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Weak nonleptonic decays of charmed hadrons in models with right-handed currents

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We analyze the two-body weak nonleptonic decays of "charmed" pseudoscalar mesons to normal hadrons, in models with the most general structure for the right-handed currents. The structure of the left-handed part of the current is chosen in accordance with the Glashow-Iliopoulos-Maiani mechanism. Implications for the SPEAR charm search are analyzed.

I. INTRODUCTION

The need for the introduction of one more degree of freedom[1] in hadronic physics was first suggested in the study of weak interactions in connection with the suppression of $\Delta S=1$ neutral currents. Although many interpretations have been advanced for the newly discovered resonances,[2] one of the most common beliefs is that they confirm this need for a new quantum number (charm) in the description of strong interactions. A direct confirmation of the charm scheme would be the discovery of particles carrying net charm.[3] A detailed study of the properties of these particles has been done, based on the standard model of weak and electromagnetic interactions in which the four-quark model is incorporated into the Weinberg-Salam gauge theory. As far as the weak interactions are concerned, of particular interest is the question of how they break the new symmetry of the strong interactions. The study of the decay characteristics of the charmed hadrons seems to provide the crucial test for the different theoretical possibilities. Since the new quantum number is supposed to be broken only by the weak interactions, the least massive of these charmed hadrons are expected to decay only weakly. Most of the previous analyses[4-6] of these decays have been done assuming a particular structure[3] for the weak interactions of these particles, namely, the one indicated by the Glashow-Iliopoulos-Maiani (GIM) charmed current. However, as has been emphasized recently,[7] the possibility of an SU(4) symmetry for the strong interactions (with its implied mass spectrum) and the problem of its symmetry breakdown by the weak interactions are two independent questions. One alternative to the GIM current consists of adding a right-handed charm-changing current. This was suggested by one of us some time ago in connection with a unified gauge model of weak and electromagnetic interactions with CP violation.[8,9]

These currents, which have been extensively discussed in recent literature,[10] are generally of the type

$$K_\mu=\bar{x}_R\gamma_\mu\mathfrak{N}_R+\bar{y}_R\gamma_\mu\lambda_R+\bar{\mathcal{P}}_R\gamma_\mu z_R+\cdots\ , \qquad (1)$$

where $\mathcal{P}_R, \mathfrak{N}_R, \lambda_R$ are the right-handed $\mathcal{P}, \mathfrak{N}, \lambda$ quarks, and x_R, y_R, z_R stand for right-handed quarks, each one of them carrying a new quantum number conserved by the strong interactions. The dots stand for other currents, not containing any of the quarks $\mathcal{P}, \mathfrak{N}, \lambda$. Either x or y will coincide with the charmed quark χ, used to implement the GIM mechanism. In the case where x coincides with the charmed quark χ, an elegant explanation for the observed enhancement of weak nonleptonic decays (with $\Delta C=0$, $\Delta I=\frac{1}{2}$) seems to emerge.[9] Whether such a mechanism is viable from the experimental point of view has recently been questioned by Golowich and Holstein.[11] Our considerations here will not be dependent on this argument and will focus on the weak nonleptonic decays of the pseudoscalar bound states of each of the x, y, z quarks with the three SU(3) quarks $\mathcal{P}, \mathfrak{N}, \lambda$. We call these bound states $D_x^0(x\bar{\mathcal{P}})$, $D_x^+(x\bar{\mathfrak{N}})$, $F_x^+(x\bar{\lambda})$, and similarly for y, z. It has in fact been recently suggested that the x, y quarks may indeed be degenerate.[12] It is clear that either D_x^0, D_x^+, F_x^+ or D_y^0, D_y^+, F_y^+ will coincide with the D's and F's of Ref. 4. As far as the decays of the charmed pseudoscalar mesons are concerned, one has two classes of models corresponding to the choices (a) $x\equiv\chi$ and (b) $y\equiv\chi$. We will deal separately with the two cases, pointing out in detail their implications to the nonleptonic decays of the charmed mesons D^0_α, D^+_α, F^+_α (where $\alpha = x$ or y) and D^0_z, D^-_z, F^0_z.

The plan of this paper is as follows. In Sec. II we analyze the properties of the nonleptonic Hamiltonian. In Sec. III we make a detailed analysis of various decay channels of the charmed pseudoscalar mesons. We will assume exact SU(3) symmetry throughout. Although it remains to be seen if this is a reliable approximation, we assume that possible corrections due to SU(3) symmetry breaking will be small enough not to alter substantially our general conclusions. In Sec. IV we discuss our results in view of the recent data from the SPEAR charm search. Boyarski *et al.*[3] have searched for inclusive production of charmed mesons by looking for narrow peaks in various invariant-mass distributions in which charmed mesons might be found. Using their results, one can deduce upper limits for the branching ratio of different decay modes of D^0, D^+, F^+. The implications of these upper limits have been analyzed[13] in the context of the "orthodox" theory of weak nonleptonic decays of charmed particles. In particular, it has been pointed out that the experimental upper bound $B(D^0 \to K^-\pi^+) \leq 2.9\%$ implies $B(D^0 \to 2$ pseudoscalars) $\leq 8.6\%$. This in turn would mean that about 90% of the nonleptonic decays of D^0 are to three or more particles, in apparent contradiction with the observed charged-particle multiplicity. We will show how this difficulty can be avoided in the modified charmed current models.

II. PROPERTIES OF THE NONLEPTONIC HAMILTONIAN

We will assume that the nonleptonic weak Hamiltonian has the usual current-current form

$$\mathcal{H}_W = \frac{G_F}{\sqrt{2}} (H_\mu H^{\dagger\mu} + H^\dagger_\mu H^\mu). \tag{2}$$

The hadronic current H_μ contains both the usual left-handed current J_μ and the new right-handed current K_μ:

$$\begin{aligned} H_\mu &= J_\mu + K_\mu, \\ J_\mu &= \bar{\mathcal{P}}_L \gamma_\mu \mathfrak{N}(\theta_C)_L + \bar{\chi}_L \gamma_\mu \lambda(\theta_C)_L, \end{aligned} \tag{3}$$

where K_μ is given by (1) and

$$\begin{pmatrix} \mathfrak{N}(\theta_C) \\ \lambda(\theta_C) \end{pmatrix} = \begin{pmatrix} \cos\theta_C & \sin\theta_C \\ -\sin\theta_C & \cos\theta_C \end{pmatrix} \begin{pmatrix} \mathfrak{N} \\ \lambda \end{pmatrix},$$

with θ_C the Cabibbo angle, and $\gamma_{L,R} = (1 \mp \gamma_5)$. Recall that the χ denotes the charm quark used to implement the GIM mechanism.

This Hamiltonian decomposes with respect to SU(4) as

$$\mathcal{H}_W = \underline{1} + \underline{15}_S + \underline{20} + \underline{84}.$$

Note that with the modified current the $\underline{15}$ is now present in $\mathcal{H}_W$, contrary to what happens in the standard GIM model.

We will analyze now the part of $\mathcal{H}_W$ contributing to the charm-changing nonleptonic decays. We will designate collectively by "charm" the new quantum numbers of the x, y, z quarks and distinguish among them by referring to C_x, C_y, C_z. We will analyze separately the two classes of models mentioned before.

A. Class (a) of models with $x \equiv \chi$

We recall that in this class of models, the $\Delta C_\chi = 0$, $\Delta S = \pm 1$ Hamiltonian is given effectively by

$$\mathcal{H}^S_W = \frac{G}{\sqrt{2}} \{ \cos\theta_C (\bar{\lambda}\chi)_L (\bar{\chi}\mathfrak{N})_R + \cos\theta_C \sin\theta_C [(\bar{\lambda}\mathcal{P})_L (\bar{\mathcal{P}}\mathfrak{N})_L - (\bar{\lambda}\chi)_L (\bar{\chi}\mathfrak{N})_L] \}, \tag{4}$$

where for convenience we have omitted Lorentz indices, and L, R stand for left- and right-handed currents, respectively. The first term[9] contributes to $\Delta I = \frac{1}{2}$ transitions only and is enhanced with respect to the second term (which contributes both to $\Delta I = \frac{1}{2}$ and $\Delta I = \frac{3}{2}$ transitions) by the Cabibbo factor $(\sin\theta)^{-1}$. The first term is further enhanced by the large mass of the charmed quark χ.

For the study of the decays of the SU(3) triplet of charmed pseudoscalar mesons D^+_χ, D^0_χ, F^+_χ, we are interested in the $\Delta C_\chi = 1$ part of the Hamiltonian which is

$$\begin{aligned} \mathcal{H}^{C_\chi}_W &= \mathcal{H}^{C_\chi}_{W(LL)} + \mathcal{H}^{C_\chi}_{W(RL)}, \\ \mathcal{H}^{C_\chi}_{W(LL)} &= \cos^2\theta_C (\bar{\chi}\lambda)_L (\bar{\mathfrak{N}}\mathcal{P})_L - \sin^2\theta_C (\bar{\chi}\mathfrak{N})_L (\bar{\lambda}\mathcal{P})_L \\ &\quad + \cos\theta_C \sin\theta_C [(\bar{\chi}\lambda)_L (\bar{\lambda}\mathcal{P})_L - (\bar{\chi}\mathfrak{N})_L (\bar{\mathfrak{N}}\mathcal{P})_L], \end{aligned} \tag{5}$$

$$\mathcal{H}^{C_\chi}_{W(RL)} = \cos\theta_C (\bar{\chi}\mathfrak{N})_R (\bar{\mathfrak{N}}\mathcal{P})_L + \sin\theta_C (\bar{\chi}\mathfrak{N})_R (\bar{\lambda}\mathcal{P})_L.$$

$\mathcal{H}^{C_\chi}_{W(LL)}$ is present in the standard gauge models with the GIM current. Its properties were fully discussed in Refs. 4, 5, and 6. We will direct our attention to the new term $\mathcal{H}^{C_\chi}_{W(RL)}$ which appears in the right-handed current models of class (a).

With respect to SU(3), $\mathcal{H}^C_W$ transforms as $\underline{8} \otimes \underline{3} = \underline{3}^* \otimes \underline{6} \otimes \underline{15}^*_M$. The states of these representations can be expressed in the basis of quark current-current products, as

$$[\underline{3}^*]_i = H^k_{ik},$$
$$[\underline{6}]^{kl} = \epsilon^{kij} H^l_{ij} + \epsilon^{lij} H^k_{ij}, \qquad (6)$$
$$[\underline{15}_M]^k_{ij} = H^k_{ij} + H^k_{ji} - \tfrac{1}{4}(\delta^k_i H^l_{jl} + \delta^k_j H^l_{il}),$$

where

$$H^k_{ij} = (q^k q_j)(\bar{\chi} q_i) - \tfrac{1}{3}\delta^k_i (q^l q_l)(\bar{\chi} q_j), \qquad (7)$$

with

$$(\chi, q_1, q_2, q_3) = (\chi, \mathcal{P}, \mathfrak{N}, \lambda).$$

Thus we find that $\mathcal{H}_W^{C_\chi}$ can be decomposed as

$$\mathcal{H}^{C_\chi}_{W(LL)} = \tfrac{1}{4}\{\cos^2\theta_C[\underline{6}]^{22} + \cos\theta_C \sin\theta_C[\underline{6}]^{23} + \sin^2\theta_C[\underline{6}]^{33}\}_{LL}$$
$$+\text{terms belonging to the } \underline{15}^*_M \text{ representation}, \qquad (8)$$

$$\mathcal{H}^{C_\chi}_{W(RL)} = \tfrac{1}{4}\{\cos\theta_C[\underline{6}]^{23} + \sin\theta_C[\underline{6}]^{33} + \tfrac{3}{2}\cos\theta_C[3^*]_1\}_{RL}$$
$$+\text{terms belonging to the } \underline{15}^*_M \text{ representation}. \qquad (9)$$

Terms in the SU(3) $\underline{15}^*_M$ representation will be discarded hereafter, since they come from the SU(4) $\underline{84}$. We recall that the $\underline{84}$ contains the SU(3) $\underline{27}$, which is known to be suppressed in strangeness-changing decays. Note that the $\underline{3}^*$ is absent in $\mathcal{H}^C_{W(LL)}$ but contributes to $\mathcal{H}^C_{W(RL)}$.

We analyze now the part of the Hamiltonian contributing to the decay of the y-type and z-type "charmed" mesons. We obtain

$$\mathcal{H}^{C_y}_{W(RL)} = (\bar{y}\lambda)_R(\bar{\mathfrak{N}}\mathcal{P})_L \cos\theta_C + (\bar{y}\lambda)_R(\bar{\lambda}\mathcal{P})_L \sin\theta_C,$$
$$\mathcal{H}^{C_z}_{W(RL)} = (\bar{z}\mathcal{P})_R(\bar{\mathcal{P}}\mathfrak{N})_L \cos\theta_C + (\bar{z}\mathcal{P})_R(\bar{\mathcal{P}}\lambda)_L \sin\theta_C. \qquad (10)$$

These expressions can in turn be composed as

$$\mathcal{H}^{C_y}_{W(RL)} = \tfrac{1}{4}\{\cos\theta_C[\underline{6}]^{22} + \sin\theta_C[\underline{6}]^{23} + \tfrac{3}{2}\sin\theta_C[3^*]_1\}_{RL} + \cdots,$$
$$\mathcal{H}^{C_z}_{W(RL)} = \tfrac{1}{4}\{\cos\theta_C[\underline{6}]^{13} + \tfrac{3}{2}\cos\theta_C[\underline{3}^*]_2 - \sin\theta_C[\underline{6}]^{12} + \tfrac{3}{2}\sin\theta_C[\underline{3}^*]_3\}_{RL} + \cdots. \qquad (11)$$

B. Class (b) of models with $y \equiv \chi$

It is apparent that models of this type conform to the constraints derived in Ref. 11. The $\Delta S = \pm 1$ Hamiltonian is now given by

$$\mathcal{H}^S_W = \frac{G}{\sqrt{2}} \cos\theta_C \sin\theta_C[(\bar{\lambda}\chi)_R(\bar{\chi}\mathfrak{N})_L + (\bar{\lambda}\mathcal{P})_L(\bar{\mathcal{P}}\mathfrak{N})_L - (\bar{\lambda}\chi)_L(\bar{\chi}\mathfrak{N})_L] + \text{H.c.} \qquad (12)$$

We see that the (RL) term, which contributes only to the $\Delta I = \frac{1}{2}$ transitions, is also suppressed by $\sin\theta_C$. Thus one has to rely only on a dynamical enhancement (which, for example, may arise either due to renormalization-group arguments[14] or otherwise[15]) to explain the $\Delta I = \frac{1}{2}$ rule. Note that this right-handed charmed current may appear in a variety of gauge models, as for example in those of Ref. 9 (in some of these models one would have to make the interchange $\mathfrak{N}_R \leftrightarrow \lambda_R$). We are interested here in the charm-changing part of the Hamiltonian. As explained in Eq. (5), there is an LL and an RL part in the Hamiltonian, the (LL) term remains unchanged [see Eq. (5)], and the (RL) term is given by

$$\mathcal{H}^{C_\chi}_{W(RL)} = \tfrac{1}{4}\{\cos\theta_C[\underline{6}]^{22} + \sin\theta_C[\underline{6}]^{23} + \tfrac{3}{2}\sin\theta_C[3^*]_1\}_{RL}. \qquad (13)$$

Obviously, this coincides with (11). The $\Delta C_x = 1$ part of the Hamiltonian is now given by

$$\mathcal{H}^{C_x}_{W(RL)} = (\bar{x}\mathfrak{N})_R(\bar{\mathfrak{N}}\mathcal{P})_L \cos\theta_C + (\bar{x}\mathfrak{N})_R(\bar{\lambda}\mathcal{P})_L \sin\theta_C,$$

which can in turn be decomposed as

$$\mathcal{H}^{C_x}_{W(RL)} = \tfrac{1}{4}\{\cos\theta_C[\underline{6}]^{23} + \sin\theta_C[\underline{6}]^{33} + \tfrac{3}{2}\cos\theta_C[3^*]_1\}_{RL}. \qquad (14)$$

Note that $\mathcal{H}^{C_z}_{W(RL)}$ is the same in both classes of models.

III. NONLEPTONIC DECAYS

A. Class (a) of models with $x \equiv \chi$

We will first analyze the various decay modes of $D^+_\chi, D^0_\chi, F^+_\chi$. From (9), it is apparent that

(i) right-handed currents contribute to $C = 1$ transitions via both the $\underline{6}_{RL}$ and the $\underline{3}^*_{RL}$;

(ii) decay modes which were suppressed by a $\sin\theta_C$ factor in the standard GIM model can now occur through $[\underline{6}]^{23}_{RL}$ without Cabibbo suppression.

We present the tables with the relative rates (apart from phase-space factors) for the decays of D^+, D^0, F^+ into two pseudoscalars, two vectors, a pseudoscalar-vector pair, and a baryon-antibaryon pair. Many of these results can be easily obtained from those of Ref. 6 by taking into account (i) and (ii). However, for completeness we present all the tables. We will now point out some of their salient features.

1. *Decays of F^+_χ, D^+_χ.* It is interesting to note that for the decays of F^+, D^+ into two pseudoscalars, the relative rates can be expressed in terms of only two parameters $P, (Q+R)$ which will be defined in Table V. This is due to the fact that the contributions of $\underline{6}_{RL}$ and $\underline{3}_{RL}$ enter

TABLE I. Relative decay rates of F^+.

$F^+\to PP$		$F^+\to VV$		$F^+\to VV$		$F^+\to \bar{B}B$	
	$\cos^4\theta\,\lvert P\rvert^2\times$		$\cos^4\theta\,\lvert\alpha\rvert^2\times$		$\cos^4\theta\times$		$\cos^4\theta\times$
$\pi^+\eta'$	$\frac{4}{9}$	$\rho^+\omega$	$\frac{2}{3}$	$\pi^+\omega$	$\left\lvert-\frac{D}{\sqrt{2}}-\sqrt{2}\,S\right\rvert^2$	$\Sigma^+\bar{\Lambda}$	$\frac{1}{6}\lvert 3D+2S\rvert^2$
$K^+\bar{K}^0$	$\frac{1}{3}$	$K^{*+}\bar{K}^{*0}$	$\frac{1}{3}$			$\Sigma^+\bar{\Sigma}^0$	$\frac{1}{2}\lvert D-2A\rvert^2$
$\pi^+\eta$	$\frac{2}{9}$			$\pi^+\rho^0$	$\left\lvert\frac{D}{\sqrt{2}}-\sqrt{2}\,A\right\rvert^2$	$\Lambda\bar{\Sigma}^+$	$\frac{1}{6}\lvert 3D-2S\rvert^2$
	$\cos^2\theta\,\lvert Q+R\rvert^2\times$		$\cos^2\theta\,\lvert\beta+\gamma\rvert^2\times$	$\pi^+\phi$	$\lvert D\rvert^2$	$\Sigma^0\bar{\Sigma}^+$	$\frac{1}{2}\lvert D-2A\rvert^2$
$K^+\eta'$	$\frac{4}{9}$	$K^{*+}\phi$	$\frac{1}{3}$	$\eta\rho^+$	$\frac{1}{6}\lvert(3D-2S)\rvert^2$	$p\bar{n}$	$\lvert D-S+A\rvert^2$
π^+K^0	$\frac{1}{3}$	$K^{*0}\rho^+$	$\frac{1}{3}$	$\pi^0\rho^+$	$\frac{1}{2}\lvert(D-2A)\rvert^2$	$\Xi^0\bar{\Xi}^+$	$\lvert D+S+A\rvert^2$
π^0K^+	$\frac{1}{6}$	$K^{*+}\rho^0$	$\frac{1}{6}$	$K^+\bar{K}^{*0}$	$\lvert D-S+A\rvert^2$		$\cos^2\theta\times$
$K^+\eta$	$\frac{1}{18}$	$K^{*+}\omega$	$\frac{1}{6}$	$\bar{K}^0K^{*+}$	$(D+S+A)^2$	$p\bar{\Sigma}^0$	$\frac{1}{2}\lvert 2D'-S'-S''-A'-A''\rvert^2$
				$\eta'\rho^+$	$\frac{4}{3}\lvert S\rvert^2$		
					$\cos^2\theta\times$	$p\bar{\Lambda}$	$\frac{1}{6}\lvert 6D'-S'-S''-A'-A''\rvert^2$
				$K^+\omega$	$\frac{1}{2}\lvert 2D'-S'-S''-A'-A''\rvert^2$	$\Sigma^0\bar{\Xi}^+$	$\frac{1}{2}\lvert 2D'+S'+S''-A'-A''\rvert^2$
				$K^+\rho^0$	$\frac{1}{2}\lvert 2D'+S'+S''+A'+A''\rvert^2$		
				$K^+\phi$	$\lvert 2D'-S'-S''+A'+A''\rvert^2$	$\Lambda\bar{\Xi}^+$	$\frac{1}{6}\lvert 6D'+S'+S''+3A'+3A''\rvert^2$
				π^0K^{*+}	$\frac{1}{2}\lvert 2D'+S'+S''-A'-A''\rvert^2$		
				ηK^{*+}	$\frac{1}{6}\lvert 6D'+S'+S''+3A'+3A''\rvert^2$	$\Sigma^+\bar{\Xi}^0$	$\lvert-2D'-S'-S''+A'+A''\rvert^2$
				π^+K^{*0}	$\lvert-2D'-S'-S''+A'+A''\rvert^2$	$n\bar{\Sigma}^+$	$\lvert 2D'-S'-S''-A'-A''\rvert^2$
				$K^0\rho^+$	$\lvert 2D'-S'-S''-A'-A''\rvert^2$		
				$\eta'K^{*+}$	$\frac{4}{3}\lvert S'+S''\rvert^2$		

always in the same combination in F^+, D^+ decays. We will see that this is no longer true for D^0. It is easy to see that the decay characteristics of D^+ are drastically changed. Recall that one of the salient features of the GIM current is that it predicts (if one assumes sextet dominance) no enhanced decays of D^+ into two pseudoscalars. Based on this fact, it has been argued that for D^+, the semileptonic modes may compete favorably with the nonleptonic ones. It is apparent from Table II that this is no longer true for the new current, since there are decays of D^+ via the $\underline{6}_{RL}$ and $\underline{3}^*_{RL}$ with no Cabibbo suppression.

2. *Decays of* D^0_χ. For the decay of D^0 into two pseudoscalars, one can measure separately the contributions of $\underline{6}_{RL}$ and $\underline{3}^*_{RL}$. Note for example that $D^0\to K^0\bar{K}^0$ goes only via $\underline{3}^*_{RL}$ and observation of this mode would thus provide a direct measure of its strength.

Of particular interest for the current search for charm are the decays of D^0 into two charged particles. While in the GIM model only π^+K^- is enhanced, now the $\pi^+\pi^-$, K^+K^- decays can go unsuppressed.

We will now analyze the decays of "y" and "z" type charmed mesons. From expression (11) for $\mathcal{H}^{Cy}_{W(RL)}$, it is clear that the Cabibbo enhancement decays of F^+_y, D^+_y, D^0_y will coincide with those of the F^+, D^+, D^0 in the GIM scheme since the SU(3) structure of the Cabibbo enhanced part for both $\mathcal{H}_{W(LL)}$ and $\mathcal{H}^{Cy}_{W(RL)}$ has the same SU(3) structure. We, therefore, do not present a separate table for these.

The decays of the z-type charmed mesons are given in Table IV.

B. Class (b) of models with $y\equiv\chi$

1. *Decays of* F^+_χ, D^+_χ, D^0_χ. From (13) one sees that the Cabibbo enhanced part of $\mathcal{H}^{C\chi}_{W(RL)}$ has the same form as $\mathcal{H}^{C\chi}_{W(LL)}$ (recall that the LL part of $\mathcal{H}_W$ is the same in all models and coincides with that of the GIM scheme). The term in $\sin\theta_C$ is, however, modified by the presence of $[3^*]_{RL}$.

TABLE II. Relative decay rates of D^+.

$D^+ \to PP$		$D^+ \to VV$		$D^+ \to PV$		$D^+ \to \bar{B}B$	
	$\cos^2\theta\,\lvert(Q+R)+\sin\theta P\rvert^2 \times$		$\cos^2\theta\,\lvert(\beta+\gamma)+\sin\theta\,\alpha\rvert^2 \times$		$\cos^4\theta \times$		$\cos^4\theta \times$
$\pi^+\eta'$	$\frac{4}{9}$	$\rho^+\omega$	$\frac{2}{3}$	$\bar{K}^0\rho^+$	$\lvert 3D\rvert^2$	$\Sigma^0\bar{\Sigma}^+$	$\lvert 3D\rvert^2$
$K^+\bar{K}^0$	$\frac{1}{3}$	$K^{*+}\bar{K}^0$	$\frac{1}{3}$	$\pi^+\bar{K}^{*0}$	$\lvert 3D\rvert^2$	$\Sigma^+\bar{n}$	$\lvert 3D\rvert^2$
$\pi^+\eta$	$\frac{2}{9}$				$\cos^2\theta \times$		$\cos^2\theta \times$
	$\sin^2\theta\,\lvert Q+\sin\theta P\rvert^2 \times$		$\sin^2\theta\,\lvert\beta+\sin\theta\,\alpha\rvert^2 \times$	$\pi^+\phi$	$4\lvert D'\rvert^2$	$\Sigma^+\bar{\Sigma}^0$	$2\lvert D'+A'+A''\rvert^2$
$K^+\eta'$	$\frac{4}{9}$	$K^{*+}\phi$	$\frac{1}{3}$	$\pi^+\rho^0$	$2\lvert D'+A'+A''\rvert^2$	$\Sigma^+\bar{\Lambda}$	$\frac{2}{3}\lvert -3D'+S'+S''\rvert^2$
$K^0\pi^+$	$\frac{1}{3}$	$K^{*0}\rho^+$	$\frac{1}{3}$	$\pi^+\omega$	$2\lvert -D'+S'+S''\rvert^2$	$\Lambda\bar{\Sigma}^+$	$\frac{2}{3}\lvert 3D'+S'+S''\rvert^2$
$K^+\pi^0$	$\frac{1}{6}$	$K^{*+}\rho^0$	$\frac{1}{6}$	$\eta\rho^+$	$\frac{2}{3}\lvert 3D'+S'+S''\rvert^2$	$\Sigma^0\bar{\Sigma}^+$	$2\lvert D'+A'+A''\rvert^2$
$K^+\eta$	$\frac{1}{18}$	$K^{*+}\omega$	$\frac{1}{6}$	$\pi^0\rho^+$	$2\lvert D'+A'+A''\rvert^2$	$p\bar{n}$	$\lvert 2D'+S'+S''-A'-A''\rvert^2$
				$K^+\bar{K}^{*0}$	$\lvert 2D'+S'+S''-A'-A''\rvert^2$	$\Xi^0\bar{\Xi}^+$	$\lvert -2D'+S'+S''+A'+A''\rvert^2$
				K^0K^{*+}	$\lvert -2D'+S'+S''+A'+A''\rvert^2$		
				$\eta'\rho^+$	$\frac{4}{3}\lvert S'+S''\rvert^2$		

It follows then that as far as the enhanced nonleptonic decays of χ-charmed mesons are concerned, the predictions of this model coincide with those of the standard model, except if the reduced matrix elements of the $[3^*]_{RL}$ are considerably larger than those of $[6]_{RL,LL}$ (i.e., larger enough to compensate the Cabibbo suppression factor).

2. *Decays of* F_χ^+, D_χ^+, D_χ^0. From (14) it is obvious that the decays of F_χ^+, D_χ^+, D_χ^0 can be obtained from those of Tables I–III by simply setting $P=D=S=A=\alpha=0$.

IV. IMPLICATION FOR SPEAR CHARM SEARCH

We will now discuss the impact of the predictions of the right-handed current models on the SPEAR charm search.

The most serious challenge to the GIM current model seems to come from the following upper limit[3]:

$$B(D^0 \to K^-\pi^+) \le 2.9\%. \quad (15)$$

In the standard GIM model (with the assumption of sextet dominance), the $K^-\pi^+$ rate is $\frac{1}{3}$ the rate for $D^0 \to PP$. Hence one gets a rather low bound:

$$B(D^0 \to PP) \le 8.7\%. \quad (16)$$

The upper bounds, relevant to our discussion, are known experimentally[3] in the following channels:

$$B(D^0 \to \pi^+\pi^-) \le 4.1\%,$$
$$B(D^0 \to K^+K^-) \le 3.8\%, \quad (17)$$
$$B(D^0 \to \pi^+\pi^- \text{ and } K^+K^-) \le 5.1\%.$$

These experimental bounds are not restrictive for all the GIM model. In the models of class (a), however, the above decays are not Cabibbo-suppressed; thus we have to examine what sort of bounds on branching ratios result from them.

First, we will show that the bound of (16) becomes less restrictive in the class (a) models. For definiteness, let us assume that only the sextet is enhanced (the final conclusion does not depend on this assumption in any important way). We will further assume that $\langle 3 \| 6_L \| 8_S\rangle \simeq \langle 3 \| 6_{RL} \| 8_S\rangle$. One sees then from Table III that in the right-handed-current model the $K^-\pi^+$ rate is about $\frac{1}{10}$ of the rate for $D^0 \to PP$. Thus one obtains the much less restrictive bound

$$B(D^0 \to PP) \le 30\%. \quad (18)$$

We further notice that since $K^-\pi^+$ goes only via the 6_{LL}, if one assumes $\langle 3 \| 6_{RL} \| 8_S\rangle > \langle 3 \| 6_{LL} \| 8_S\rangle$, following renormalization-group

TABLE III. Relative decays of D^0.

$D^0 \to PP$		$D^0 \to VV$		$D^0 \to PV$		$D^0 \to B\bar{B}$	
	$\cos^4\theta\, \lvert P\rvert^2 \times$		$\cos^4\theta \times$		$\cos^4\theta \times$		$\cos^4\theta \times$
$\eta'\bar{K}^0$	$\frac{4}{3}$	$K^{*-}\rho^+$	$\lvert\alpha\rvert^2$	$\bar{K}^0\rho^0$	$\frac{1}{2}\lvert 2D+S-A\rvert^2$	$\Xi^0\bar{\Sigma}^0$	$\frac{1}{2}\lvert 2D+S-A\rvert^2$
π^+K^-	1	$K^{*0}\phi$	$\lvert\alpha\rvert^2$	$\pi^0\bar{K}^{*0}$	$\frac{1}{2}\lvert 2D-S-A\rvert^2$	$\Xi^0\bar{\Lambda}$	$\frac{1}{6}\lvert S+3A\rvert^2$
$\pi^0\bar{K}^0$	$\frac{1}{2}$	$K^{*0}\rho^0$	$\frac{1}{2}\lvert\alpha\rvert^2$	$K^-\rho^+$	$\lvert -D+S-A\rvert^2$	$\Sigma^0\bar{n}$	$\frac{1}{2}\lvert 2D-S-A\rvert^2$
$\eta\bar{K}^0$	$\frac{1}{6}$	$\bar{K}^{*0}\omega$	$\frac{1}{2}\lvert\alpha\rvert^2$	π^+K^{*-}	$\lvert D+S+A\rvert^2$	$\Lambda\bar{n}$	$\frac{1}{6}\lvert -S+3A\rvert^2$
	$\cos^2\theta \times$		$\cos^2\theta \times$	$\eta\bar{K}^{*0}$	$\frac{1}{6}\lvert -S+3A\rvert^2$	$\Xi^-\bar{\Sigma}^+$	$\lvert -D+S-A\rvert^2$
$\eta'\eta'$	$\lvert R'\rvert^2$	$\phi\phi$	$\lvert -\beta-\frac{1}{3}\gamma+\gamma'\rvert^2$	$\eta'\bar{K}^{*0}$	$\frac{4}{3}\lvert S\rvert^2$	$\Sigma^+\bar{p}$	$\lvert -D+S+A\rvert^2$
$\pi^+\pi^-$	$\lvert Q+\frac{1}{3}R+R'\rvert^2$	$\rho^+\rho^-$	$\lvert \beta+\frac{1}{3}\gamma+\gamma'\rvert^2$	$\bar{K}^0\phi$	$\lvert S+A\rvert^2$		$\cos^2\theta \times$
K^+K^-	$\lvert -Q+\frac{1}{3}R+R'\rvert^2$	$K^{*+}K^{*-}$	$\lvert -\beta+\frac{1}{3}\gamma+\gamma'\rvert^2$	$\bar{K}^0\omega$	$\frac{1}{2}\lvert S-A\rvert^2$	$n\bar{n}$	$\lvert -D'+2A'-\frac{2}{3}S''+S''\rvert^2$
$\pi^0\pi^0$	$\lvert \frac{1}{2}Q+\frac{1}{6}R+R'\rvert^2$	$\rho^0\rho^0$	$\lvert \frac{1}{2}\beta+\frac{1}{3}\gamma+\gamma'\rvert^2$		$\cos^2\theta \times$		$\lvert D'-2A'-\frac{2}{3}S''+S''\rvert^2$
$\eta\eta$	$\lvert \frac{1}{2}Q-\frac{1}{6}R+R'\rvert^2$	$\omega\omega$	$\lvert \frac{1}{2}\beta+\frac{1}{6}\gamma+\gamma'\rvert^2$	$\pi^-\rho^+$	$\lvert D'-S'+A'+\frac{1}{3}S''+A''+S''\rvert^2$		$\lvert -D'+S'-A'+\frac{1}{3}S''+S''\rvert^2$
$\eta\eta'$	$\left\lvert \sqrt{2}Q+\frac{\sqrt{2}}{3}R\right\rvert^2$	$\omega\rho^0$	$\lvert -\beta+\gamma\rvert^2$	$K^0\bar{K}^{*0}$	$\lvert -D'+2A'-\frac{2}{3}S''+S''\rvert^2$	$p\bar{p}$	$\lvert D'+S'+A'+\frac{1}{3}S''+A''+S''\rvert^2$
$\eta\pi^0$	$\left\lvert -\frac{Q}{\sqrt{3}}+\frac{1}{3\sqrt{3}}R\right\rvert^2$	$K^{0*}\bar{K}^{0*}$	$\lvert -\frac{2}{3}\gamma+\gamma'\rvert^2$	K^+K^{*-}	$\lvert D'+S'+A'-\frac{1}{3}S''-A''+S''\rvert^2$	$\Sigma^+\bar{\Sigma}^-$	$\lvert -D'-S-A'+\frac{1}{3}S''+A''+S\rvert^2$
$\eta'\pi^0$	$\left\lvert -\frac{2}{\sqrt{6}}Q+\frac{2}{3\sqrt{6}}R\right\rvert^2$			K^-K^{*+}	$\lvert -D'+S'-A'+\frac{1}{3}S''+A''+S''\rvert^2$	$\Sigma^-\bar{\Sigma}^+$	$\lvert D'-S'+A'+\frac{1}{3}S'-A''+S''\rvert^2$
$K^0\bar{K}^0$	$\lvert -\frac{2}{3}R+R'\rvert^2$			$\bar{K}^0K^{*0}$	$\lvert D'-2A'-\frac{2}{3}S''+S''\rvert^2$	$\Sigma^0\bar{\Lambda}$	$\frac{1}{3}\lvert -3D'+S'+S''\rvert^2$
				$\pi^+\rho^-$	$\lvert -D'-S'-A'+\frac{1}{3}S''-A''+S''\rvert^2$	$\Lambda\bar{\Sigma}^0$	$\frac{1}{3}\lvert -D'+S'+S''\rvert^2$
				$\pi^0\phi$	$2\lvert D'\rvert^2$	$\Sigma^0\bar{\Sigma}^0$	$\lvert -S'+\frac{1}{3}S''+S''\rvert^2$
				$\eta\phi$	$\frac{4}{3}\left\lvert \sqrt{2}S'+\frac{\sqrt{2}}{3}S''-\frac{S''}{\sqrt{2}}\right\rvert^2$	$\Lambda\bar{\Lambda}$	$\lvert S'-\frac{1}{3}S''+S''\rvert^2$
				$\eta'\phi$	$\frac{1}{3}\lvert 2S'-\frac{2}{3}S''+S''\rvert^2$		
				$\eta'\rho^0$	$\frac{2}{3}\lvert -S'+S''\rvert^2$		
				$\eta'\omega$	$\frac{2}{3}\lvert S'+\frac{1}{3}S''+S''\rvert^2$		
				$\eta\rho^0$	$\frac{1}{3}\lvert -3D'+S'+S''\rvert^2$		
				$\eta\omega$	$\frac{1}{3}\lvert -S'+\frac{1}{3}S''+S''\rvert^2$		
				$\pi^0\rho^0$	$\lvert -S'+\frac{1}{3}S''+S''\rvert^2$		
				$\pi^0\omega$	$\lvert -D'+S'+S''\rvert^2$		

TABLE IV. Relative decay rates of the "z" type heavy mesons.

$F_z^0 \to PP$	$\cos^2\theta\,\lvert Q+R\rvert^2\times$	$D_z^- \to PP$	$\cos^2\theta\,\lvert Q+R\rvert^2\times$	$D_z^0 \to PP$	$\cos^2\theta\times$
$K^+\pi^-$	$\frac{1}{3}$	$\pi^-\eta$	$\frac{2}{9}$	K^+K^-	$\lvert\frac{1}{3}Q-\frac{2}{3}R+R'\rvert^2$
$K^0\pi^0$	$\frac{1}{6}$	$\pi^-\eta'$	$\frac{4}{9}$	$K^0\bar{K}^0$	$\lvert\frac{1}{3}Q+\frac{1}{3}R+R'\rvert^2$
K^0	$\frac{1}{18}$	K^-K^0	$\frac{1}{3}$	$\pi^+\pi^-$	$\lvert-\frac{2}{3}Q+\frac{1}{3}R+R'\rvert^2$
$K^0\eta'$	$\frac{4}{9}$			$\eta\eta$	$\lvert\frac{1}{3}Q-\frac{1}{6}R+R'\rvert^2$
				$\pi^0\pi^0$	$\lvert-\frac{1}{3}Q+\frac{1}{6}R+R'\rvert^2$
				$\eta\eta'$	$\left\lvert\frac{4}{3\sqrt{2}}Q+\frac{4}{9\sqrt{2}}R\right\rvert^2$
				$\eta'\eta'$	$\lvert R'\rvert^2$
				$\pi^0\eta$	$\frac{1}{3}\lvert R\rvert^2$
				$\pi^0\eta'$	$\frac{2}{3}\lvert R\rvert^2$

arguments,[16] then the bound becomes even less restrictive.

We analyze, now, the significance of the upper limit (15) for the models of class (b). As mentioned earlier, the situation will not differ significantly from the GIM scheme, unless the $[3^*]_{RL}$ is substantially enhanced with respect to the $[6]_{RL,LL}$. If this enhancement does not take place, then one gets still a rather low upper limit for $B(D^0\to PP)$.

Let us now examine the significance of the bounds, given by (17), under various assumptions. First, in the case of 3^* dominance,[7] one obtains

$$B(D^0\to\pi^+\pi^-)=B(D^0\to K^+K^-)\approx\tfrac{1}{4}B(D^0\to PP). \qquad (19)$$

From (17), it follows that

$$B(D^0\to PP;\,\underline{3}^*)\lesssim 10.2\%.$$

On the other hand, if we assume sextet dominance with $\langle 3\|6_{RL}\|8_S\rangle\simeq\langle 3\|6_{LL}\|8_S\rangle$, one gets

TABLE V. Table for reduced matrix elements. For simplicity, we have used the same symbols for PV and $B\bar{B}$ reduced matrix elements. Although *a priori* they are unrelated and could be accidentally identical, in using the above tables they should be treated as different quantities.

$P,\alpha\equiv\langle 3\|6_{LL}\|8_S\rangle_{PP,VV}$	$S,S'\equiv\langle 3\|6_{LL,RL}\|8_S\rangle_{PV}$
$Q,\beta\equiv\langle 3\|6_{RL}\|8_S\rangle_{PP,VV}$	$S''\equiv\langle 3\|3^*_{RL}\|8_S\rangle_{PV}$
$R,\gamma\equiv\langle 3\|3^*_{RL}\|8_S\rangle_{PP,VV}$	$A,A'\equiv\langle 3\|6_{LL,RL}\|8_A\rangle_{PV}$
$R',\gamma'\equiv\langle 3\|3^*_{RL}\|1\rangle_{PP,VV}$	$A''\equiv\langle 3\|3^*_{RL}\|8_A\rangle_{PV}$
$D,D'\equiv\langle 3\|6_{LL,RL}\|10\rangle_{PV}$	$S'''\equiv\langle 3\|3^*_{RL}\|1\rangle_{PV}$

$$B(D^0\to\pi^+\pi^-)=B(D^0\to K^+K^-)\approx\tfrac{1}{9}B(D^0\to PP),$$

and we are led to the much less restrictive bound,

$$B(D^0\to PP;\,\underline{6})\lesssim 23\%.$$

Finally, if we assume equal contributions of all reduced matrix elements, we find

$$B(D^0\to PP;\,\underline{3}^*\text{ and }\underline{6})\lesssim 16\%.$$

From the experimental bound

$$B(D^+\to\bar{K}^0K^+)+B(F^+\to\bar{K}^0K^+)\lesssim 10\%,$$

and taking into account that (assuming again sextet dominance and $\langle 3\|6_{LL}\|8_S\rangle\simeq\langle 3\|6_{RL}\|8_S\rangle$)

$$B(D^+\to\bar{K}^0K^+)=\tfrac{1}{3}B(D^+\to PP),$$

$$B(F^+\to\bar{K}^0K^+)=\tfrac{1}{6}B(D^+\to PP),$$

one obtains

$$B(F^+\to PP)+2B(D^+\to PP)\lesssim 60\%,$$

which is not a very restrictive bound.

The upper limit obtained by Boyarski *et al.* on $B(D^+\to\bar{K}^0\pi^+$ or $K^-\pi^+\pi^+)$,

$$B(D^+\to\bar{K}^0\pi^+\text{ or }K^-\pi^+\pi^+)\lesssim 7.2\%,$$

does not present any serious difficulty. The point is that these decay modes are not supposed to be enhanced either in the GIM model or in the present right-handed current models.

V. CONCLUSION

We have analyzed the properties of the nonleptonic weak Hamiltonian for models with right-handed currents. Assuming SU(3) invariance, we have analyzed the dominant decay modes of the charmed mesons, contrasting the present predictions with those of the standard GIM model. In particular we have pointed out that there are now decays of D^+ unsuppressed by the Cabibbo factor and we predict that the nonleptonic decays of D^+ are expected to dominate over the semi-leptonic ones. With respect to the results of the SPEAR charm search, we have shown that the new current $\bar{\chi}_R\gamma_\mu n_R$ can accommodate the low upper limit set for $B(D^0\to K^-\pi^+)$. For models of class (a) one notices that there are in general more enhanced decays (each one with small branching ratio) than in the GIM case. This has implications for the current search for charm, since it makes the detection of a charmed particle somewhat more difficult.

We would like to comment now about our expectations for the masses of the new type of charm particles, i.e., those which do not coincide with χ, the GIM quark. As already stated, theoretical prejudices would lead us to believe that x and y particles are degenerate to lowest order in weak and electromagnetic interactions. So, the charm particles corresponding to these quarks should roughly have the same mass, i.e., around 2 to 4 GeV.[17] As regards the charm particles corresponding to z, the situation is quite different. First of all, present neutrino experiments would tend to suggest that, up to $E_\nu \simeq 100$ GeV, the z threshold is not open, because if it were the ratio $\sigma^{\bar{\nu}}/\sigma^{\nu}$ would go to unity.[18] Here, $\sigma^{\nu,\bar{\nu}}$ represent the charged-current total cross sections for ν or $\bar{\nu}$ scattering on nucleon targets. In such a case, our consideration of the z current is at the moment mostly academic and will be relevant only after the z threshold has been exceeded.

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