TABLE I. The charge states of Xe^{131} following the isomeric transition of Xe^{131m}.[a]

Charge	Percent intensity	Charge	Percent intensity	Charge	Percent intensity
1	0.62 ± 0.12	11	6.22 ± 0.10	21	0.012 ± 0.001
2	0.97 ± 0.03	12	3.01 ± 0.04	22	0.013 ± 0.001
3	1.62 ± 0.06	13	1.78 ± 0.03	23	0.006 ± 0.001
4	4.26 ± 0.06	14	1.09 ± 0.01	24	0.010 ± 0.001
5	5.36 ± 0.10	15	0.55 ± 0.01	25	0.004 ± 0.002
6	10.42 ± 0.13	16	0.242 ± 0.004	26	0.004 ± 0.001
7	15.69 ± 0.25	17	0.098 ± 0.002	27	0.002 ± 0.001
8	20.88 ± 0.19	18	0.036 ± 0.001	28	0.007 ± 0.001
9	15.74 ± 0.24	19	0.023 ± 0.001	29	0.006 ± 0.002
10	11.32 ± 0.19	20	0.014 ± 0.001	30	0.001 ± 0.001

[a] Average charge $= 8.04 \pm 0.04$. This may be compared with the value 8.5 ± 0.3 reported by M. L. Perlman and J. A. Miskel [Phys. Rev. **91**, 899 (1953)].

tion of this kind, there is of course no electrostatic perturbation arising from a change in nuclear charge, and what is observed is the atomic consequence of the simple removal of an innershell electron through internal conversion.

The charge 1 state corresponds to the emission of the original conversion electron, and its intensity shows how rarely the vacancy is filled by purely radiative transitions. What usually happens (charge 8) is that the whole outer electron shell is depopulated. Charge 27 corresponds to the removal of half of the normal complement of electrons in xenon. We believe that there is real intensity in the charge spectrum out to charge 29.

Note added in proof.—Since these measurements were made, an experimental effect has been found that throws the figures for the higher charges into doubt. The relative intensities up to charge 20 are probably approximately correct, but pending remeasurement, those above charge 20 must be viewed with caution.

[1] A. H. Snell and F. Pleasanton, Phys. Rev. **100**, 1396 (1955).

Interpretation of *K*-Meson Decays*

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THIS note is concerned with a possible explanation of the K-meson lifetimes that may be subject to experimental verification. The only previous explanation[1] of the near equality of the τ and θ lifetimes involves the assumption of electromagnetic conversion of one heavy meson into the other. If τ and θ are not both spin-zero particles, then Lee and Orear require either different strangeness assignments for τ and θ, or a finite but exceptionally small mass difference ($\lesssim 100$ kev) between τ and θ. While this latter possibility would be very hard to rule out experimentally, such an accidental near-coincidence of masses would be exceptional. If τ and θ are 0^- and 0^+ particles, respectively, a large mass difference ($\gtrsim 10$ Mev) is needed. The mass differences, if any, are considerably smaller than this.[2]

Unless τ and θ have rather high spins or a very peculiar internal structure, only one can decay into three pions and only the other into two pions. They are both, however, capable of decaying in the $K_{\mu2}$, $K_{\mu3}$, and K_β modes. (For simplicity, in what follows, $K_{\mu3}$ and K_β will be included in the $K_{\mu2}$ mode.) The explanation of the nearly equal lifetimes to be proposed here is as follows: τ and θ are two parents that decay into $\mu+\nu$ at the same rate $R_{\tau\mu}=R_{\theta\mu}$; for both τ and θ this is the predominant decay mode. Then the τ and θ lifetimes are bound to be approximately equal, the differences in lifetime arising through the remaining $K_{\pi3}$ and $K_{\pi2}$ modes.

On this hypothesis, $K_{\mu2}$ should show a composite differential decay curve unless the τ and θ lifetimes are equal. The present data show no significant difference in $R_\tau=1/\tau_\tau$ and $R_\theta=1/\tau_\theta$, the total transition rates for τ and θ. From the fractional abundances N_μ, N_2, N_3 of $\mu2$, $\pi2$, $\pi3$, respectively, one finds:

(a) If, as proposed here, $R_{\tau\mu}=R_{\theta\mu}=R_\mu$, then the fraction f of τ's in the K beam at production must be small,

$$f=N_3/(1-N_\mu)=0.18,$$

and

$$R_\mu/R=N_\mu=0.62,$$

where $R=R_\theta=R_\tau$.

(b) Lee and Yang[3] and Gell-Mann[4] have proposed describing the strange particles by a new operator $\mathcal{R}$, which is conserved in strong interactions. A prediction of their theory is that τ and θ should be produced in equal fractions $f=1-f=\frac{1}{2}$. Then the small amount N_3 of $\pi3$ observed requires that τ decay predominantly into the $\mu2$ mode,

$$R_{\theta\mu}/R=1-2N_2=0.38, \quad R_{\tau\mu}/R=1-2N_3=0.87.$$

The conjecture of equal θ and τ transition rates into $\mu+\nu$ and the $\mathcal{R}$-degeneracy picture are definitely incompatible. In fact, if one assumes both $R_{\theta\mu}=R_{\tau\mu}$ and $f=\frac{1}{2}$, the observed ratios N_μ:N_2:N_3 can be obtained[5] only if the τ lifetime is unreasonably short ($\sim$ one-third the θ lifetime).

An attractive feature of possibility (a) is that one can assume a universal meson-lepton interaction.[6] If this interaction is vector or axial (which helps explain the absence of $e+\nu$ decays) and if the universal coupling constant is fitted to the pion lifetime, one then obtains $\tau_\theta=\tau_\tau=1/R=1.1\times10^{-8}$ second. This quantity is in good agreement with experiment.[7]

The whole question admits of an experimental determination. If different transition rates R_τ and R_θ are detected, then an analysis of the $K_{\mu2}$ decay rate

$$dN_\mu/dt=R_{\tau\mu}f\exp(-R_\tau t)+R_{\theta\mu}(1-f)\exp(-R_\theta t),$$

and of the abundance overdetermines the quantities f, $R_{\tau\mu}$, $R_{\theta\mu}$. It may turn out that neither $f=\frac{1}{2}$ nor $R_{\tau\mu}=R_{\theta\mu}$ is tenable!

The $K_{\mu2}$ lifetime now observed is the mean,

$$\bar{\tau}_\mu=\int t\,dN_\mu \Big/ \int dN_\mu.$$

If $\tau_\tau \approx \tau_\theta$, then

$$\bar{\tau}_\mu \approx \tau_\theta + f(\tau_\tau - \tau_\theta) = \tau_\tau + (1-f)(\tau_\theta - \tau_\tau),$$

independently of $R_{\tau\mu}$ and $R_{\theta\mu}$. At present, the lifetime differences are not known accurately enough to determine f at all.[8] Of course, once f is determined, the abundance ratios

$$N_\mu : N_2 : N_3 = f\frac{R_{\tau\mu}}{R_\tau}\exp(-R_\tau t) + (1-f)\frac{R_{\theta\mu}}{R_\theta}\exp(-R_\theta t)$$
$$:(1-f)\frac{R_\theta - R_{\theta\mu}}{R_\theta}\exp(-R_\theta t) : f\frac{R_\tau - R_{\tau\mu}}{R_\tau}\exp(-R_\tau t)$$

enables one to find $R_{\tau\mu}$ and $R_{\theta\mu}$.

The theoretical reason for conjecturing equal transition rates for $\tau \rightarrow \mu + \nu$ and $\theta \rightarrow \mu + \nu$ is that, since the neutrino is massless, the pseudovector decay of a pseudoscalar meson proceeds at the same rate as the vector decay of a scalar meson.[6] In fact, it is precisely the neutrino processes that are unaffected by the presence or absence of a γ_5.[9] If τ and θ turn out to have equal masses and spins and to differ only in parity, it seems reasonable to assume that they have similar couplings to the lepton field. The rate and equality of the lifetimes of K-meson decay are then directly connected with the slow pion decay. Only the "minor" K decays involving pions then remain strange.

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[1] T. D. Lee and J. Orear, Phys. Rev. **100**, 932 (1955).
[2] M. N. Whitehead *et al.* and W. H. Barkas *et al.*, private communication. I am indebted to members of the emulsion groups at the Radiation Laboratory for information on the K-meson masses and abundances in advance of publication.
[3] T. D. Lee and C. N. Yang (to be published).
[4] M. Gell-Mann, private communication.
[5] M. Lynn Stevenson, private communication.
[6] S. A. Bludman and M. A. Ruderman, Phys. Rev. **101**, 910 (1956).
[7] L. W. Alverez *et al.*, Phys. Rev. **100**, 1264 (A) (1955); V. Fitch and R. Motley (to be published).
[8] L. W. Alverez and S. Goldhaber, Nuovo cimento **2**, 344 (1955); Harris, Orear, and Taylor, Phys. Rev. **100**, 932 (1955).
[9] See C. N. Yang and J. Tiomno, Phys. Rev. **79**, 495 (1950), on nucleon β decay.

Crossing Symmetry in Meson-Nucleon Scattering

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In recent years, attention has been drawn to the "crossing" symmetry[1] applied to the meson-nucleon scattering amplitude. This symmetry, as we shall see, is nothing but an expression of the Bose statistics of the mesons. One can find another "crossing" symmetry which is deduced from the Fermi statistics of the nucleons and their behavior under charge conjugation. It is shown that the restrictions placed upon the meson-nucleon scattering amplitude by these symmetries are identical on the energy shell.

In a charge-independent, relativistically invariant field theory, the meson-nucleon propagator may be written as

$$T_{ij}(p,p';q,q') = \int \{\exp[i(px - p'x' + qy - q'y')]\} \times \langle T(\psi(x)\bar{\psi}(x')\phi_i(y)\phi_j(y'))\rangle dxdx'dydy', \quad (1)$$

where p', q' are the nucleon and meson incoming four momenta, respectively, and p, q are the outgoing momenta, and $\langle T(\;)\rangle$ means the true vacuum expectation value of the T-product of the Heisenberg operators. Implied in the definition (1) is invariance under charge conjugation, Schwinger space-time reversal, and the crossing symmetry. Making use of these invariance properties places restrictions on the form of T_{ij}. For example, we shall deduce the restriction of the crossing symmetry.[2] Since $\phi_i(y)$, $\phi_j(y')$ are Bose fields, they commute inside a T-product, and since y and y' are dummy variables, the right side of (1) is equal to

$$\int \{\exp[i(px - p'x' - q'y + qy')]\} \times \langle T(\psi(x)\bar{\psi}(x')\phi_j(y)\phi_i(y'))\rangle dxdx'dydy' = T_{ji}(p,p';-q',-q). \quad (2)$$

We may similarly deduce restrictions placed upon T_{ij} by using charge conjugation[3] and the Fermi statistics of the nucleons. Another restriction is that of Schwinger space-time reversal.[4] Charge conjugation invariance leads to

$$T_{ij}(p,p';q,q') = CT_{ij}{}^t(-p',-p;q,q')C^{-1} \quad (3)$$

where in the usual representation $C = \alpha_2\tau_2$ and t means spinor and isotopic spinor transpose.

Invariance under Schwinger space-time reversal leads to[5]

$$T_{ij}(p,p';q,q') = \gamma_5 T_{ij}(-p,-p';-q,-q')\gamma_5. \quad (4)$$

The crossing symmetry (2) is

$$T_{ij}(p,p';q,q') = T_{ji}(p,p';-q',-q). \quad (5)$$

Since the theory is invariant under Lorentz transformations and spatial reflections, T_{ij} can be considered as a function of the quantities γp, etc.; and since γ_5 anticommutes with γ_μ, (4) is just an identity. Thus Schwinger space-time reversal puts no restriction on T_{ij}.[6]

We now maintain that the restrictions placed on T_{ij} by (3) are identical to those of (5). We prove this