PHYSICAL REVIEW LETTERS

VOLUME 45 | 24 NOVEMBER 1980 | NUMBER 21

Signature $F^+ \rightarrow p\bar{n}$ as an Unambiguous Proof of the Annihilation Mechanism

Xuan Yem Pham
Laboratoire de Physique Théorique et Hautes Energies, Associé au Centre National de Recherche Scientifique, Université Pierre et Marie Curie, F-75230 Paris 05, France
(Received 16 June 1980)

The spectacular decay mode $F^+ \rightarrow p\bar{n}$ would be easily observed if the annihilation mechanism actually dominates the standard one (light-quark spectator) in the disintegrations of charmed particles. Its branching ratio would be about 10^{-2}. Otherwise, it should be very rare ($\simeq 10^{-6}$) because of partial conservation of axial current.

PACS numbers: 14.40.Pe, 12.30.Ez, 13.25.+m

Suppose one observes a nice bump in the nucleon-pair invariant mass with extremely narrow width (less than 10^{-4} eV, which is partially hidden by experimental resolutions); for the first time one will have a meson that decays weakly into a baryon pair. Up to now, all baryons decay into other baryons, all mesons into other mesons.

From the theoretical point of view, we find the mode $F^+ \rightarrow p\bar{n}$ interesting in two aspects because of its relation with the following:

(i) the partial conservation of axial current[1] (PCAC);

(ii) the annihilation mechanism[2] recently suggested as the dominant one that governs particle disintegrations, in particular charmed meson decays. This model was proposed in order to understand the large difference in the D^+ and D^0 lifetimes recently reported.[3]

First I remark that the mode $F^+ \rightarrow p\bar{n}$ must occur for any decay mechanism. The only question raised is how large is its branching ratio. The purpose of this note is to point out that measurement of this ratio will provide an unambiguous way to test the two theoretical points mentioned above.

The key of the annihilation mechanism as proposed in Ref. 2 is to realize that the valence pair $c\bar{q}$ of the charmed pseudoscalar mesons is not necessarily in the $^1S_0$ (spectroscopic notation) bound state as implicitly assumed in the standard mechanism.[4] Because of gluons, the valence-quark pair can be in all possible configurations, in particular the $^3S_1$ state. Colorless and spinless properties of physical mesons can be arranged with the aid of gluons.

This idea means that the decay width $F^+ \rightarrow p\bar{n}$ can be written as

$$\Gamma(F^+ \rightarrow p\bar{n}) = M(G^2/8\pi)\cos^4\theta_c(1-4m^2/M^2)^{1/2}(F^{\mu\nu}G_{\mu\nu}). \tag{1}$$

Here phase-space factor, coupling constant, etc., have been factored out, and the dynamics are focused in the matrix element squared, $F^{\mu\nu}G_{\mu\nu}$. The tensor $G_{\mu\nu}$ corresponding to the final $p\bar{n}$ state is calculated to be

$$G_{\mu\nu} = g_V^2(t)\left(p_\mu n_\nu + p_\nu n_\mu - \frac{M^2}{2} g_{\mu\nu}\right) + g_A^2(t)\left(p_\mu n_\nu + p_\nu n_\mu - \frac{M^2 - 4m^2}{2} g_{\mu\nu}\right)$$
$$+ g_T^2(t)\left[\frac{M^2 + 4m^2}{8m^2} q_\mu q_\nu - \frac{M^2}{4m^2}(p_\mu n_\nu + p_\nu n_\mu) - \frac{M^2}{2} g_{\mu\nu}\right] + g_P^2(t)\frac{M^2}{8m^2} q_\mu q_\nu$$
$$+ g_A(t) g_P(t) q_\mu q_\nu + g_V(t) g_T(t)(M^2 g_{\mu\nu} - q_\mu q_\nu). \qquad (2)$$

The antisymmetric part of $G_{\mu\nu}$ proportional to $i\epsilon_{\mu\nu\alpha\beta}p^\alpha n^\beta$ does not contribute to the width since it is contracted with the symmetric tensor $F^{\mu\nu}$ as will be discussed later.

We denote by q, p, and n the four-momentum of F meson, proton, and neutron, respectively. M and m are their masses. Also $q^2 = (p+n)^2 = t$. The four dimensionless form factors g_V, g_T, g_A, and g_P are defined in a similar way as in nucleon β decay:

$$\langle 0 | V_\mu + A_\mu | p\bar{n}\rangle = \bar{u}(p)\left[\gamma_\mu \tilde{g}_V(t) + i\delta_{\mu\nu}\frac{q_\nu}{2m}\tilde{g}_T(t) + \gamma_\mu\gamma_5 \tilde{g}_A(t) + \frac{q_\mu}{2m}\gamma_5 \tilde{g}_P(t)\right] v(n).$$

In principle $g_{V,T,A,P}$ might be different[5] from $\tilde{g}_{V,T,A,P}$ as can be seen in Fig. 1. The form factors $\tilde{g}_{V,T,A,P}$ are associated indeed to the vacuum-nucleon pair transition, while for the form factors $g_{V,T,A,P}$ the effect of gluons has to be smeared out when we calculate the width. These effects are put into the falloff of form factors. After all, in the timelike region considered here, the t dependence of $\tilde{g}_{V,T,A,P}$ is not well known either.

At zero momentum transfer, from the conserved vector current (CVC), we have $\tilde{g}_V(0) = 1$, $\tilde{g}_T(0) = \mu_p - \mu_n = 3.7$. Also $\tilde{g}_A(0) \simeq 1.2$. The induced pseudoscalar term $\tilde{g}_P(t)$ has a pole at the pion mass, and the residue at the pion pole can be fixed via the Goldberger-Treiman (G.T.) relation:

$$\tilde{g}_P(t) = 4m^2 \tilde{g}_A(0)/(m_\pi^2 - t) + \cdots .$$

This term is known to play an important role in muon capture and the relative sign between $\tilde{g}_A$ and g_P is fixed by PCAC.

We now come to the tensor $F^{\mu\nu}$ associated to the initial state, the F^+ meson. Since F^+, as a spin-0 particle, has only one degree of freedom, i.e., its four-momentum q, the most general form[6] of $F^{\mu\nu}$ is $-aM^2 g^{\mu\nu} + b q^\mu q^\nu$, where a and b are two positive constants related somehow to the wave function[7] of the F^+ meson. In the standard mechanism[4] where the $c\bar{s}$ valence-quark pair is in the ${}^1S_0$ configuration, then $a = 0$ and $b = f_F^2/M^2$ corresponding to the vacuum-F^+ transition. Here f_F, the F^+ decay coupling constant, is defined in a way similar to that of the pion f_π.

Putting $a = 0$ and $b = f_F^2/M^2$ in Eqs. (1) and (2), we recover the result[8]

$$\Gamma(F^+ \to p\bar{n}) = (G^2 \cos^4\theta_c/16\pi) f_F^2 M (1 - 4m^2/M^2)^{1/2} |D(t = M^2)|^2, \qquad (3)$$

where

$$D(t) = 2m\tilde{g}_A(t) + (t/2m)\tilde{g}_P(t). \qquad (4)$$

$D(t)$ is the matrix element between nucleons of the operator $\partial^\mu A_\mu$. In the standard mechanism, the decay matrix element $F^+ \to p\bar{n}$ is proportional to the divergence of the hadronic axial current; therefore measurement of this mode is equivalent to measurement—at the charmed mass—of the deviation from the *asymptotic conserved axial current,* $D(t) \to 0$ as $t \to \infty$, a crucial point for having an unsubtracted dispersion relation for $D(t)$. If one writes

$$D(t) = \frac{\sqrt{2}\, m_\pi^2 f_\pi g_{\pi NN}}{m_\pi^2 - t} + \frac{1}{\pi}\int_{(3m_\pi)^2}^{\infty} dx \frac{\rho(x)}{x - t}, \qquad (5)$$

where the pion pole is shown explicitly in the first term, puts $t = 0$ on both sides of Eqs. (4) and (5), and neglects the contribution from higher states $\beta \equiv \pi^{-1}\int dx\, \rho(x)/x$, then the G.T. relation $2m\tilde{g}_A(0)$

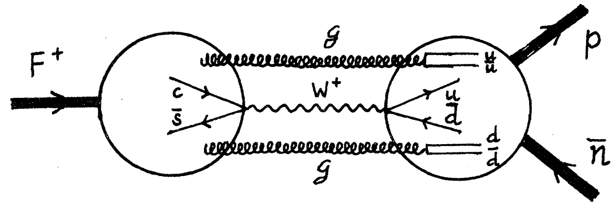

FIG. 1. The decay $F^+ \to p\bar{n}$ by the annihilation mechanism.

$=\sqrt{2}f_\pi g_{\pi NN}$ is obtained.[9]

Putting now Eq. (5) into Eq. (3) and neglecting the integral term $\gamma \equiv \pi^{-1}\int dx\,\rho(x)/(x-M^2)$, then we get, using the G. T. relation,

$$\Gamma(F^+\to p\bar{n}) = \frac{G^2}{4\pi}\cos^4\theta_c f_F{}^2 M m^2\left(1-\frac{4m^2}{M^2}\right)^{1/2}\tilde{g}_A{}^2(0)\left(\frac{m_\pi}{M}\right)^4 \simeq 4.5\times10^{+6}\left(\frac{f_F}{f_\pi}\right)^2\ \mathrm{sec}^{-1}. \tag{6}$$

For a total lifetime of F^+ between 1 and 5 (10^{-13} sec), the value commonly accepted, the branching ratio $B(F^+\to p\bar{n})$ is found to be extremely small ($\simeq 10^{-6}$).

As suggested previously,[8] if there exists some heavy resonances $J^P=0^-$, $I^G=1^-$ with masses not far from the charmed mass and coupled strongly to nucleons, for example a heavy pion π' (and/or some baryonium states), then the branching ratio $B(F^+\to p\bar{n})$ might be larger. However, it is unlikely that it can reach 10^{-4} as can be seen by the following explicit example. Since the pion-pole term contributes negligibly, a possible enhancement for $F^+\to p\bar{n}$ can only come from the $\gamma \equiv \pi^{-1}\int dx\rho(x)/(x-M^2)$ term. Let us approximate $\rho(x)$ by

$$\rho(x)=\sqrt{2}\,\pi f_{\pi'} g_{\pi' NN} m_{\pi'}{}^2\delta(m_{\pi'}{}^2-x)$$

with $m_{\pi'}\simeq M$. In this case,

$$\Gamma(F^+\to p\bar{n}) = \frac{G^2}{4\pi}\cos^4\theta_c f_F{}^2 M m^2\left(1-\frac{4m^2}{M^2}\right)^{1/2}\epsilon^2\frac{m_{\pi'}{}^4}{(m_{\pi'}{}^2-M^2)^2+\Gamma_{\pi'}{}^2 m_{\pi'}{}^2}, \tag{7}$$

where ϵ denotes the ratio $\epsilon\equiv\sqrt{2}f_{\pi'}g_{\pi' NN}/[2m\tilde{g}_A(0)]=[2m\tilde{g}_A(0)/\sqrt{2}f_\pi g_{\pi NN}]-1$. The coefficient ϵ is, however, very small because the G. T. relation is known to agree with experiment within a 5% limit. This fact puts an upper bound on β and consequently on $B(F^+\to p\bar{n})$. Presumably an upper limit 10^{-4} for $B(F^+\to p\bar{n})$ is reasonable in the standard mechanism ($a=0$).

The situation will be completely changed when the annihilation mechanism[2] is switched on ($a\neq 0$). If this model actually dominates the standard one, then obviously the width $\Gamma(F^+\to p\bar{n})$ should be also enhanced with respect to the canonical value of Eq. (6). In this case the decay matrix element is no longer related to PCAC, and the suppression factor $(m_\pi/M)^4$ disappears in the width. It is now the $-aM^2g^{\mu\nu}$ term of the tensor $F^{\mu\nu}$ that gives rise to the large nonleptonic widths of D^0 and F^+ (the nonleptonic width of D^+ in this model is suppressed by $\tan^2\theta_c$ with respect to that of F^+). We assume the constant a to be the same for all charmed mesons F and D. This parameter a can then be determined by the lifetimes of D^0, F^+, or by the D^0 semileptonic branching ratio ($B_{sl}{}^{D^0}$). A straightforward calculation gives

$$\Gamma(F^+\to \mathrm{hadrons})=3[\Gamma_0+a(G^2M^5/4\pi)\cos^2\theta_c], \tag{8}$$

where Γ_0 is the standard width $G^2m_c{}^5/192\pi^3$. Also

$$(B_{sl}{}^{D^0})^{-1}=2+3[1+48\pi^2a(M_D/m_c)^5\cos^2\theta_c]. \tag{9}$$

In this model, the width $\Gamma(F^+\to p\bar{n})$ becomes

$$\begin{aligned}\Gamma(F^+\to p\bar{n}) &= a(G^2M^5/8\pi)(1-4m^2/M^2)^{1/2}\cos^4\theta_c\\ &\times\left\{g_V{}^2(M^2)\left(1+\frac{2m^2}{M^2}\right)+g_A{}^2(M^2)\left(1-\frac{6m^2}{M^2}\right)+g_T{}^2(M^2)\left(1+\frac{M^2}{8m^2}\right)-g_P{}^2(M^2)\frac{M^2}{8m^2}\right.\\ &\left.\qquad -g_A(M^2)g_P(M^2)-3g_V(M^2)g_T(M^2)\right\}.\end{aligned} \tag{10}$$

Assuming a common dipole form $f(t)=(1-t/\Delta^2)^{-2}$ for the form factors

$$g_V(t)=f(t),\quad g_A(t)=\tilde{g}_A(0)f(t)\simeq 1.2f(t),\quad g_T(t)=(\mu_p-\mu_n)f(t)\simeq 3.7f(t),\quad g_P(t)=[4m^2\tilde{g}_A(0)/(m_\pi{}^2-t)]f(t),$$

with $\Lambda^2=1\ \mathrm{GeV}^2$, one finds, from Eqs. (8) and (10), $B(F^+\to p\bar{n})\simeq 5\times10^{-3}$–$10^{-2}$, which is mostly independent of the parameter a, and quite intensitive to Λ^2 (between 0.71 and 1.2 GeV^2).

Concerning a possible enhancement by some heavy resonances with masses near the charmed mass, my previous remark, of course, applies here also. Moreover, in this case, since not only a heavy pion π' but also any heavy resonances $J^P=0^+$ can contribute, the possibility for enhancement becomes doubled.

Four years ago,[10] the charmed D mesons were identified through the mode $D \to K\pi$ with the branching ratio of only 2%; we therefore believe that the search for the mode $F^+ \to p\bar{n}$ is not hopeless,[11] provided that the annihilation mechanism correctly describes weak hadronic decays.

I am grateful to Professor M. Gourdin for helpful discussions and encouragement.

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