QUARK AND PARTICLE HELICITIES IN HADRONIC CHARMED PARTICLE DECAYS

J.G. KÖRNER
Deutsches Elektronen-Synchrotron DESY, Hamburg, Germany

and

Gary R. GOLDSTEIN [1]
Rutherford Laboratory, Chilton, Didcot, UK[2]

Received 25 September 1979

We investigate possibilities of obtaining information on the chirality of the charm-changing current $(\bar{c}s)$ by measuring the helicities of hadrons with spin in exclusive hadronic charmed particle decays.

As more data on charmed particle decays are becoming available, it is desirable to obtain more detailed information on the decays than is obtainable from measurements of rates alone. In particular, since the charm changing current $(\bar{c}s)$ is believed to be left-handed, it would be interesting to find out whether this helicity information is being handed down to hadrons with spin in the final state.

For semileptonic charmed particle decays this has been answered in the affirmative [1,2]. However, in semi-leptonic decays useful helicity information can only be collected in a limited q^2-range (q_μ being the momentum transfer to the leptons) due to angular momentum constraints at the kinematic boundaries. Compared to these the hadronic two-body decay modes are concentrated at a "fixed" $q^2 = m^2$ which is more favourable to a helicity analysis from rate considerations alone.

We start our considerations with charmed meson decays [3,4]. The $\Delta C = \Delta S = 1$ transitions are generated by an effective hamiltonian [3],

$$\mathcal{H}_{\text{eff}} = (G/\sqrt{2}) \cos^2\theta_c \{\tfrac{1}{2}(f_+ + f_-)(\bar{c}s)_L(\bar{d}u)_L + \tfrac{1}{2}(f_+ - f_-)(\bar{c}u)_L(\bar{d}s)_L\} . \tag{1}$$

[1] This work was supported in part by the U.S. Department of Energy.

[2] Permanent address: Tufts University, Medford, MA 02155, USA.

$\mathcal{H}_{\text{eff}}$ is the usual GIM current–current form renormalized at short distances by gluon exchange effects [5]. Without renormalization one has $f_+ = f_- = 1$. The renormalized values for charmed particle decays have been calculated in ref. [3] and are given by $f_- = f_+^{-2} = 2.15$.

The effective hamiltonian (1) leads to the quark transitions shown in the inset of fig. 1. These can be calculated in the quark model since the local form of (1) allows one to relate the diagrams of fig. 1 to one-meson current matrix elements and the meson wave-function at the origin [3,4].

To save on generality of notation we discuss the specific process $D^0 \to K^{*-}\rho^+$ (there is only one p.c. helicity amplitude in $D^0 \to K^{*-}\pi^+(K^-\rho^+)$). Our subsequent

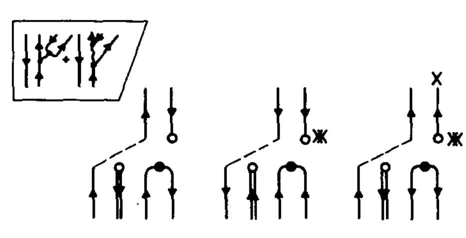

Fig. 1. Leading helicity diagrams for $D \to K^*\rho$. Arrows denote helicities. Inset shows corresponding quark line diagrams where arrows denote direction of quarks. × denotes chirality suppression and ⋇ helicity flip. Double line is charmed quark. Current–current vertices denoted by circles.

conclusions are, however, general. From the helicity diagrams in fig. 1 we expect the longitudinal helicity transition H_{00} to dominate. For H_{--} to occur quark spin has to be flipped. The (flip/nonflip) factor is $\sqrt{2}M_\rho/M_c$ on the quark level and $\sqrt{2}M_\rho/M_D$ at the particle level [6]. For H_{++} to occur one has a chirality suppression in addition to the spin flip. The chirality suppression is $\approx(M_s/M_c)$ at the quark level and $\approx(M_K/M_D)$ at the particle level [6]. Thus one expects $|H_{00}|^2 \gtrsim |H_{--}|^2 > |H_{++}|^2$ ($|H_{00}|^2 \gtrsim |H_{++}|^2 > |H_{--}|^2$ for (V+A) currents) independent of the details of the quark model. For example, using U(2,2) quark model wavefunctions as in ref. [6] one has $|H_{00}|^2 : |H_{--}|^2 : |H_{++}|^2 = 0.78 : 0.65 : 0.26$ for $D^0 \to K^{*-}\rho^+$.

It is clear that the ratio $|H_{00}|^2/\frac{1}{2}(|H_{--}|^2+|H_{++}|^2)$ tests the quark model dynamics whereas $|H_{++}|^2/|H_{--}|^2$ tests the chirality of the $(\bar{c}s)$ current. Compared to the semileptonic decay $D^0 \to K^{*-} + e^+ + \nu$ the hadronic decay is concentrated at a $q^2(=m_\rho^2)$-value which is maximally sensitive to this chirality effect.

It is, of course, a different question whether one can define an experimental measurement which is sensitive to this ratio. In order to be able to answer this question we write down the full double decay distribution for $D \to K^*(\to K\pi)\rho(\to\pi\pi)$. In terms of the double decay density matrix $\rho_{\lambda\mu,\lambda\mu} = H^*_{\lambda\lambda}H_{\mu\mu}$ one has

$$\begin{aligned} W(\theta'\varphi', \theta\varphi) &= \rho_{00,00}\cos^2\theta\cos^2\theta' \\ &+ \tfrac{1}{4}(\rho_{11,11} + \rho_{-1-1,-1-1})\sin^2\theta\sin^2\theta' \\ &+ \sin\theta\cos\theta\sin\theta'\cos\theta' \\ &\times [\cos(\varphi+\varphi')\,\mathrm{Re}(\rho_{10,10} + \rho_{-10,-10}) \\ &+ \sin(\varphi+\varphi')\,\mathrm{Im}(\rho_{10,10} - \rho_{-10,-10})] \\ &+ \tfrac{1}{2}\sin^2\theta\sin^2\theta'[\cos 2(\varphi+\varphi')\,\mathrm{Re}\,\rho_{1-1,1-1} \\ &+ \sin 2(\varphi+\varphi')\,\mathrm{Im}\,\rho_{1-1,1-1}], \end{aligned} \quad (2)$$

where $\theta\varphi$ and $\theta'\varphi'$ are the polar and azimuthal angles of $K^* \to K\pi$ and $\rho \to \pi\pi$ in their respective rest systems.

The only terms in the double differential decay distribution (2) that change sign under parity reflections can be seen to be proportional to the *imaginary* parts of the double decay density matrix. Equivalent to this statement is the observation that only the elements $(\rho_{10,10} - \rho_{-10,-10})$ and $\rho_{1-1,1-1}$ contain interferences between the p.v. amplitudes $H_T^{\text{p.v.}} = (H_{++} + H_{--})/\sqrt{2}$ and H_{00}, and the p.c. amplitude $H_T^{\text{p.c.}} = (H_{++} - H_{--})/\sqrt{2}$. However, within the factorizing approach of fig. 1 and from *CP*-invariance one predicts that the p.v. and p.c. amplitudes are all relatively real. This means that the information contained in the naive quark model prediction $H_{++} < H_{--}$ cannot be measured unless the model is extended to include final state interactions. Turning the argument around, measurements on p.v. azimuthal correlations in the decays $D \to K^*(\to K\pi)\rho(\to\pi\pi)$ would test the viability of the factorization approach which is at the basis of all present-day phenemenology of the hadronic D-decays.

Only qualitative estimates can be given at present for the strength of the final state interactions. In fig. 2a we have drawn the *t*- and *s*-channel representation of the final state interaction in the particle picture. The two descriptions are related by duality and should be considered complementary. In the *s*-channel picture the presence of 0^- and 0^+ resonances with masses $m \approx m_D$ coupling weakly to D and strongly to $K^*\rho$ could in principle lead to a strong final state interaction. However, the quark model predicts that high mass, low spin recurrences practically decouple from channels containing ground state mesons due to the presence of nodes in the wave functions of the recurrences [7]. Also the weak matrix elements coupling the D to these recurrences are not expected to be strong. This means that one does not expect any substantial final state interaction effects in the particle picture ‡1. In the quark picture imaginary parts could result from additional gluon final state interactions as depicted in fig. 2b. However, these are believed to be suppressed since they can be shown to correspond to non-leading contributions in the $1/N_c$ expansion [8].

Needless to say, it is important to check the strength of the imaginary contributions. If these turn out to be small and if their phase can be calculated reliably

‡1 This is particularly true for D^0 ($\bar{D}^0$) decays since the relevant *s*-channel is exotic in this case and thus there will be no *s*-channel resonance contributions.

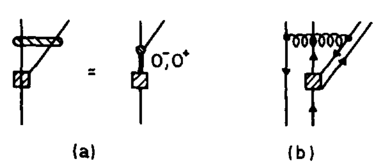

Fig. 2. (a) *s*- and *t*-channel diagrams responsible for final state interaction. (b) Representative final state interaction diagram from QCD.

from theory, the p.v. azimuthal correlations in eq. (2) can then in principle be utilized to separate H_{++} and H_{--}. Apart from this, measurements on the polar angle distribution would reveal information on the strength of H_{00} compared to $(H_{++} + H_{--})$. As remarked earlier, this separation is also useful, since it provides a further test of the dynamics of the quark model.

The absence of p.v. azimuthal correlations in the factorizing quark model approach can also be understood from a simple heuristic argument. One is trying to measure p.v. correlations of the $(\boldsymbol{p}_1 \times \boldsymbol{p}_2)\cdot\boldsymbol{p}_3$ type, which, from time reversal invariance, must receive contributions only from the imaginary parts of the transition amplitudes. The conclusion is the same as drawn from eq. (2). In this respect charmed baryon decays offer an advantage. The p.v. correlations can now be obtained from $(\boldsymbol{\sigma}\cdot\boldsymbol{p})$-type measurements from the spin orientation of the final baryon which can usually be reconstructed from its subsequent weak decay ‡2. The angular analysis is quite standard (see e.g. ref. [9]) and we can concentrate on dynamical aspects.

Our analysis is based again on quark model helicity diagrams. Unfortunately the number of quark line diagrams is larger in the baryon case than in the meson case which makes the extraction of information on the $(\bar{c}s)$-chirality more difficult in this case.

Before turning to the charmed baryon decays we justify our approach by briefly discussing the corresponding non-leptonic hyperon decays. Usually these are treated in the framework of the PCAC-current-algebra approach of ref. [10] where the main features of the data can be successfully accounted for (see e.g. ref. [11]). Since the latter approach can be seen to be equivalent to an analysis based on quark model diagrams [12–14] the quark model approach can be used as an alternative description of the hyperon decays. Also, experiments have shown that the hyperon decay amplitudes are very near to being relatively real. This means that for hyperon decays the quark model results are valid without having to include final state interaction effects.

From the above considerations we feel confident that the quark model approach can be extended to the charmed baryon case. We shall concentrate on the $(1/2^+ \to 1/2^+ + 0^-)$ decays of the three lowest ground state charmed baryon states $C_0^+(2.26)$ [15] and $A^{0,+}$ (≈ 2.47 [13]) with flavour content c [du], c [su] and c [sd]. These and the T^0(c{ss}, ≈ 2.73 [13]) are the stable members of the charmed baryon family.

The helicity diagram analysis is simplified by the observation that the non-charmed quarks in C_0^+ and $A^{0,+}$ are flavour-antisymmetric (a.s.) which implies that the two non-charmed quarks are also in an a.s. helicity configuration (see fig. 3). Further, no two quarks of the same baryon with the same helicity should end (or originate) in the weak vertex. This follows from the fact that coloured quarks effectively behave as bosons in baryons [16] and thus, from the (V – A) (V – A) structure, the ground state quarks may only have the a.s. spin and a.s. flavour configuration in this case.

One now has the necessary ingredients to draw the leading helicity diagrams (see fig. 3). The respective

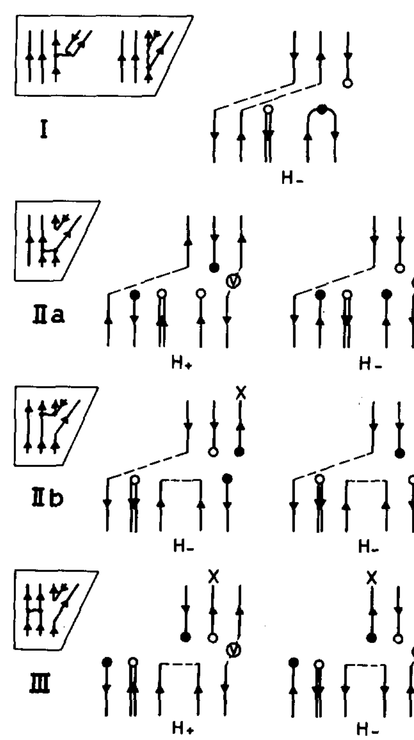

Fig. 3. Leading helicity diagrams for $C_0^+ \to 1/2^+ + 0^-$. Quark line diagrams are shown in inset. Notation as in fig. 1. V denotes pair creation from vacuum. $H_\pm$ denote helicity of final baryon in c.m. system.

‡2 For the same reason the transverse helicity separation is possible in the semileptonic decay $D \to K^* + e + \nu$, where spin information on the (off-shell) W-boson is obtained from its subsequent weak "decay".

diagrams have the following dominant helicity structure. The final baryon has: (I) negative helicity, (IIa) negative or positive helicity depending on the relative weight of the two contributions, (IIb) negative helicity, (III) negative and positive helicity with equal weights, i.e. a purely p.c. transition following from *CP*-invariance. Transitions IIb and III are chirality suppressed [‡3].

Apart from the helicity coefficients, which are given later, the relative weights of the 4 diagrams cannot be fixed in a model independent way. Except for diagrams IIa and IIb, which, from *CP*-invariance, contribute equally in the equal mass limit, the quark dynamics is different in each diagram. Diagram I is of the factorizable form which allows an absolute estimate in terms of the meson wavefunction at origin and the normalized baryon current matrix element. Diagrams IIa and IIb test the baryon wavefunction at the origin once, whereas diagram III tests the baryon wavefunction at the origin twice. From these remarks it should be clear that model independent $(\bar{c}s)$ chirality tests can only be expected from decay modes where only as few as possible diagrams contribute.

The spin-flavour content of diagrams I–III can easily be computed either by using the explicit quark model wavefunctions as listed e.g. in ref. [18] or else by using the closed expressions in terms of SU(4) invariants as derived in ref. [13]. The results of the spin-flavour contractions are listed in table 1.

From the entries of table 1 it is clear that there is only one candidate for C_0^+ decays, namely $C_0^+ \to p\bar{K}^0$,

[‡3] In ref. [17] it was argued that contributions from diagrams IIa and III should be neglected since they involve quark pair creation from the vacuum. Since the momenta of the charmed baryon decay products are not very high we do not believe that this suppression is already operative. Corresponding strong decays involving quark pair creation with similar final state momenta do not appear to be suppressed.

Table 1

Leading helicity coefficients for charmed baryon decays. Coefficients are calculated from the SU(4) invariants I_i and $\hat{I}_i$ as defined in ref. [13]. η_1 and η_8 are SU(3) singlet and octet states as in ref. [13]. The relative normalization of different diagrams is given by the leading M_1-coefficient of the explicit quark model expressions in ref. [13]. Contributions from diagrams IIb and III are chirality suppressed (x).

	H_-^{I}	H_+^{I}	H_-^{IIa}	H_+^{IIa}	$H_-^{\mathrm{IIb}\,\mathrm{x}}$	$H_+^{\mathrm{IIb}\,\mathrm{x}}$	$H_-^{\mathrm{III}\,\mathrm{x}}$	$H_+^{\mathrm{III}\,\mathrm{x}}$
	$-5I_1-4I_2$	I_1-I_2	$-\frac{3}{2}(I_3-I_4)$	$\frac{3}{2}(2I_3+I_4)$	$-\frac{3}{2}(\hat{I}_3-\hat{I}_4)$	$\frac{3}{2}(2\hat{I}_3+\hat{I}_4)$	$18I_5$	$18I_5$
$6C_0^+ \to \Lambda\pi^+$	18	0	9	0	9	0	18	18
$\sqrt{12}C_0^+ \to \Sigma^0\pi^+$	0	0	3	−6	−9	0	−18	−18
$\sqrt{12}C_0^+ \to \Sigma^+\pi^0$	0	0	−3	6	9	0	18	18
$6C_0^+ \to \Sigma^+\eta_8$	0	0	9	0	−9	0	18	18
$\sqrt{18}C_0^+ \to \Sigma^+\eta_1$	0	0	−9	9	−9	0	18	18
$\sqrt{6}C_0^+ \to p\bar{K}^0$	9	0	6	−3	0	0	0	0
$\sqrt{6}C_0^+ \to \Xi^0K^+$	0	0	3	3	0	0	18	18
$\sqrt{6}A^+ \to \Sigma^+\bar{K}^0$	9	0	0	0	9	0	0	0
$\sqrt{6}A^+ \to \Xi^0\pi^+$	−9	0	0	0	−9	0	0	0
$6A^0 \to \Lambda\bar{K}^0$	−9	0	−9	9	9	0	18	18
$\sqrt{12}A^0 \to \Sigma^0\bar{K}^0$	−9	0	−3	−3	−9	0	−18	−18
$\sqrt{6}A^0 \to \Sigma^+K^-$	0	0	3	3	0	0	18	18
$\sqrt{12}A^0 \to \Xi^0\pi^0$	0	0	−6	3	9	0	0	0
$6A^0 \to \Xi^0\eta_8$	0	0	0	−9	−9	0	−36	−36
$\sqrt{18}A^0 \to \Xi^0\eta_1$	0	0	−9	9	−9	0	18	18
$\sqrt{6}A^0 \to \Xi^-\pi^+$	9	0	6	−3	0	0	0	0
$T^0 \to \Xi^0\bar{K}^0$	−1	2	0	0	−3	6	0	0
$6C_0^+ \to \Lambda K^+$	18	0	0	9	−9	0	36	36
$\sqrt{6}C_0^+ \to \Sigma^+K^0$	0	0	0	−9	9	0	0	0

in which the $(\bar{c}s)$ chirality is directly transmitted to the final baryon. The decay is dominated by diagram IIa since diagram I is strongly suppressed in this case due to the colour–flavour factor $\chi_- = (2f_+ - f_-)/3 = -0.26$ multiplying it. The proton is thus predicted to emerge with dominant negative helicity, i.e. $\alpha = (|H_+|^2 - |H_-|^2)/(|H_+|^2 + |H_-|^2)$ = negative. Unfortunately, measurements of the proton's polarization require double scattering and are very difficult.

The situation improves for the $A^{+,0}$ decays. $A^+ \to \Sigma^+\bar{K}^0$ decays predominantly via IIb and thus the asymmetry α is predicted to be negative. This would not be too difficult to measure in the maximally ($\alpha \approx -1$) p.v. decay mode $\Sigma^+ \to p\pi^0$. The decay $A^+ \to \Xi^0\pi^+$ occurs via I and IIb with a predicted negative asymmetry. Again there is enough asymmetry in the Ξ^0 decay modes to detect this.

For the sake of completeness we have also included in table 1 the Cabibbo-suppressed decay modes. Of these $C_0^+ \to \Sigma^+K^0$ is predicted to marginally have a positive asymmetry since IIb is chirality suppressed relative to IIa. We have also calculated the only T^0 decay mode $T^0 \to \Xi^0K^0$. The corresponding helicity diagrams are not included in fig. 3 but can easily be drawn. In this case, the quark symmetries differ. We predict a positive asymmetry in this case.

All other decay modes obtain (coherent!) contributions from several diagrams and model independent predictions on their decay asymmetries are no longer possible. Asymmetry measurements on these decay modes would nevertheless be important since they would provide further information on the details of the quark model dynamics. A good example is the decay $C_0^+ \to \Lambda\pi^+$ which is predicted to be purely p.c. ($\alpha = 0$) in the SU(4) symmetry limit [13,19]. An asymmetry measurement in this decay would thus provide information on the mechanism of SU(4) breaking. Similarly $C_0^+ \to \Xi^0K^+$ is predicted to be purely p.c. in the quark model approach of ref. [13] which includes SU(4) breaking effects. It would be interesting to check this prediction.

A similar helicity analysis can be done on the decay channels $1/2^+ \to 1/2^+ + 1^-, 3/2^+ + 0^-, 3/2^+ + 1^-$. These decays provide some additional information on the helicity structure. However, there are no principally new features. They will be treated in a separate publication. Among these there are the two interesting decay modes $A^0 \to \Omega^-K^+$ and $T^0 \to \Omega^-\pi^+$ which, in models with (V – A) (V – A) structure, are predicted to be purely p.c. [13]. Unfortunately this will be hard to confirm since the same dynamics also predicts the Ω^- decay modes to be purely p.c. [‡4] and thus the Ω^--polarization cannot be reconstructed.

Apart from the above-mentioned asymmetry measurements it is of course quite important to check on the relative phase of p.v. and p.c. amplitudes by measuring the β-parameter [9] ($\beta = \mathrm{Im}(H_+H_-^*)/(|H_+|^2 + |H_-|^2)$). Similar to the meson case discussed earlier, nonvanishing relative phases would indicate that final state interactions are present. If these should turn out to be strong and strongly influence the helicities of the final state particles, the present approach based on simple quark transitions would have to be modified.

To summarize, we have used simple quark transition diagrams and a left-handed $(\bar{c}s)$-current in an attempt to make model-independent predictions for the helicity structure of charmed meson and charmed baryon decays into channels with spins. We have discussed possibilities of measuring the predicted helicity structure. For charmed meson decays the proposed angular correlation measurements should not be too difficult using the clean sample of D-mesons from the decay of the $\psi''(3.77)$. For charmed baryon decays definite helicity predictions were obtained only for a few decay channels. For C_0^+-decays there is only the decay $C_0^+ \to p\bar{K}^0$ where the proton's helicity is hard to measure. Some good and measurable candidates were found among $A^{+,0}$ decays. Since $A^{+,0}$ are not likely to be produced copiously in neutrino interactions the proposed measurements would have to be done in e^+e^--interactions. These measurements will have to wait for the early 80's when the ARGUS and MARK III detectors become operative at DORIS and SPEAR, respectively.

J.G.K. acknowledges the hospitality of the Rutherford Laboratory where this work was initiated. We would like to thank M. Fugukita and J. Kühn for helpful discussions.

‡4 This has recently been confirmed in the decay $\Omega^- \to \Lambda K^-$ [20].

References

[1] R. Nabavi, X.Y. Pham and W.N. Cottingham, J. Phys. G3 (1978) 1485.

[2] A.J. Buras and J. Ellis, Nucl. Phys. B111 (1976) 341; G. Köpp, L.M. Sehgal and P.M. Zerwas, Nucl. Phys. B123 (1977) 61; B123 (1977) 77.

[3] N. Cabibbo and L. Maiani, Phys. Lett. 73B (1978) 418.
[4] D. Fakirov and B. Stech, Nucl. Phys. B133 (1978) 315.
[5] G. Altarelli and L. Maiani, Phys. Lett. 52B (1974) 351; M.K. Gaillard and B.W. Lee, Phys. Rev. Lett. 33 (1974) 108.
[6] A. Ali, J.G. Körner, G. Kramer and J. Willrodt, Z. Phys. C1 (1979) 269.
[7] M. Böhm, H. Joos and M. Krammer, Nucl. Phys. B69 (1974) 349;
E. Eichten, K. Gottfried, T. Kinoshita, K.D. Lane and T.M. Yan, Phys. Rev. D17 (1978) 3090;
A. Le Youanc, L. Oliver, O. Pene and J.C. Roynal, Phys. Lett. 71B (1977) 397.
[8] M. Fukugita, T. Inami, N. Sakai and S. Yazaki, Phys. Lett. 72B (1977) 237.
[9] R.E. Marshak, Riazuddin and C.P. Ryan, Theory of weak interactions (Wiley–Interscience).
[10] L.S. Brown and C.M. Sommerfield, Phys. Rev. Lett. 16 (1966) 751;
Y. Hara, Y. Nambu and J. Schechter, Phys. Rev. Lett. 16 (1966) 380.
[11] H. Galic, D. Tadic and J. Trampetic, Zagreb preprint IRB-TP-1-79 (1979).
[12] J.G. Körner and T. Gudehus, Nuovo Cimento 11A (1972) 597.
[13] J.G. Körner, G. Kramer and J. Willrodt, Z. Phys. C2 (1979) 117.
[14] B. Desplanques, J.F. Donoghue and B.R. Holstein, MIT Preprint CTP 754 (1979).
[15] E.G. Cazzoli et al., Phys. Rev. Lett. 34 (1975) 1125;
B. Knapp et al., Phys. Rev. Lett. 37 (1976) 882;
C. Baltay et al., Phys. Rev. Lett. 42 (1979) 1721.
[16] K. Miura and T. Minamikawa, Prog. Theor. Phys. 38 (1967) 954;
J.G. Körner, Nucl. Phys. 25B (1970) 282;
J.C. Pati and C.W. Woo, Phys. Rev. D3 (1971) 2920.
[17] L. Maiani, INFN preprint, Rome (1978).
[18] W. Thirring, in: Proc. Schladming Winter School (1965), ed. P. Urban (Springer, 1965).
[19] H.R. Rubinstein, L. Stodolsky, F. Wagner and U. Wolff, Phys. Lett. 73B (1978) 433.
[20] M. Bourquin et al., CERN preprint (1978).