by J. L. Osborne and A. W. Wolfendale (D. Reidel, Dordecht, 1974), p. 97.
[10]N. N. Kalmykov *et al.*, in *Proceedings of the Sixteenth International Conference on Cosmic Rays, Kyoto, Japan, 1979* (University of Tokyo, Tokyo, Japan, 1979), Vol. 9, p. 73.
[11]R. J. Protheroe and K. E. Turver, Nuovo Cimento 51A, 277 (1979).
[12]R. Walker and A. A. Watson, private communication.
[13]G. Cunningham *et al.*, Astrophys. J. 236, L71 (1980).
[14]W. V. Jones, in *Proceedings of 1978 Dumand Summer Workshop,* edited by A. Roberts (Fermilab, Batavia, Ill., 1979), Vol. 1, p. 313.
[15]R. A. Antonov *et al.*, in *Proceedings of the Sixteenth International Conference on Cosmic Rays, Kyoto, Japan, 1979* (University of Tokyo, Tokyo, Japan, 1979), Vol. 9, p. 258.
[16]This composition would arise if an extragalactic proton component was added to the mass distribution from galactic sources, which dominate at low energies.
[17]J. A. Goodman *et al.*, Phys. Rev. Lett. 42, 854, 1246(E) (1979).
[18]S. Thornton and R. Clay, Phys. Rev. Lett. 45, 1463(E) (1980).

Proton Decay Rate

Yukio Tomozawa
Randall Laboratory of Physics, University of Michigan, Ann Arbor, Michigan 48109
(Received 27 October 1980)

The pionic two-body decay amplitude of the proton in the SU(5) grand unified gauge theory is computed by using the soft-pion method in the reference frame of the pion at rest. The normalization condition of the relativistic three-body Bethe-Salpeter wave function of the proton is used in the computation. It is shown that the partial decay rate of the process, proton $\to e^+\pi^0$, is $(0.86\times 10^{30}\text{ yr})^{-1}$ and $(1.4\times 10^{31}\text{ yr})^{-1}$ for $m_X = 4.0\times 10^{14}$ GeV and 8.0×10^{14} GeV, respectively. These values are on the boundary of the present experimental limit.

PACS numbers: 13.30.Ce, 11.40.Ha, 12.20.Hx, 14.20.Ei

One of the most remarkable consequences of grand unified gauge theories is that the proton decays. Several authors have made estimates of the proton decay rate in the SU(5) and SO(10) models[1,2] and have found the lifetime to be in the range of $10^{31}-10^{33}$ yr,[3,4] which is close to the present experimental lower bound[5] ($\sim 10^{30}$ yr). Renewed attempts to observe such decays are under way and are expected to give some results soon if the proton lifetime is in the above-mentioned range.

Among various decay modes, the mode $p\to e^+ + \pi^0$ seems to be the most appropriate one for detection in most experiments in progress if its branching ratio is significantly large. However, there is some uncertainty in the theoretical estimates for this process since the estimates are made based on either the SU(6) wave functions or the Massachusetts Institute of Technology bag model wave functions, which are either nonrelativistic or noncovariant. The proton, consisting of three light quarks, is likely to be a relativistic system and the relativistic corrections may not be negligible. In this article, I present a relativistic calculation of the pionic two-body decay of protons based on the partially conserved axial-vector current hypothesis (PCAC) and current algebra.

In order to use the soft-pion method, I shall calculate the decay amplitude in the rest frame of the pion. If the soft-pion limit is taken literally, the momentum of the incident proton becomes infinite:

$$(p^2+M^2)^{1/2} = \frac{M^2+m^2+\mu_\pi^2}{2\mu_\pi}\to\infty, \text{ as } \mu_\pi\to 0, \qquad (1)$$

where M, m, and μ_π are the masses of the proton, the positron, and the pion, respectively. However, it is important to note that we can use the soft-pion method either in the ordinary reference frame of the pion at rest or in the infinite-momentum reference frame of the incident proton.[6] The two situations are identical in the limit $\mu_\pi\to 0$. In fact, this method has been used in the analysis of nonleptonic hyperon weak decays and has led to successful sum rules for the p-wave amplitudes as well as for the s-wave amplitudes.[7]

The interaction Lagrangian in the SU(5) gauge model, which is relevant to the proton decay, is given by[1]

$$\mathcal{L} = (4G\lambda/\sqrt{2})\epsilon_{ijk}[\bar{u}_{kL}{}^c\gamma_\mu u_{jL}(2\bar{e}_L{}^\dagger\gamma_\mu d_{iL} + \bar{e}_R{}^\dagger\gamma_\mu d_{iR} + \bar{\mu}_L{}^\dagger\gamma_\mu s_{iL} + \bar{\mu}_R{}^\dagger\gamma_\mu s_{iR}) + \bar{u}_{kL}{}^c\gamma_\mu d_{jL}(\bar{\nu}_{eR}{}^c\gamma_\mu d_{iR} + \bar{\nu}_{\mu R}{}^c\gamma_\mu s_{iR})], \quad (2)$$

where $(u_k, d_k, s_k)_{L(R)}$ and $(e^-, \nu_e, \mu^-, \nu_\mu)_{L(R)}$ are the left- (right-) handed quarks of color index k and the leptons, respectively. The coupling constant G is given in terms of the grand unified coupling constant $g_{\rm GUT}$ and the heavy gauge field mass m_X, by

$$G = (g_{\rm GUT}{}^2/4\sqrt{2}m_X{}^2), \quad (3)$$

while λ is the quantum-chromodynamic enhancement factor[1,3] ($\lambda = 3.7$).

Using the PCAC relation for pions,

$$\varphi^\alpha = (F_\pi\mu_\pi{}^2)^{-1}\partial A_\mu{}^\alpha/\partial x_\mu, \quad \alpha = 1, 2, 3, \quad (4)$$

where $\sqrt{2}F_\pi$ is the decay constant of $\pi^\pm \to \mu^\pm\nu_\mu$ ($\sqrt{2}F_\pi = 0.945\mu_\pi$) and where $A_\mu{}^\alpha$ is the axial-vector current, and applying the standard technique of the Lehmann-Symanzik-Zimmermann reduction formalism in the rest frame of the pion, I obtain the following expression for the decay amplitude for $p \to e^+\pi^0$:

$$(2q_0)^{1/2}\langle e(\vec{p}), \pi^0(\vec{0})|\mathcal{L}|p(\vec{p})\rangle = -i\,\frac{\mu_\pi{}^2 - q_0{}^2}{F_\pi\mu_\pi{}^2}\left\{\langle e|[Q_5{}^3, \mathcal{L}]|p\rangle - q_0\sum_l\left[\frac{\langle e|A_0{}^3(0)|l\rangle\langle l|\mathcal{L}|p\rangle}{q_0 + p_0{}^e - p_0{}^l + i\epsilon} - \frac{\langle e|\mathcal{L}|l\rangle\langle l|A_0{}^3(0)|p\rangle}{q_0 + p_0{}^l - p_0{}^p - i\epsilon}\right]_{\vec{p}_l = \vec{p}}\right\}. \quad (5)$$

The limit $q_0 \to 0$ selects the intermediate state l in the summation of Eq. (5) to be equal to either the initial or the final particle state. Thus, I obtain

$$f = (2q_0)^{1/2}\langle e\pi^0|\mathcal{L}|p\rangle = \frac{-i}{F_\pi}\{\langle e|[Q_5{}^3, \mathcal{L}]|p\rangle + \sum_{\substack{\text{intermediate}\\ \text{spin sum}}}\langle e|\mathcal{L}|p\rangle\langle p|A_0{}^3(0)|p\rangle\}. \quad (6)$$

Here I have used the identity $\langle e|A_0{}^3(0)|e\rangle = 0$. The first term of Eq. (6) is computed by using

$$Q_5{}^3 = \tfrac{1}{2}\int(u_i{}^\dagger\gamma_5 u_i - d_i{}^\dagger\gamma_5 d_i)d^3x, \quad (7)$$

giving,[8]

$$[Q_5{}^3, \mathcal{L}] = (4G\lambda/\sqrt{2})\epsilon_{ijk}\tfrac{1}{2}[\bar{u}_{kL}{}^c\gamma_\mu u_{jL}(2\bar{e}_L{}^\dagger\gamma_\mu d_{iL} - \bar{e}_R{}^\dagger\gamma_\mu d_{iR}) + \bar{u}_{kL}{}^c\gamma_\mu d_{jL}\bar{\nu}_{eR}{}^c\gamma_\mu d_{iR}]. \quad (8)$$

Then Eq. (6) becomes

$$f = \frac{-i}{F_\pi}\frac{4G\lambda}{\sqrt{2}}\epsilon_{ijk}\left(\frac{m}{\mathcal{E}_p}\right)^{1/2}\bar{u}_e{}^{(e)}(p)\langle 0|u_{ka}d_{ic}u_{jb}|p\rangle\left(c^{-1}\gamma_\mu\frac{1+\gamma_5}{2}\right)_{ab} \times\left[\frac{1}{2}\left(\gamma_\mu\frac{1+3\gamma_5}{2}\right)_{ec} + \left(\gamma_\mu\frac{3+\gamma_5}{2}\right)_{ec}\frac{M}{2E_p}g_A\bar{u}(p)\gamma_4\gamma_5 u(p)\right], \quad (9)$$

where $\mathcal{E}_p$, $u^{(e)}(p)$ and E_p, $u(p)$ are the energy and the Dirac spinor for the positron and the proton, respectively, and g_A is the axial-vector coupling constant of the neutron β decay.

In order to find the expression for the three-quark wave function of the proton, $\langle 0|u_{ka}(0)d_{ic}(0)u_{jb}(0)|p\rangle$, we define the three-body Bethe-Salpeter amplitude for the octet baryons,[9] B,

$$\langle 0|T\psi_{ia}{}^\alpha(x_1)\psi_{jb}{}^\beta(x_2)\psi_{kc}{}^\delta(x_3)|B\rangle = (M/E_p)^{1/2}\epsilon^{ijk}\tfrac{1}{2}(\chi_{abc}{}^{(\xi)}U_{\alpha\beta\gamma}{}^{(\xi)} + \chi_{abc}{}^{(\eta)}U_{\alpha\beta\gamma}{}^{(\eta)})\varphi(\xi, \eta, p)e^{ipX}, \quad (10)$$

where

$$X = \tfrac{1}{3}(x_1 + x_2 + x_3), \quad p = p_1 + p_2 + p_3, \quad \xi = x_1 - x_2,$$
$$p_\xi = \tfrac{1}{2}(p_1 - p_2), \quad \eta = \tfrac{1}{2}(x_1 + x_2 - 2x_3), \quad p_\eta = \tfrac{1}{3}(p_1 + p_2 - 2p_3), \quad (11)$$

$$\chi_{abc}{}^{(\xi)} = \left(\frac{-i\gamma p + M}{2M}\gamma_5 C\right)_{ab}u_c(p), \quad \chi_{abc}{}^{(\eta)} = (1/\sqrt{3})(\chi_{bca}{}^{(\xi)} - \chi_{cab}{}^{(\xi)}), \quad (12)$$

and

$$U_{\alpha\beta\gamma}^{(\xi)} = \epsilon_{\alpha\beta\gamma} B_\gamma^{\;\delta}, \quad U_{\alpha\beta\gamma}^{(\eta)} = (1/\sqrt{3})(U_{\beta\gamma\alpha}^{(\xi)} - U_{\gamma\alpha\beta}^{(\xi)}). \tag{13}$$

The Greek index in Eq. (10) represents the ordinary SU(3) index of the strong interactions (not the color index) and B stands for the 3×3 matrix of the baryon octet.

The normalization of the above amplitude can be written as

$$-iM\int \frac{d^4p_\xi' d^4p_\eta' d^4p_\xi d^4p_\eta}{(2\pi)^{16}} \bar{\chi}_p(p_\xi' p_\eta') \frac{\partial}{\partial p_0}[I(p_\xi', p_\eta', p_\xi, p_\eta; p) + \bar{G}(p_\xi', p_\eta', p_\xi, p_\eta; p)]\chi_p(p_\xi, p_\eta) = p_0, \tag{14}$$

where $\chi_p(p_\xi, p_\eta)$ is the Fourier transform of

$$\chi_p(\xi,\eta) = \tfrac{1}{2}\epsilon^{ijk}(\chi_{abc}^{(\xi)} U_{\alpha\beta\gamma}^{(\xi)} + \chi_{abc}^{(\eta)} U_{\alpha\beta\gamma}^{(\eta)})\varphi(\xi,\eta,p). \tag{15}$$

$\bar{G}$ is the kernel for the six-point Green's function and

$$I(p_\xi', p_\eta', p_\xi, p_\eta; p) = (2\pi)^8 \delta(p_\xi - p_\xi')\delta(p_\eta - p_\eta')[S_F^A(\tfrac{1}{3}p + p_\xi + \tfrac{1}{2}p_\eta) S_F^B(\tfrac{1}{3}p - p_\xi + \tfrac{1}{2}p_\eta) S_F^C(\tfrac{1}{3}p - p_\eta)]^{-1} \tag{16}$$

with

$$[S_F^A(k)]^{-1} = i(i\gamma k + m_A). \tag{17}$$

Equation (14) is a generalization of the normalization condition for the two-body Bethe-Salpeter amplitude.[10] The explicit computation of Eq. (10) gives

$$\int \frac{d^4p_\xi d^4p_\eta}{(2\pi)^8}|\tilde{\varphi}|^2\left[\frac{1}{2}\frac{(pp_\eta)^2}{M^2} + \frac{2}{3}\frac{(pp_\xi)^2}{M^2} - 2(\tfrac{1}{3}M - m_q)^2\right] = 1, \tag{18}$$

where $\tilde{\varphi}(p_\xi, p_\eta; p)$ is the Fourier transform of $\varphi(\xi,\eta;p)$. Then, for the wave function for the relativistic harmonic-oscillator potential,[9]

$$\varphi(\xi,\eta;p) = N\exp\left\{-\frac{\alpha}{6}\left[2\left(\frac{p\hat{\eta}}{M}\right)^2 + \hat{\eta} + 2\left(\frac{p\hat{\xi}}{M}\right)^2 + \hat{\xi}^2\right]\right\}, \tag{19}$$

where

$$\hat{\xi} = \xi/\sqrt{2}, \quad \hat{\eta} = (2/3)^{1/2}\eta, \tag{20}$$

the normalization condition gives

$$N = (\alpha/3\pi)^2[\alpha - 2(M - 3m_q)^2]^{-1/2}. \tag{21}$$

Assuming that $M \approx 3m_q$, I obtain

$$\varphi(0,0;p) \equiv N = (3\pi)^{-1/2}(\alpha/3\pi)^{3/2}. \tag{22}$$

Using Eqs. (10), (12), (13), and (19), I can compute the matrix element of Eq. (9), giving

$$f = \frac{-i}{F_\pi}\frac{4G\lambda}{\sqrt{2}} 6N\left(\frac{mM}{\mathcal{E}_p E_p}\right)^{1/2}\frac{1}{2}\left[\bar{u}^{(e)}(p)\Gamma_1 u(p) + \bar{u}^{(e)}(p)\Gamma_2 u(p)\frac{M}{2E_p} g_A \bar{u}(p)\gamma_4\gamma_5 u(p)\right], \tag{23}$$

where

$$\begin{aligned}\Gamma_i = {}& -\gamma_\mu \frac{a_i + b_i\gamma_5}{2} C\gamma_\mu^T \frac{-i\gamma^T p + M}{2M} C^{-1}\gamma_\mu \frac{1+\gamma_5}{2} + \tfrac{1}{3}\gamma_\mu \frac{a_i + b_i\gamma_5}{2}\frac{-i\gamma p + M}{2M}\gamma_5 C \frac{1+\gamma_5^T}{2}\gamma_\mu^T C^{-1} \\ & + \tfrac{1}{3}\gamma_\mu \frac{a_i + b_i\gamma_5}{2}\mathrm{Tr}\left(\frac{-i\gamma p + M}{2M}\gamma_5\gamma_\mu\frac{1+\gamma_5}{2}\right) \\ = {}& -b_i + a_i\gamma_5 + \frac{a_i + b_i}{3}(1-\gamma_5)\frac{i\gamma p + M}{M} \to -b_i + a_i\gamma_5, \quad i = 1,2.\end{aligned} \tag{24}$$

Using the values

$$a_1 = \tfrac{1}{2}, \quad b_1 = \tfrac{3}{2}, \quad a_2 = 3, \quad b_2 = 1, \tag{25}$$

I obtain

$$\Gamma_1 = -\tfrac{3}{2} + \tfrac{1}{2}\gamma_5, \quad \Gamma_2 = -1 + 3\gamma_5. \tag{26}$$

The bilinear Dirac spinors in the rest frame of the pion which appear in Eq. (23) can be computed by using the formulas

$$\bar{u}(p)\gamma_4\gamma_5 u(p) = -\frac{|p|}{M}\xi^\dagger\vec{\sigma}\cdot\hat{\vec{p}}\xi, \quad \bar{u}^{(e)}(p)u(p) = \tfrac{1}{2}(Mm)^{-1/2}[(M+m)^2-\mu_\pi^{\ 2}]^{1/2}\xi^{(e)\dagger}\xi,$$
$$\bar{u}^{(e)}(p)\gamma_5 u(p) = \tfrac{1}{2}(Mm)^{-1/2}[(M-m)^2-\mu_\pi^{\ 2}]^{1/2}\xi^{(e)\dagger}\vec{\sigma}\cdot\hat{\vec{p}}\xi, \tag{27}$$

where ξ and $\xi^{(e)}$ are the two-component spinor of proton and positron, respectively, and $\hat{\vec{p}}$ is the unit vector in the direction of $\vec{p}$. Alternatively, one can use Table III of Ref. 7 for the infinite-momentum reference frame of the incident proton. Both methods give the same result upon neglecting the pion mass. I get the final form for the decay amplitude

$$\frac{f}{(2q_0)^{1/2}} = \frac{-i}{F_\pi}\frac{1}{(2q_0)^{1/2}}\frac{12G\lambda N}{\sqrt{2}}\left(\frac{mM}{\mathcal{E}_p E_p}\right)^{1/2}\bar{u}^{(e)}(p)(A+B\gamma_5)u(p), \tag{28}$$

where

$$A = -\tfrac{3}{2}\{1+[(M-m)/(M+m)]g_A\} \cong -\tfrac{3}{2}(1+g_A) = -3.39,$$
$$B = \tfrac{1}{2}\{1+[(M+m)/(M-m)]g_A\} \cong \tfrac{1}{2}(1+g_A) = 1.13. \tag{29}$$

(Here, use is made of the empirical value, $g_A = 1.260 \pm 0.007$.)

The decay rate of $p \to e^+\pi^0$, then, can be computed from Eq. (28), giving

$$\Gamma(p\to e^+\pi^0) = \frac{M}{16\pi}\left(\frac{12G\lambda N}{\sqrt{2}F_\pi}\right)^2(|A|^2+|B|^2) = \frac{M}{16\pi}\left(\frac{3g_{\rm GUT}^{\ 2}\lambda N}{2F_\pi m_X^{\ 2}}\right)^2\frac{5}{2}(1+g_A)^2. \tag{30}$$

For typical values of the parameters[4,9,11]

$$g_{\rm GUT}^{\ 2}/4\pi = 0.024, \quad \lambda = 3.7, \quad \alpha = 0.4\ \mathrm{GeV}^2, \tag{31}$$

I obtain the partial decay time

$$\tau(p\to e^+\pi^0) = [\Gamma(p\to e^+\pi^0)]^{-1} = \begin{cases} 0.86\times 10^{30}\ \mathrm{yr\ for}\ m_X = 4.0\times 10^{14}\ \mathrm{GeV} \\ 1.4\times 10^{31}\ \mathrm{yr\ for}\ m_X = 8.0\times 10^{14}\ \mathrm{GeV}. \end{cases} \tag{32}$$

These values are on the boundary of the experimental limit,[5] and are reasonably close to those obtained by Gavela *et al.*[12] $\{[\Gamma(p\to e^+\pi^0)]^{-1} = 0.53\times 10^{30}$ y for $m_X = 4\times 10^{14}$ GeV$\}$. The results obtained by other groups may be compared with Eq. (32) by assuming a typical branching ratio[4,12-14] $\Gamma(p\to e^+\pi^0)/\Gamma(p\to \text{two bodies}) = \frac{1}{3}$; Goldman and Ross (Ref. 3), Din, Girardi, and Sorba (Ref. 13), Buras *et al.* (Ref. 1), and Donoghue (Ref. 4) give, respectively $10^{-30}[\Gamma(p\to e^+\pi^0)]^{-1} = 1.8$, 2.7, 3.7, and 30 yr for $m_X = 4\times 10^{14}$ GeV. For a comprehensive review and many references on this subject, see Ref. 14. Note also that if we take the value of the parameter α of the proton wave function to be 0.5 GeV^2, which is an alternative choice suggested in Ref. 9, we would have obtained $[\Gamma(p\to e^+\pi^0)]^{-1} = 0.44\times 10^{30}$ yr for $m_X = 4.0\times 10^{14}$ GeV. This is very close indeed to the value obtained by Gavela *et al.*[12]

From Eqs. (28) and (29), we can deduce that the decay amplitude of $p\to e^+\pi^0$ is predominantly s wave and the asymmetry parameter is given by

$$\alpha = 2\,\mathrm{Re}(A^*B)/(|A|^2+|B|^2) = -0.60. \tag{33}$$

It is a pleasure to thank David Unger for useful comments and for reading the manuscript. The work is supported in part by the U. S. Department of Energy.

[1]H. Georgi and S. L. Glashow, Phys. Rev. Lett. 32, 438 (1974); H. Georgi, H. R. Quinn, and S. Weinberg, Phys. Rev. Lett. 33, 451 (1974); A. J. Buras, J. Ellis, M. K. Gaillard, and D. V. Nanopoulos, Nucl. Phys. B135, 66 (1978).

[2]H. Georgi and D. V. Nanopoulos, Phys. Lett. 82B, 392 (1979), and Nucl. Phys. B155, 52 (1979).

[3]T. J. Goldman and D. A. Ross, Phys. Lett. 84B, 208 (1979), and Nucl. Phys. B171, 273 (1980); F. Wilczek and A. Zee, Phys. Rev. Lett. 43, 1571 (1979); L. Hill, Harvard University Report No. HUTP-80/A024, 1980 (to be published).

[4]M. Machacek, Nucl. Phys. B159, 37 (1979); J. F. Donoghue, CTP 824 (1979); G. Kane and G. Karl, Phys. Rev. D 22, 2808 (1980); G. Karl and H. J. Lipkin, FNAL Report No. 80/54-THY, 1980 (to be published).

[5]J. Learned, F. Reines, and A. Soni, Phys. Rev. Lett. 43, 907, 1626(E) (1979).

[6]Of course, the amplitude are Lorentz scalars and their values should not depend on the reference frame chosen. However, approximations which are made could be more effective in some reference frames than

in others.
[7]Y. Tomozawa, Phys. Lett. 32B, 485 (1970); M. Machacek and Y. Tomozawa, Phys. Rev. D 12, 3711 (1975); Y. Tomozawa, University of Michigan Report No. UM-HE-80-20, 1980 (to be published).
[8]Notice the relations, $[\int u_i^\dagger \gamma_5 u_i d^3x, u_{jL\,(jR)}] = (\mp) u_{jL\,(jR)}$, $[\int u_i^\dagger \gamma_5 u_i d^3x, \bar{u}_{jL\,(jR)}{}^c] = (\pm)\bar{u}_{jL\,(jR)}{}^c$, etc., where $u^c = C\bar{u}^T$, $\bar{u}^c = -u^T C^{-1}$, $CC^\dagger = 1$, $C^T = -C$, and $C^{-1}\gamma_\mu C = \gamma_\mu{}^T$.
[9]R. P. Feynman, M. Kislinger, and F. Ravndal, Phys. Rev. D 3, 2706 (1971); R. Lipes, Phys. Rev. D 5, 2849 (1972).
[10]D. Lurié, *Particle and Fields* (Wiley, New York, 1968), p. 433; N. Nakanishi, Prog. Theor. Phys. Suppl. 43, 1 (1969), and earlier references quoted therein.
[11]P. Binétruy and T. Schücker, CERN Report No. CERN-TH-2857, 1980 (to be published).
[12]M. G. Gavela *et al.*, Laboratoire de Physique Théorique et des Particules Elémentaires, Orsay, Report No. LPTHE-80/15, 1980 (to be published).
[13]A. M. Din, G. Girardi, and P. Sorba, Phys. Lett. 91B, 77 (1980).
[14]J. Ellis, M. K. Gaillard, and D. V. Nanopoulos, CERN Report No. CERN-TH-2833, 1980 (to be published).

Analysis of Reactor Experiments for Neutrino Oscillations

D. Silverman
Department of Physics, University of California, Irvine, California 92717

and

A. Soni
Department of Physics, University of California, Los Angeles, California 90024

(Received 17 November 1980)

Data from four reactor experiments is analyzed without using any calculated $\bar{\nu}_e$ spectrum. N_{e1} and N_{e2}, for e^+ observed with $2.2 < E_{e1} < 6.7$ MeV and $4.4 < E_{e2} < 6.7$ MeV, are extracted and N_{e1}/N_{e2} is found to be 2.7 ± 0.5, 5.6 ± 0.6, and 8.20 ± 0.35 for the 6.5-, 8.7-, and 11.2-m experiments, respectively. In pairs, these numbers differ by 3–8 standard deviations. No distance-independent $\bar{\nu}_e$ spectrum accounts for all the data with a confidence level (C.L.) > 0.0028. Oscillations with three (two) ν's yield fits to all data with C.L. $= 0.061$ (0.033) and to the high-statistics experiments with C.L. $\gtrsim 0.31$ (0.18).

PACS numbers: 14.60.Gh, 13.15.+g

Since the phenomena of neutrino oscillations was first discussed[1] there have been several experimental suggestions in support of that possibility.[2,3] The recent round of discussions on this subject was intensified by the experimental findings of Reines, Sobel, and Pasierb (RSP), who measured the rates for neutral current deuteron (ncd) and charge current deuteron (ccd) reactions initiated by reactor $\bar{\nu}_e$.[3] Over the years, the energy spectrum of reactor $\bar{\nu}_e$ has been experimentally measured by the inverse beta (IB) reaction $\bar{\nu}_e + p \to n + e^+$ at 6.5,[4] 8.7,[5] and 11.2 m (Ref. 6) from reactor sources. We shall study the e^+ energy spectra measured in those three experiments in conjunction with the deuteron experiment of RSP to examine the hypothesis of oscillations.

In the IB reaction the differential rate for e^+ with observed kinetic energy E_e at a distance L from a reactor source is given by

$$\frac{dR}{dE_e} = 0.203 \times (9.24 \times 10^{-44}\ \mathrm{cm}^2)^{-1} \times \left(\frac{P}{1\ \mathrm{MW}}\right)\left(\frac{n_p}{10^{26}}\right)\left(\frac{L}{1\ \mathrm{m}}\right)^{-2} \eta_s \int \sigma(E_\nu) R_e(E_e, E_e') \eta(E_e') n(E_\nu, L) dE_\nu\ (\mathrm{MeV\ d})^{-1}, \quad (1)$$

where n_p is the number of protons in the target, P is the reactor power, $E_e' = E_\nu - 1.8$ MeV, $\sigma(E_\nu) = 9.24 \times 10^{-44}(E_\nu - 1.29)[(E_\nu - 1.29)^2 - 0.26]^{1/2}$ cm², $R_e(E_e, E_e')$ is the experimental energy resolution function, $\eta(E_e')$ is the energy-dependent detection efficiency, η_s is the energy-independent systematic efficiency, and $n(E_\nu, L)$ is the spectrum (number of $\bar{\nu}_e$'s per fission per megaelectronvolt) of $\bar{\nu}_e$ with