN THE FUTURE

With every new resonance in e^+e^- or $\mu^+\mu^-$ reported and the new quarks attributed to it, a new model may be constructed having a neutral current changing known quarks of the same charge into this new quark following the rules given in Sec. I. What will be the chances to contradict or support such a theory? This section deals with the following tests on FCNC:

(1) measurements of X^0-$\overline{X}^0$ mixing,

(2) flavor-change signatures in deep-inelastic neutrino scattering,

(3) direct coupling of vector particles to the neutral current.

A. Measurements of $X^0 - \overline{X}^0$ mixing

The FCNC has the most general form

$$j_\mu^{\rm NFC} = \bar{q}\gamma_\mu (d_V - d_A \gamma_5) q' + \bar{q}'\gamma_\mu (d_V^* - d_A^* \gamma_5) q , \qquad (3.1)$$

where Hemiticity has been imposed because the Z^0 also couples to the diagonal NC, being therefore Hermitian. (This would be different in a theory with two neutral weak bosons apart from the photon; then we would have d_V' and d_A'). There we are to expect neutral off-diagonal particles X^0 with the quark content: $\bar{q}, q'$, $q \neq q'$. For simplicity only pseudoscalars are considered for mixing. In principle, also all other kinds of particles can mix, but most probably their strong decays will dominate. As in Sec. II there will be a large mass difference between X_1 and X_2:

$$\Delta m = \frac{G}{\sqrt{2}} m_{X^0} f_{X^0}{}^2 d_A{}^2 , \qquad (3.2)$$

unless d_A is very small, f_{X^0} being the pseudoscalar form factor, analogous to f_D in Sec. II. (For some estimates on pseudoscalar form factors for systems involving higher quarks, see Ref. 16.) As in Sec. II, (3.2) will lead to complete mixing similar to (2.4), yielding for the ratio R_M, which is always something like

$$R_M=\frac{\Gamma(X^0\to\overline{X}^0\to\overline{X})}{\Gamma(X^0\to X)}\,,$$

$R_M\approx 1$, unless d_A is very small, to be measured e.g. in e^+e^- from the ratio of equal pairs ($K^\pm K^\pm$, $D^\pm D^\pm$, or $e^\pm\mu^\pm$ etc.) to opposite pairs in a region where X^0 and $\overline{X}^0$ are produced as pairs. If the experimental detection efficiency (K-separation, lepton acceptance, knowledge of em background etc.) is good, already an upper limit of some 10–20% on R_M gives marvelous upper limits on the flavor change in the NC, but only if the second-order contribution from charged currents is small enough [in the D^0-$\overline{D}^0$ system almost 5%, see (2.6′)].

In contrast to the K^0-$\overline{K}^0$ system, where the parameter $\Delta m/(\Gamma_1+\Gamma_2)$ can be directly observed because the particles live long enough to show the oscillations, we are now, as in D^0-$\overline{D}^0$, in the situation to determine it at practically infinite times from

$$R_M=\frac{\left(\frac{\Delta\Gamma}{2(\Gamma_1+\Gamma_2)}\right)^2+\left(\frac{\Delta m}{\Gamma_1+\Gamma_2}\right)^2}{\frac{1}{2}-\left(\frac{\Delta\Gamma}{2\Gamma}\right)^2+\left(\frac{\Delta m}{\Gamma_1+\Gamma_2}\right)^2}\,.$$

Now Ellis *et al.* have pointed out that the parameter $\Delta m/(\Gamma_1+\Gamma_2)$ can become $m_t^2/(700\ \text{GeV}^2)$, giving for a quark of 30 GeV and Cabibbo-type angles neglected, for R_M about

$$R_M=0.7\text{–}0.9\,,$$

which of course always has an experimental error. In a recent paper, Ali and Aydin[29] point out that in the T_c^0 and B_s^0 systems ($\{t,\overline{c}\}$ and $\{b,\overline{s}\}$) this mixing is still enhanced by $\cot^2\theta_C$, already giving large mixing also for lower quark masses. Therefore already from second-order (charged-current) transitions, R_M can approach the value 1, up to now being rather unambiguously interpreted owing to first-order flavor change. A second difficulty for the application of R_M is that higher quarks, such as b for instance, decay in a cascade

$$b\to cX^-,\ c\to s\overline{Y}^+ \quad\text{and}\quad \overline{b}\to\overline{c}\overline{X}^+,\ \overline{c}\to\overline{s}Y^-\,,$$

where X^- may be $(\overline{u},d)$, $(\overline{c},s)$—the remaining $(\overline{u},s)$ and $(\overline{c},d)$ are suppressed—or $(\mu^-,\overline{\nu}_\mu)$, $(e^-,\overline{\nu}_e)$ and Y^- the same without $(\overline{c},s)$ because of phase space. Now if the mass of the b is not too near that of the c, from one cascade decay two K's of the same sign can be produced whereas the conjugate $\overline{b}$ could also give no K (for $b\to c$ need not be the dominant transition). Also $\mu^\pm e^\pm$ is not a good signature because one does not know from which step the leptons come; it will appear without flavor change as well. These difficulties are described in more detail in Ref. 29. Multilepton states are proposed as indication of mixing, having, however, very small branching ratios, and working only if the mass differences are such that momentum cutoffs can indicate the cascade step. Under some circumstances, I therefore conclude that it may be very difficult to verify mixing at all, and it may be impossible to filter out or to rule out first-order flavor-changing NC from these assignments in e^+e^-, used up to now in connection with R_M, in which case another kind of test would be desirable.

B. Flavor-change signatures in deep-inelastic neutrino-nucleon scattering

With the neutral current [(+) for ν and (−) for $\overline{\nu}$ in (3.4) and (3.5)]

$$\frac{1}{\sqrt{2}}\overline{\nu}\gamma_\mu[1-(\pm)\gamma_5]\,\nu+\frac{d}{\sqrt{2}}[\overline{c}\gamma_\mu(1-f\gamma_5)u+\text{H.c.}] \tag{3.4}$$

the process of neutral charm production [Fig. 1(a)], charm sea in the nucleon neglected, has in the parton model for deep-inelastic scattering[30] the cross section

$$\frac{d\sigma^{N\Delta C}}{dx\,dy}=\frac{G^2ME}{\pi}\frac{d^2}{4}\frac{1}{2}[u(x')+d(x')]\times\left\{\left[(1-y)2x'+\frac{y^2}{2}2x\right]\frac{1+f^2}{2}\pm 2xy\left(1-\frac{y}{2}\right)f\right\}\theta(Ey-m_c)\,,$$

$$\frac{d\sigma_{\overline{c}}^{N\Delta C}}{dx\,dy}=\frac{\overline{u}(x')+\overline{d}(x')}{u(x')+d(x')}\,\frac{d\sigma_c^{N\Delta C}}{dx\,dy}\ (\text{replace } f\to -f)\,, \tag{3.5}$$

where $x'=x-m_c^2/2M(E-E')$. For all other variables see the literature. What we call "charm" in

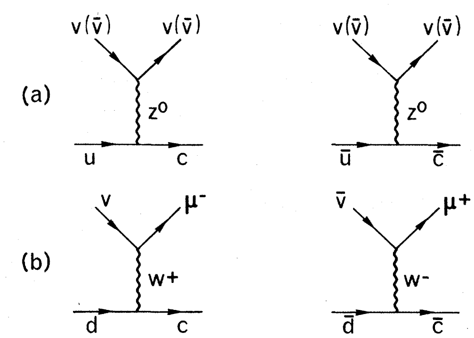

FIG. 1. (a) Neutral charm production, (b) charged charm production in ν or $\overline{\nu}$.

this section may be immediately replaced by any other flavor of $q=\frac{2}{3}$. For a $q=-\frac{1}{3}$ flavor the roles of μ^+ and μ^- just have to be interchanged.

This charm ($_c$) [anticharm ($\sigma_{\bar{c}}$)] production with the subsequent semileptonic decay contributes to the one-muon [one-antimuon] rate. Briefly I note that in principle the kinematics of such muons can be calculated, e.g., by the method described by Seghal and Zerwas,[31] which has already been successfully applied to dimuon kinematics. In order to get the gross features however, we integrate over all variables leaving in the formulas the semileptonic branching ratio B for charmed particles, which is known to be about 10% [CERN-Dortmund-Heidelberg-Saclay (CDHS)].[32]

Neglecting the effect of mixing ($D^0 \to \bar{D}^0$) which changes the μ^+ or μ^- production at most by about 10%,[18] the contribution from ν and $\bar{\nu}$ (every beam has both, one as contamination) to the muon rate by the neutral current (3.4) is

$$\sigma_{\mu^+}^{N\Delta C}=n(\nu)\,\sigma_c^{N\Delta C,\nu}+n(\bar{\nu})\,\sigma_c^{N\Delta C,\bar{\nu}}\,,$$

$$\sigma_{\mu^-}^{N\Delta C}=n(\nu)\,\sigma_{\bar{c}}^{N\Delta C,\nu}+n(\bar{\nu})\,\sigma_{\bar{c}}^{N\Delta C,\bar{\nu}}\,,$$

which in terms of (3.5) is

$$\sigma_{\mu^+}^{N\Delta C}=\frac{G^2ME}{\pi}\frac{d^2}{4}BN_{\rm val}\tfrac{1}{4}\left[\left(n(\bar{\nu})+\frac{n(\nu)}{3}\right)(1-f)^2+\left(n(\nu)+\frac{n(\bar{\nu})}{3}\right)(1+f)^2\right], \tag{3.6}$$

$$\sigma_{\mu^-}^{N\Delta C}=\frac{N_{\rm sea}}{N_{\rm val}}\,\sigma_{\mu^+}^{N\Delta C}[\text{with } n(\nu) \text{ and } n(\bar{\nu}) \text{ interchanged}]\,.$$

We now discuss the implications of (3.6).

$f=+1$ (V-A): μ^+ appear in a beam with dominantly *neutrinos* as "wrong-sign" muons (in contrast to charged-current induced μ^-) whereas faked μ^+ in an antineutrino beam appear with a factor of $\frac{1}{3}$ (relative to the wrong-sign μ^+ contribution).

$f=-1$ ($V+A$): The wrong-sign μ^+ rate in a neutrino beam is three times smaller than the contribution of faked μ^+ to the normal (charged current) μ^+ production in an antineutrino beam.

μ^- (neutrally induced) come only from anticharm and therefore from the sea. μ^+ in neutrino (faked) could be distinguished from their different kinematics. Looking up the results of Seghal and Zerwas for dimuon production, the following characteristics are obtained:

(1) faked muon production rises logarithmically in energy compared to the normal one-muon rate:

$$\frac{\sigma(\mu^+,N\Delta C)}{\sigma(\mu^-,\ \text{charged current})}\sim\ln\frac{E}{m_c}\,. \tag{3.7}$$

The logarithm comes from the used fragmentation function $1-Z/Z$, see Ref. 31.

(2) The distribution of the visible y,

$$y_{\rm vis}=\frac{E_{\rm hadr}}{E_{\rm hadr}+E_\mu}$$

would be as in Fig. 2 (drawing only schematic).

The contribution to high y comes from the small energy of the muon, faking a high energy transfer to the hadronic system from the neutrino vertex. But the contribution to a "high-y anomaly", reported in Ref. 33 and recently cast into doubt,[32] would be

$f=-1$ ($V+A$):

$$\frac{\sigma(\mu^+,y\approx 1,N\Delta C)}{\sigma(\mu^+,y\approx 0,\ \text{charged current})}\approx\frac{B}{4}\ln\frac{E}{m_c}\,,$$

giving ≈10% at $E=100$ GeV, if $d=1$. Therefore at most a high-y anomaly of 10% could be explained.

Interesting could be the ratio

$$\frac{\sigma(\mu^+\text{in}\nu)}{\sigma(\mu^-\text{in}\bar{\nu})}=\frac{N_{\rm val}}{N_{\rm sea}}\frac{\frac{1}{3}(1-f)^2+(1+f)^2}{(1-f)^2+\frac{1}{3}(1+f)^2}\,. \tag{3.8}$$

Because there is always an background of wrong-sign muons, the ratio (3.8), if large (of the order of 10–30), would be an indication of a charm change in the NC. The CDHS group gave numbers on wrong-sign muons (Tittel in Ref. 32):

$$\frac{\sigma(\mu^+)}{\sigma(\mu^-)}\leq 1.2\times 10^{-3}\,,\ \text{whence } d\leq\begin{cases}0.2 & (f=+1)\,,\\ 0.37 & (f=-1)\,.\end{cases} \tag{3.9}$$

We see that in general from the one-wrong-sign rate we cannot expect bounds as good as those found from mixing measurements (if these are well defined, see remarks above).

Effects of associated charm production have been assumed to be negligible.

Next, the contribution of charm change (3.4) to the *dimuon production* is considered. For charged (see GIM model) charm production we have

$$\frac{d\sigma^{\Delta Q,\nu}}{dx\,dy}=\frac{G^2ME}{\pi}\left[(u(x)+d(x))\sin^2\theta_C+2s\cos^2\theta_C\right],$$
$$\frac{d\sigma^{\Delta Q,\bar{\nu}}}{dx\,dy}=\frac{G^2ME}{\pi}\,2\bar{s}\cos^2\theta_C\,, \tag{3.10}$$

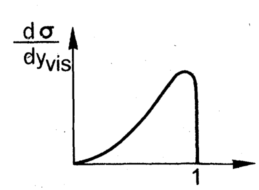

FIG. 2. *Schematic* visible $y_{\rm vis}$ distribution caused by "faked" μ^+ from charm change in ν [in addition to the usual $(1-y)^2$]; for numbers see the text.

coming from the diagrams in Fig. 1(b). When a D^0 is produced, it can undergo mixing and produce like-sign muons. Again formula (2.6) is used:

$$\left.\frac{\sigma(\mu^-\mu^-)}{\sigma(\mu^-\mu^+)}\right|_\nu = R_M \times F(D^0)\,,$$
$$\left.\frac{\sigma(\mu^+\mu^+)}{\sigma(\mu^+\mu^-)}\right|_{\bar\nu} = R_M \times F(\bar D^0)\,, \tag{3.11}$$

always having multiplied the fraction of real D^0 ($\bar D^0$) produced at the vertex: $F(D^0)\,(F(\bar D^0))$. The data of CDHS[32]

$$\left.\frac{\sigma(\mu^+\mu^+)}{\sigma(\mu^+\mu^-)}\right|_{\bar\nu} = \frac{1.5\pm 2.7}{40}\,,\quad \left.\frac{\sigma(\mu^-\mu^-)}{\sigma(\mu^-\mu^+)}\right|_\nu = \frac{9\pm 12}{256}$$

are consistent with 5% and therefore

$$\sqrt{n}\,\frac{d}{\sqrt{2}} \leq 5.2\times 10^{-4}\,, \tag{3.12}$$

which is [assuming $F(D^0)=0.1$] of the same order as the result obtained from e^+e^-, see Sec. II, Eq. (2.7). Note the uncertainty of n (Sec. II) and $F(D^0)$. The above-mentioned value for $F(D^0)$ may be guessed from meson production off nuclei: The system $\{c,\bar q\}$, $\bar q=\bar u$, $\bar d$, $\bar s$ can be in the states D^0, D^{*0}, D^+, D^{*+}, F^+ and F^{*+}. Counting all the mesons together with their spin yields 12. From Ref. 4 one knows that, e.g., $D^{*0}\to D^0\pi^0$ with 85–90% and also $D^{*+}\to D^0\pi^+$ (seen, but branching ratio as yet unknown). Therefore

$$F(D^0)=F(\bar D^0)\approx 0.1 \tag{3.13}$$

might be a guess.

Finally look at *trimuons*: When a FCNC like (3.4) is present, the decays

(A) $D^0\to \pi^0\mu^+\mu^-$,

(B) $D^0\to \mu^+\mu^-$,

(C) $D^{*0}\to \mu^+\mu^-$

will contribute to the trimuon rate. Because of the low statistics in trimuons observed up to now it will be difficult to filter out process A. By means of the FCNC form factor, estimated in Appendix C it can be evaluated. Processes B and C could possibly be observed in the invariant-mass spectrum of $\mu^+\mu^-$ pairs. However, process C will not contribute as long as strong decays like $D^{*0}\to D^0\pi^0$ (85%–95 % in the D^{*0} case) can still take place. Only for very-high-mass vector mesons, where possibly the vector-pseudoscalar mass difference becomes small, only em $M1$ transitions could take place.[9] Also weak decays could become concurrent.[18] For the moment we are left with process B: $D^0\to\mu^+\mu^-$. The formula is given in Appendix A, where the pion form factor has been inserted (which is an underestimate, as can be seen from recent considerations on pseudoscalar form factors).[16] Together with the neutral leptonic current $\frac{1}{2}\bar\psi\gamma_\mu(C_V-C_A\gamma_5)\psi$ [the Weinberg-Salam predictions are $C_V=-\frac{1}{2}+2\sin^2\theta_W$, $C_A=-\frac{1}{2}$ (Ref. 10)] one obtains in terms of (3.4)

$$\frac{\Gamma(D^0 \xrightarrow{N\Delta C} \mu^+\mu^-)}{\Gamma(D^0\to K^-\pi^+)} \approx 0.25\times 10^{-2} C_A{}^2 d^2 f^2\,. \tag{3.14}$$

For $D^0\to K^-\pi^+$, see Ref. 23 (multipion decays are included). The production of $\mu^+\mu^-$ pairs at the invariant mass m_{D^0} is then

$$\frac{\sigma(\mu+(\mu^+\mu^-)_{M_{D^0}})}{\sigma(\mu)} = \frac{\sigma_{\text{charm}}^{\text{charged current}}}{\delta(\mu)} F(D^0)\,\frac{\Gamma(D^0 \xrightarrow{N\Delta C}\mu^+\mu^-)}{\Gamma(D^0\to X)}$$
$$\approx 0.3\times 10^{-3} F(D^0) C_A{}^2 d^2\,. \tag{3.15}$$

If for instance in 10^5 single-muon events, N were found, compatible with a D^0 decay, one would get the upper limit

$$C_A d \leq 0.2\left(\frac{N}{F(D^0)}\right)^{1/2}.$$

Note that N is limited by the total trimuon rate (three events in 10^5, Kleinknecht CDHS[32]). This limit will be improved by better statistics, subtraction of process A, and a better understanding of normal trimuon production.

To summarize, in deep-inelastic neutrino scattering, dimuon production sets the best upper limit on the effective coupling of a charm-changing NC because D^0-$\bar D^0$ mixing is the main effect. As it has been argued in paragraph 1 of this section, for higher vector mesons the signatures of X^0-$\bar X^0$ mixing might become ambiguous, so that other tests can gain in interest: wrong-sign muons for flavor change involving nucleon quarks and the invariant-mass spectrum of $\mu^+\mu^-$ pairs in trimuons (if the mass of the corresponding meson is known) as well as the test discussed in the following paragraph.

C. Direct coupling of vector particles to the NC and signatures of flavor change

In e^+e^- flavor-changing NC should give rise to processes[9] like in Fig. 3(a). The case where $X^0=\{q,\bar q'\}$ or $\bar X^0$ are in a resonant state, Fig. 3(b). will dominate. In principle, there are as intermediate states J^{PC}: 1^{--}(vector), 1^{++} and 1^{+-} (axial vectors), 0^{++} (scalar), and 0^{-+} (pseudoscalar). But, first of all, scalars will be difficult to observe in e^+e^- because, in the Breith-Wigner cross section formula, the leptonic width $\Gamma_{ee}=\Gamma(X^0\to e^+e^-)$ will appear, which, for scalars, is suppressed by helicity conservation ($\sim m_e{}^2$). Second, axial vectors are rather broad. Therefore we are left with vector mesons X^{*0}. The cross section at resonance

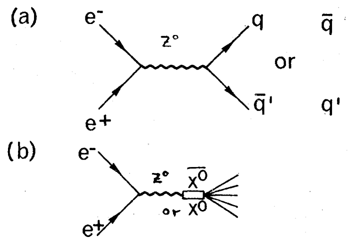

FIG. 3. (a) Flavor-nonconserving graph in e^+e^-, (b) $\bar{q},q'$ and $\bar{q}',q$ in resonant states.

is

$$\sigma(e^+e^- \to X^{*0} \to \text{anything})\,\big|_{\sqrt{s}=M_{X^{*0}}} = \frac{12\pi\Gamma_{ee}(X^{*0}\to e^+e^-)}{M_{X^{*0}}{}^2\,\Gamma_V}\,, \quad (3.16)$$

where Γ_V is the total width of the X^{*0} vector meson is given in Appendix A. Taken into account the probably small total widths of X^{*0} vector mesons corrected for the experimental resolution ($\Delta\sqrt{s}$) and multiplied by two (because X^0 and $\bar{X}^0$ are produced) the ratio of the integrated signal versus the total $e^+e^-\to\mu^+\mu^-$ cross section is[9]

$$\sigma_{\text{peak}}(e^+e^-\to Z^0\to X^{*0}\to\text{anything}) = \sigma(e^+e^-\to\mu^+\mu^-)\frac{M_V}{\Gamma_V}\left(\frac{1}{M_{\text{proton}}}\right)^4 16\pi f_V{}^2\times 2.7\times 10^{-9} \quad (3.17)$$

and

$$\frac{\int\sigma(e^+e^-\to Z^0\to X^{*0}\to\text{anything})\,d\sqrt{s}}{\sigma(e^+e^-\to\mu^+\mu^-)\Delta\sqrt{s}} = \frac{\pi}{2}\frac{\sigma_{\text{peak}}}{\sigma\mu^+\mu^-}\frac{\Gamma_V}{\Delta\sqrt{s}}\,. \quad (3.18)$$

f_V is the vector form factor defined in Appendix A; estimates are found in Appendices B1 and B2. Some numerical values for the signal size[9] are given in Table I.

TABLE I. Sizes of signals: for neutral-vector-meson production in e^+e^- versus μ-pair production, in the case of maximal [$d_V=1$ in Eq. (3.1)] flavor change in the vector part of the neutral current.

Mass of the vector meson (GeV)	Assumed experimental resolution (MeV)	Signal size B^{FC} (%)
2 (D^{*0})	5	1.4×10^{-3}
6	7	0.25
8	14	0.5
15.5	27	3
30	48	5

The signature would be at fixed energy $\sqrt{s}=m_V$: (1) An apparent strangeness violation: production of single K's (D's) (2) Monochromatic γ's or π^0's (possibly only γ's) coming from the decay into a pseudoscalar with emission energies of 140–10 MeV depending on m_V. The upper limits on $(d^2/\sqrt{2})(G/\sqrt{2})$ for the strength of the off-diagonal NC would be, for the resolutions assumed in Table I, when signals of (0.5–5% are absent, of the order $\frac{1}{2}d^2\leq 0.7$–0.3. But in principle they could be made much smaller by limiting $\Delta\sqrt{s}$, because the vector-meson width Γ_V is very small.[9] Theoretical estimates in the literature for the D^{*0} (2 GeV) range from 100 eV to 1 keV,[34] coming from the small mass difference[4]:

$$m(D^{*0})-m(D^0)=141\pm5\text{ MeV}\,. \quad (3.19)$$

Therefore, in principle, the resonance peak is very high, the factor $\Delta\sqrt{s}/\Gamma_V$ being of the order $>10^3$. But it seems as yet unrealistic to assume that $\Delta\sqrt{s}$ can be pushed down to values of 0.1–1 keV, at beam energies of several GeV. Nevertheless, in future $\bar{p}p$ facilities very high beam resolutions are expected. The expected small mass differences[9] suggest the application of the nonrelativistic quark model.[35] Up to now, these calculations gave only rather unsatisfactory results for the decays into pseudoscalars and a pion or a gamma, a reason being perhaps the relativistic motion of the emitted particle. Possibly now the time has come, with high-mass mesons and low-energy emission, to really test the nonrelativistic quark model. The predictions are

$$\frac{g_{VP\pi}}{(2\pi)^3}=\tfrac{3}{5}g_{\pi N}\approx 8.2\,,$$

$$g_{VP\gamma}=\begin{cases}\frac{32}{27}(2\mu)^2 & \text{for quarks with charge } \frac{2}{3}\,,\\ \frac{8}{27}(2\mu)^2 & \text{for quarks with charge } -\frac{1}{3}\,,\end{cases} \quad (3.20)$$

with $g_{\pi n}$ = pion-nucleon coupling constant and $\mu=(2.79/2M_{\text{proton}})(\alpha 4\pi)^{1/2}$ the proton magnetic, moment. The decay rates are given by

$$\Gamma(V\to P\pi)=\frac{g_{VP\pi}{}^2}{12\pi M_V}\left(\frac{(M_V{}^2+m_P{}^2-m_\pi{}^2)^2}{4M_V{}^2}-m_P{}^2\right)^{1/2},$$

$$\Gamma(V\to P\gamma)=\frac{g_{VP\gamma}{}^2}{12\pi}\left(\frac{M_V{}^2-m_P{}^2}{2M_V}\right)^{3/2}.$$

To summarize, the widths of the X^0-like vector mesons are expected to be very small. Consequently the experimental resolution $\Delta\sqrt{s}$ is the crucial parameter on which the upper limits (if the signal is absent) for the effective coupling of FCNC will depend.

In conclusion, I make some remarks on bremsstrahlung background and inclusive γ-ray spec-

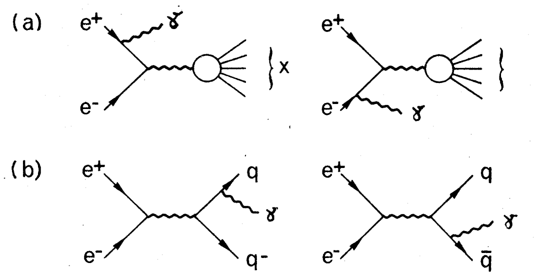

FIG. 4. Bremsstrahlung corrections to the proposed test of flavor conservation (by monochromatic γ's in e^+e^-).

troscopy. For the decay $V\to P\gamma$ one has two important background sources: (1) bremsstrahlung [Fig. 4(a)], (2) π^0, η decay into $\gamma\gamma$. Bremsstrahlung can be expressed by the total cross section of $e^+e^-\to$ hadrons times a factor[36]:

$$F=\frac{\alpha}{\pi}\ln\frac{k_2}{k_1}\left\{\frac{E}{|\vec{p}|}\ln\frac{\left(\frac{E}{|\vec{p}|}+\cos\theta_0\right)^2}{\left(\frac{E}{|\vec{p}|}-\cos\theta_0\right)^2}-\frac{m_e^2}{E^2}\frac{\cos\theta_0}{1-(|\vec{p}|/E)\cos\theta_0}\right\}. \qquad (3.21)$$

Here it has been integrated over the γ energy between k_1 and k_2, over the azimuth (2π), and over the polar angle from $\pi-\theta_0$ to $\pi+\theta_0$, corresponding to a cone-like blind region in both directions of the beam axis (in e^+e^- the beam pipe for instance). E and $\vec{p}$ are the energy and the spatial momentum of the electron or positron, m_e the electron mass. Some examples ($E=3.75$ GeV $\hat{=}\sqrt{s}=7.5$ GeV) for the factor F are given in Table II.

We see that when the forward/backward direction is avoided, the effect at $\sqrt{s}=7$ GeV is reducible to ~0.1%. In Ref. 9 it is shown that the signal itself increases as $M_V{}^4$, therefore improving the situation at higher masses. The emission of γ's from the hadronic blob is expected to be much smaller, but only if there is no other resonance, because then one can argue in terms of bremsstrahlung off quarks (e.g. Ref. 37), with the diagrams in figure 4(b), which are suppressed by the larger mass of the quark.

TABLE II. Background from bremsstrahlung to the γ emission of the flavor-change-produced neutral vector meson in e^+e^- ($\sqrt{s}=7.5$ GeV) in percent of the (nonresonant) $\sigma_{\rm tot}$, γ energies between k_1 and k_2, with different conelike regions (θ_0 in radians) the forward/backward direction excluded.

k_1	k_2 (MeV)	Whole solid angle: $\theta_0=0$	$\theta_0=0.01$	$\theta_0=0.05$	$\theta_0=0.1$
56	60	1%	0.17%	0.12%	0.1%
54	58	1%	0.18%	0.12%	0.1%
45	49	1%	0.21%	0.16%	0.12%

The background from produced π^0's and η's is much more serious. It is always possible that only one of the γ's is detected. But here some experimental progress has been reported. It turned out to be possible to observe directly[38] monochromatic γ's from ψ' decays with branching fractions of 2–10%, the minimal detection energy being about 50 MeV. In connection with this, it seems possible to account for the expected π^0 production by means of a Monte Carlo program based on a so-called "all pion-invariant-phase-space model"[38,39] and therefore the expected gamma background.

In conclusion, signatures of vector mesons produced by an off-diagonal neutral current may provide an important test for flavor-changing neutral couplings. Their strength is a crucial criterion for determining the structure of weak and electromagnetic interactions.

ACKNOWLEDGMENTS

Part of this work has been completed at the Laboratoire de Physique Théorique et Hautes Energies, Paris VI et VII (Jussieu), and I would like to thank Professor Michel Gourdin for welcoming me to his laboratory and for supporting my work. I am indebted to Dr. Jean Kaplan for his friendly criticism and keen discussions. For their helpful advice and discussions many thanks to Professor Xuan Yem Pham and Dr. Guy Bonneau, and from the Institute für Theoretische Physik in Karlsruhe (Germany) to Dr. Richard Fischer, as well to Patricia Meralli, Annette Holtkamp, and Rakshanda Faizi for technical support. This work was supported by the Studienstiftung des deutschen Volkes, Bonn–Bad Godesberg (Germany), and a French Government Scolarship.

APPENDIX A: NEUTRAL LEPTONIC DECAYS OF D^0 AND D^{*0}

By means of the leptonic neutral current $\frac{1}{2}\bar{\psi}\gamma_\mu(C_V-C_A\gamma_5)\psi$ and

$$\langle 0\,|j_\mu^{N,\Delta C}(0)\,|D^0\rangle=i(2\pi)^{-3/2}\frac{d}{2}f_{D^0}(p^2)P_\mu\,,$$

the D^0 partial width becomes

$$\Gamma(D^0\xrightarrow{N\Delta C}\mu^+\mu^-)$$

$$=\left(\frac{G}{\sqrt{2}}\right)^2\frac{d^2}{4}\frac{f_D{}^2M}{4\pi}\,2C_A{}^2m_\mu{}^2\left(1-\frac{4m_\mu{}^2}{M^2}\right)^{1/2},$$

and with the vector form factor

$$\langle 0\,|j_\mu^{N,\Delta C}(0)\,|D^{*0}\rangle = (2\pi)^{-3/2}\frac{d}{2} f_V(p^2)\epsilon_\mu^\lambda(p)$$

for D^{*0} one has

$$\Gamma(D^{*0} \xrightarrow{N\Delta C} \mu^+\mu^-)$$

$$=\left(\frac{G}{\sqrt{2}}\right)^2 \frac{d^2}{4}\frac{f_V M}{12\pi}\left(1-\frac{4m_\mu{}^2}{M^2}\right)^{1/2}$$

$$\times\left[(C_V{}^2+C_A{}^2)\left(1-\frac{m_\mu{}^2}{M^2}\right)+(C_V{}^2-C_A{}^2)\frac{3m_\mu{}^2}{M^2}\right]$$

[d is the coupling strength defined in Sec. II, Eq. (2.1)].

APPENDIX B: ESTIMATES ON THE FCNC VECTOR FORM FACTOR

(1) From $D^{*0}\to D^0\pi^0$ (50%[4]) follows by charge normalization and from vector-meson dominance, together with the usual smoothness assumption[18]:

$$f_V(0) = \frac{M^2}{g_{D^*D\pi}}(f_{D^0D^*\pi} + d_{D^0D^*\pi})$$

$g_{D^*D\pi}$ is the strong coupling constant usually determined from the width (3.20); $(f_{DD^*\pi})$ $d_{DD^*\pi}$ are the (anti-) symmetric structure constants of the symmetry. Corrections are to be expected, because the derivation holds only in the limit of exact symmetry ($m_D = m_\pi$).

(2) f_V can be related to the em form factor: Identification of the current $j_\mu^{N\Delta C}$ with W spin, e.g., in SU(4) as a member of the 15-plet, leads to

$$\langle 0\,|j_\mu^{N\Delta C}(0)\,|D^{*0}\rangle = \tfrac{1}{2} d\langle 0\,|\cdots + \bar{c}\gamma_\mu c\,|\psi\rangle\,,$$

relating f_V to the photon–vector-meson coupling of the nearest (in mass) resonance. From the best fit for this quantity[40] one obtains

$$16\pi f_V{}^2 = M_V{}^4\, 2.8\left(\frac{\text{GeV}}{M_V}\right)^{1+1/8}.$$

APPENDIX C: FLAVOR-CHANGING SEMILEPTONIC DECAYS

For the matrix element $\langle 0|j_\mu^{N\Delta C}(0)|D^0\pi^0\rangle$, which is related to $\langle D^0|j_\mu^{N\Delta C}|\pi^0\rangle$, and important for flavor-changing semileptonic decays, vector dominance predicts

$$\langle 0\,|j_\mu^{N\Delta C}(0)\,|D^0\pi^0\rangle = \frac{f_V(p^2)}{(2\pi)^3}\frac{g_{D^*D\pi}}{p^2 - M_{D^*}{}^2}$$

$$\times\, i\left[(p_\pi - p_0)_\mu - \frac{m_\pi{}^2 - m_D{}^2}{M_{D^*}{}^2}(p_\pi + p_D)_\mu\right].$$

---

*Formerly at the Institut für Theoretische Physik der Universität Karlsruhe, Physikhochhaus, Unitversität, D-7500 Karlsruhe, Germany.

[1]S. L. Glashow, J. Iliopoulos, L. Maiani, Phys. Rev. D 2, 1285 (1970).

[2]S. Weinberg, Phys. Rev. Lett. 19, 1264 (1967).

[3]A. Salam, in *Elementary Particle Physics*, edited by N. Svartholm (Almqvist and Wilksells, Stockholm, 1968).

[4]Particle Data Group, Rev. Mod. Phys. 48, S1 (1976); Phys. Lett. 68B, 1 (1977).

[5]J. Bjorken, in *Weak Interactions at High Energy and the Production of New Particles*, proceedings of the SLAC Summer Institute on Particle Physics, 1976, edited by M. Zipf (SLAC, Stanford, 1977), p. 1.

[6]E. Paschos, Phys. Rev. Lett. 39, 858 (1977).

[7]S. W. Herb *et al.*, Phys. Rev. Lett. 39, 252 (1977); W. R. Innes *et al.*, *ibid.* 39, 1240 (1977).

[8]M. Kobayashi and M. Maskawa, Prog. Theor. Phys. 49, 652 (1973).

[9]H. Genz and M. Gorn, Nucl. Phys. B140, 327 (1978).

[10]E. S. Abers and B. W. Lee, Phys. Rep., 9C, 1 (1973).

[11]E. Paschos, Phys. Rev. D 15, 1966 (1977).

[12]M. Gorn, Karlsruhe report, 1977 (unpublished).

[13]R. Mohapatra, J. Rao, and M. Marshak, Phys. Rev. Lett. 20, 1081 (1968); T. Appelquist, J. Bjorken, and M. Chanowitz, Phys. Rev. D 7, 2225 (1973); G. Preparata, Phys. Lett. B40, 253 (1972).

[14]L. Wolfenstein, Phys. Rev. Lett. 13, 562 (1964).

[15]J. Ellis, M. Gaillard, and D. Nanopoulos, Nucl. Phys. B109, 213 (1976).

[16]J. Ellis, M. Gaillard, D. V. Nanopoulos, and S. Rudaz, Nucl. Phys. B131, 285 (1977).

[17]Y. Achiman and T. Walsh, Phys. Lett. 66B, 174 (1977).

[18]M. Gorn, Ph.D. thesis, Karlsruhe University, 1978 (unpublished).

[19]G. Ecker, Phys. Lett. 72B, 450 (1978).

[20]S. Glashow and S. Weinberg, Phys. Rev. D 15, 1958 (1977).

[21]M. Chanowitz, J. Ellis, and M. Gaillard, Nucl. Phys. B128, 506 (1977).

[22]M. L. Goldberger and K. M. Watson, *Collision Theory* (Wiley, New York, 1967).

[23]M. Gaillard, B. W. Lee, and J. Rosner, Rev. Mod. Phys. 47, 277 (1975).

[24]J. Ellis, M. Gaillard, and D. Nanopoulos, Nucl. Phys. B100, 313 (1975).

[25]L. B. Okun, V. I. Zakharov, and B. M. Pontecorvo, Lett. Nuovo Cimento 13, 218 (1975); A. Pais and S. B. Treiman, Phys. Rev. D 12, 2744 (1975).

[26]G. Goldhaber *et al.*, Phys. Rev. Lett. 37, 255 (1976).

[27]T. Feldman *et al.*, Phys. Rev. Lett. 38, 1313 (1976).

[28]B. W. Lee and S. Weinberg, Phys. Rev. Lett. 38, 1237 (1977); P. Langacker, G. Segrè, and M. Golshani, Phys. Rev. D 17, 1402 (1978); R. Fischer, Acta Phys. Austriaca 50, 13 (1978).

[29]A. Ali and Z. Z. Aydin, Nucl. Phys. B148, 165 (1979).
[30]J. D. Bjorken and E. Paschos, Phys. Rev. D 1, 3151 (1970). R. P. Feynman, *Photon-Hadron Interactions* (Benjamin, New York, 1972).
[31]L. M. Seghal, P. Zerwas, Nucl. Phys. B108, 483 (1976).
[32]K. Kleinknecht, in *Proceedings of the International Symposium on Lepton and Photon Interactions at High Energy*, edited by F. Gutbrod (DESY, Hamburg, 1977).
[33]A. Benvenuti *et al.*, Phys. Rev. Lett. 36, 1478 (1976); 37, 189 (1976); B. Barish *et al.*, Phys. Rev. Lett. 38, 314 (1977).
[34]A. Bohm and R. B. Teese, Phys. Rev. Lett. 38, 629 (1977).
[35]R. Van Royen and V. Weisskopf, Nuovo Cimento 50A, 4781 (1967).
[36]S. Gasiorowicz, *Elementary Particle Physics* (Wiley, New York, 1966).
[37]R. Farrar and B. Joffe, Phys. Lett. 71B, 118 (1977).
[38]C. Biddik *et al.*, Phys. Rev. Lett. 38, 1324 (1977).
[39]G. Feldman and M. Perl, Phys. Rep. 19C, 235 (1975); 33C, 285 (1977).
[40]R. Barnett, in *Proceedings of the European Conference on Particle Physics*, edited by L. Jenik and I. Montvag, (CRIP, Budapest, 1978).