ASTROPHYSICAL CONSEQUENCES OF THE EXISTENCE OF CHARGED INTERMEDIATE VECTOR BOSONS

By J. A. CAMPBELL*

[*Manuscript received November* 16, 1967]

Summary

Two-photon annihilation into a neutrino–antineutrino pair, which is forbidden to the lowest order in the coupling constant for weak interactions if the conventional form of weak interaction is assumed, is permitted at that order if it is supposed that the reaction is carried by a charged intermediate vector boson (W). The rate of loss of energy in stellar evolution through this process is calculated. Evolutionary time scales derived with and without this rate are compared with results from astronomical observations. Comparisons with present data are inconclusive, but further observations and calculations that may give more accurate information concerning the question of the existence of charged W mesons are suggested.

I. INTRODUCTION

The present status of weak interactions in elementary particle physics differs from that of the strong and electromagnetic interactions, in that experimental results for weak processes involving four fermions can still be treated adequately by a theory in which the four particles are assumed to interact strictly locally, at one point. By contrast, the simplest pictures of electromagnetic and strong interactions require the exchange of one photon and one pion respectively between the interacting particles. Each of these processes is therefore nonlocal, having a scale of range equal to the Compton wavelength of the exchanged particle. As a means of bringing weak-interaction theory into line with theories for the stronger interactions, Lee and Yang (1960*a*) have proposed the existence of a particle, the W meson or intermediate boson, which can be exchanged between weakly interacting particles. To account for the weak interactions that have been observed, it is necessary to postulate only two states of charge, W^+ and W^-, and to suppose that this intermediate boson is a vector particle, having a spin of 1. Varieties of neutral W mesons have been suggested (Lee and Yang 1960*b*; Marshak 1966) on additional theoretical grounds, but in the present paper only the implications of the original proposal are considered.

Since the range of an interaction is inversely proportional, through the Compton wavelength, to the mass of the exchanged particle, a strictly local interaction is one in which a particle of infinite mass is "exchanged". Therefore any theoretical description of a weak process involving W mesons must reproduce exactly the results of the point-interaction theory in the limit that the mass of the W meson tends

* School of Physical Sciences, Flinders University of South Australia, Bedford Park, S.A. 5042.

to infinity. This requirement provides a condition of consistency on the nonlocal theory, and in practice it relates the square of the coupling constant of W mesons with other weakly interacting particles to the known coupling constant for beta decay.

Laboratory experiments that seek to observe W production by proton beams explore progressively shorter ranges of interaction as the proton energy increases. If the W meson is not observed when the experiment measures a given transfer of four-momentum, a lower limit for the W mass can therefore be set. It has been reported, by means of an argument of this type, that the mass must be greater than 2 GeV (Burns *et al.* 1965) and that it is unlikely that a W meson with mass between 2 and 3 GeV exists (Lamb *et al.* 1965). Recently Glashow, Schnitzer, and Weinberg (1967) have found that certain amplitudes for the decays of K mesons, which are divergent if the local weak interaction is used, become convergent if a W meson is assumed to take part in the decays. The convergent theoretical calculation contains the W mass as a free parameter. From a comparison with the experimentally determined rate of decay of the state K_1^0, they conclude that the mass is approximately 8 GeV. This value is exceptionally large by the standards of the elementary particles that have been detected in past experiments.

Without positive evidence of the W meson (which is of crucial interest for the structure of weak interaction theory) from the laboratory, one may still turn to astrophysics to see what effects the existence of the particle implies in stellar evolution and to compare these effects with what is observed astronomically. Moreover, since the numerical limits quoted above are appreciably higher than the mass of the heaviest meson detected to date, any astrophysical test that can in principle distinguish merely between the existence and non-existence of the W meson (and not make stronger statements about limits in mass) is welcome. In Section II are listed several neutrino-producing reactions that are the most efficient mechanisms for loss of energy from stars at central temperatures above about 5×10^8 °K. One of these reactions, neutrino–antineutrino production by the annihilation of pairs of photons, depends critically on the existence of the W meson. Quantitative estimates of the effect of that reaction in stars are made in Section III, and are compared with the consequences of the other neutrino processes in Section IV.

II. Astrophysical Neutrino Processes and W Mesons

The present theoretical description of weak interactions is based in effect on the current–current hypothesis of interaction put forward by Feynman and Gell-Mann (1958). In this theory, a vertex of interaction is predicted at which an electron and a positron can annihilate to give an electron neutrino and an electron antineutrino. That interaction has not yet been seen in the laboratory, but Chin, Chiu, and Stothers (1966) conclude from astrophysical evidence that the vertex actually occurs, and that its coupling constant is probably about the same as the coupling constant for other observed vertices permitted by the theory of Feynman and Gell-Mann. That value for the coupling constant is assumed in equation (4) below.

Existence of the $(e\nu_e)(e\nu_e)$ local interaction allows three astrophysically important reactions involving neutrinos to take place: the pair-annihilation process,

$$e^+ + e^- \rightarrow \nu_e + \bar{\nu}_e\,, \tag{1}$$

the photoneutrino process,

$$\gamma + e^- \rightarrow \nu_e + \bar{\nu}_e + e^-\,, \tag{2}$$

and the plasma process,

$$\gamma \rightarrow e^+ + e^- \rightarrow \nu_e + \bar{\nu}_e\,. \tag{3}$$

The process (3) is forbidden by gauge invariance for a free photon, but is permitted in the degenerate electron gas that occurs in the core of a star at some stages of stellar evolution. There, the symmetrical disposition of filled and unfilled electron states about an energy of zero in the basic Dirac theory is removed by the filling of all states between an energy of $+mc^2$ (where m is the mass of the electron and c the velocity of light) and a Fermi energy E_F. Photons propagating through the electron gas are not free photons and the process (3) takes place.

Chiu (1966) has summarized the expected rates of loss of energy in stars for each of the three processes listed above. Generally, the summary shows that the plasma process is the most important mechanism for loss of energy in degenerate material, while in nondegenerate matter the photoneutrino process is most important below about 7×10^8 °K and pair annihilation is the dominant process above that temperature.

If the coupling constant for interactions of W mesons with leptons is f, the condition of consistency (mentioned in Section I) on the weak interaction in the limit of a strictly local four-fermion process implies (Lee and Wu 1965) that:

$$f^2/W^2 = G/\sqrt{2}\,, \tag{4}$$

where W is the mass of the W meson and G is the coupling constant for beta decay. A practical consequence of equation (4) is that, if the cross section σ_W for a weak four-fermion reaction mediated by a W meson and the cross section σ_0 for the corresponding local reaction are calculated separately, they can be related by the equation

$$\sigma_W \approx \sigma_0(1+2q^2/W^2)\,, \tag{5}$$

where q is the four-momentum transferred in the process and terms containing higher powers of q^2/W^2 have been neglected. Since an energy of 1 MeV is associated with a temperature of $1\cdot16\times10^{10}$ °K (temperatures reached in stars during ordinary stellar evolution scarcely attain this value) and the experimental lower limit quoted for W in Section I is 2 GeV, it is not possible to measure any significant astrophysical differences between σ_W and σ_0 for the processes (1), (2), and (3).

However, there is a fourth process,

$$\gamma + \gamma \rightarrow \nu + \bar{\nu}\,, \tag{6}$$

which is not subject to the limitations of equation (5). Both σ_W and σ_0 for the process of two-photon annihilation can be calculated by an application of the conventional rules that relate Feynman diagrams to matrix elements, but Gell-Mann (1961) has shown separately, in the lowest order of G, that reaction (6) cannot proceed through a local interaction. His argument uses invariance under charge conjugation, the inability of a pseudovector particle to decay into two photons, and the fact that $\gamma.p\, v(p)$ and $\bar{v}(p)\,\gamma.p$ are zero from the Dirac equation, where $v(p)$ is an antineutrino spinor and γ represents a Dirac matrix. Nevertheless, as Gell-Mann himself has pointed out, this argument does not hold if the reaction is mediated by a W meson, as in the Feynman diagram in Figure 1. Therefore the presence or absence of the process (6) in stellar evolution depends on the existence of charged W mesons. The question of their existence may be examined by comparisons of predictions about stellar evolution with and without (6) to actual astronomical observations.

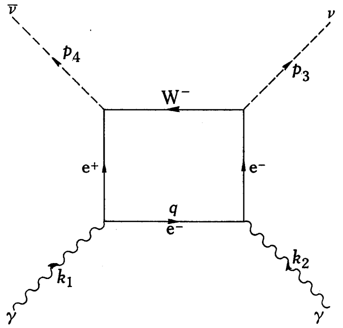

Fig. 1.—Feynman diagram for the process of two-photon annihilation, with the inclusion of a virtual charged W meson.

III. Quantitative Results for the Two-photon Reaction

For the reaction illustrated by the Feynman diagram in Figure 1, the matrix element is

$$\frac{\mathrm{i}}{(2\pi)^6}\frac{e^2f^2}{(16E_1E_2E_3E_4)^{\frac{1}{2}}}\,\delta(p_3+p_4-k_1-k_2)$$
$$\times\int \mathrm{d}^4q\,\bar{v}(p_4)\,\gamma_\mu(1+\gamma_5)\frac{\gamma.(k_1-q)+m}{q^2-2k_1.q-m^2}\,\gamma.e_1\frac{\gamma.q+m}{q^2-m^2}\,\gamma.e_2\frac{\gamma.(k_2+q)+m}{q^2+2k_2.q-m^2}$$
$$\times\gamma_\nu(1+\gamma_5)\,v(p_3)\,\frac{\delta_{\mu\nu}-W^{-2}(q+k_2-p_3)_\mu\,(q+k_2-p_3)_\nu}{q^2+2(k_2-p_3).q-W^2}. \tag{7}$$

In (7), which results from a straightforward application of Feynman's rules, e_1 and e_2 are polarization four-vectors for the photons, δ is the four-dimensional Dirac delta function, and natural units with $\hbar = c = 1$ are used. Diagrams with γWW vertices are ignored because of the extra factors of W^{-2} that they contribute to (7). For the parts of (7) in which five factors of q occur in the numerator of the integrand, the integral over q is zero because a function that is odd about $q = 0$ is involved. In each of the five possible combinations of four factors of q, the integral is logarithmically divergent and may be expressed in terms of a cutoff Λ with the dimension of a mass. The introduction of a cutoff is somewhat artificial, but it is usually justified by the statement that Λc^2 is the highest energy for which the theory is valid. In this example, uncertainties introduced into the calculation by the different possible numerical choices of Λ are not important on an astrophysical scale.

The dominant terms in the cross section derived from (7) contain the logarithmic dependence on Λ. If only these terms are retained, scalar products of q with four-momenta for external lines of Figure 1 are neglected in the denominator of (7) (because they are considerably smaller than W^2 for small q, and smaller than q^2 for large q), and factors of $\hbar$ and c are re-introduced, the cross section σ for the process is given by

$$\sigma \approx \frac{811G^2M^4}{262\,144}\frac{\alpha^2}{\pi^5}\left(\frac{11}{6}+\log\frac{\Lambda^2}{W^2}\right)^2\left(\frac{\hbar}{mc}\right)^2\left(\frac{m}{M}\right)^4\frac{E_1E_2}{(mc^2)^2} \qquad \mathrm{cm}^2, \tag{8}$$

where α is the fine-structure constant and M is the mass of the proton. To obtain equation (8), sums over final-state quantities and averages over all initial-state quantities except the photon energies E_1 and E_2 have been performed. Further, the coupling constant f has been replaced in terms of G from equation (4). The derivation of equation (8) has been carried out by computer programmes (Hearn 1966; Campbell 1967) designed for such an application, and checked by hand.

The cross section in equation (8) is an integral of a differential cross section with respect to the momenta $\mathbf{p}_3$ and $\mathbf{p}_4$. The energy lost per event of the type (6) can be calculated if the differential cross section is multiplied by (E_3+E_4) prior to the integration. Then, if this energy loss is represented by $E(\mathbf{p}_1, \mathbf{p}_2)$, the total rate of loss of energy through (6) in a star is given by the integral of $E(\mathbf{p}_1, \mathbf{p}_2)$ with respect to the differentials of the number densities of the particles with momenta $\mathbf{p}_1$ and $\mathbf{p}_2$. In general, it is possible to write (Chiu and Stabler 1961)

$$-\rho\,\mathrm{d}E/\mathrm{d}t = \int |\mathbf{v}|\,E(\mathbf{p}_1,\mathbf{p}_2)\,\mathrm{d}n(\mathbf{p}_1)\,\mathrm{d}n(\mathbf{p}_2)\,, \tag{9}$$

where $\mathbf{v}$ is the relative velocity of particles 1 and 2 and ρ is the density of material in which the reaction takes place. Because both particles are photons, $|\mathbf{v}| = c$ and the number densities refer to particles that obey Bose–Einstein statistics. For dimensional consistency in equation (9), dn must have the dimension of a number per unit volume. The correct substitution for inclusion in equation (9) is therefore

$$\mathrm{d}n(\mathbf{p}) = \mathrm{d}n(E) = \{\pi^2(\hbar c)^3\}^{-1}\{\mathrm{e}^{E/kT}-1\}^{-1}E^2\,\mathrm{d}E\,. \tag{10}$$

The neutrinos in Figure 1 and the process (6) may equally well be muon neutrinos as electron neutrinos, provided that the electron and positron lines in Figure 1 are replaced by negative and positive muon lines in the first case. In this respect, the process (6) is unlike the reactions (1), (2), and (3), which involve only electrons and positrons, because stellar temperatures below about $1\cdot2\times10^{12}$ °K are insufficient for the production of appreciable concentrations of muons. Muons may still make virtual contributions to (6) at lower temperatures by their presence as internal lines in Figure 1. Admittedly electrons and positrons occur only as internal lines in the Feynman diagram for the plasma process (3), but these virtual particles accompanying photon propagation come from the degenerate electron gas that supports it. There is no corresponding "muon gas". Therefore, in the evaluation of equation (9), $E(\mathbf{p}_1, \mathbf{p}_2)$ should include contributions from both the processes

$$\gamma+\gamma \rightarrow \nu_{\mathrm{e}}+\bar{\nu}_{\mathrm{e}} \qquad \text{and} \qquad \gamma+\gamma \rightarrow \nu_\mu+\bar{\nu}_\mu.$$

Both contributions are covered adequately if the differential cross section that leads to equation (8) is multiplied by a factor of two.

The expression for the rate of loss of energy when equation (10) is substituted into equation (9) is then

$$-\frac{\mathrm{d}E}{\mathrm{d}t} = \frac{811}{131\,072}\left(\frac{kT}{\pi c}\right)^9 \frac{G^2\alpha^2}{\hbar^4\rho}\left(\frac{11}{6}+\log\frac{\Lambda^2}{W^2}\right)^2 \int_0^\infty \mathrm{d}x \int_0^\infty \mathrm{d}y\, \frac{(x+y)x^3y^3}{(\mathrm{e}^x-1)(\mathrm{e}^y-1)}$$

$$= \frac{811}{40\,960}\left(\frac{kT}{c}\right)^9 \frac{G^2\alpha^2}{\pi^5\hbar^4\rho}\left(\frac{11}{6}+\log\frac{\Lambda^2}{W^2}\right)^2 \zeta(5) \qquad \mathrm{ergs\,g^{-1}\,sec^{-1}}, \tag{11}$$

where k is Boltzmann's constant and ζ is the Riemann zeta function. If it is supposed that $\Lambda \approx 50W$ in equation (11), the rate of loss in terms of the two variables ρ and T_9 ($= T/10^9$) is

$$-\mathrm{d}E/\mathrm{d}t \approx 2\cdot5\times10^9\, T_9^9/\rho \qquad \mathrm{ergs\,g^{-1}\,sec^{-1}}. \tag{12}$$

This rate is independent of both the degeneracy of the material in which it occurs and any approximations concerning nonrelativistic or extreme-relativistic kinematics. Moreover, it agrees with the rate calculated approximately by Boccaletti, de Sabbata, and Gualdi (1964) for annihilation of pairs of optical photons.

Chiu and Morrison (1960) have estimated losses of energy in the additional astrophysical process $\gamma+\gamma \rightarrow \gamma+\nu+\bar{\nu}$. If the subscript B is used for this reaction and A for the process (6), an order-of-magnitude calculation based on equations (8) and (11) and Feynman diagrams suggests that

$$\left(-\frac{\mathrm{d}E}{\mathrm{d}t}\right)_{\mathrm{B}} < \frac{2\pi\alpha(4a+1)^2}{9a^2}\left(-\frac{\mathrm{d}E}{\mathrm{d}t}\right)_{\mathrm{A}}, \tag{13}$$

where

$$a = 11/6 + \log(\Lambda^2/W^2).$$

The inequality (13) implies an insignificant rate of loss of energy for the reaction B, which is not examined further. The results of the examination of a three-photon astrophysical process by de Graaf and Tolhoek (1966) also satisfy the inequality.

IV. Discussion

Rates of loss of energy for the processes (1), (2), and (3) have been presented collectively, in forms suitable for comparison with equation (12), by Chiu and Stabler (1961), Zaidi (1965), and Chiu (1966). Provided that charged W mesons exist, examination of these results shows that, as a medium of loss of energy in stars, the process of two-photon annihilation is competitive with:

(i) The plasma process, when the quantity $\hbar\omega_{\mathrm{p}}/kT$ (where ω_{p} is the plasma frequency) is either less than $0\cdot01$ or greater than 50. The first case corresponds to an extremely high temperature, and the second to an extremely high degree of degeneracy. Although either of these conditions may be realised in an exceptional situation such as the gravitational collapse of a star, it is unlikely that they can be the subjects of widespread observation.

(ii) All photoneutrino processes except those taking place in degenerate matter at nonrelativistic energies.

(iii) Pair annihilation in degenerate matter, provided either that the energies involved are nonrelativistic or that $T_9 \geqslant 3$. The definition of degeneracy used in this example is that the Fermi energy E_F should be greater than about $5mc^2$.

Chiu (1966) and Chin, Chiu, and Stothers (1966) have considered altogether six situations in astrophysics in which neutrino production may have a significant influence. Of these, the behaviour of intergalactic neutrinos is not notably affected by the reaction (6) because of low temperature, and the evolution of degenerate white dwarf stars and semidegenerate "ultraviolet" dwarf stars is believed to depend only on details of the plasma process. Two-photon annihilation may be expected to have an important place in models describing supernovae or gravitational collapse, but in these cases there is no chance that astronomers can make a sufficient number of observations to provide indirect evidence of the existence of charged W mesons.

In the two remaining examples, counts of numbers N of stars in different regions of colour–magnitude diagrams are used to determine characteristic evolutionary time scales for the different regions, according to relations of the form

$$\tau_A/\tau_B = N_A/N_B \tag{14}$$

Overall normalization of times is provided by relevant models that describe stellar evolution. The time scale for any one region is connected with the output of energy from a star by the relation

$$\tau \propto \int \left(\sum_i L_i(T)\right)^{-1} \mathrm{d}T\,, \tag{15}$$

where each $L_i(T)$ is the luminosity contributed by one mechanism for loss of energy which assists the evolution of the star through that region. Luminosity is given in terms of the rate of loss of energy per unit volume by

$$\mathrm{d}L_i/\mathrm{d}r = 4\pi r^2(-\mathrm{d}E/\mathrm{d}t)_i\,. \tag{16}$$

The rate on the right-hand side of equation (16) is a function of r to the extent that stellar models usually divide a star into a core and a small number of outer shells, the material in each section having a different temperature, density, and chemical composition. The integral over r in equation (15) can thus be written as a sum of integrals over different ranges of r, $\mathrm{d}E/\mathrm{d}t$ being constant within each range.

Central stars of planetary nebulae occupy a well-defined region on a colour–magnitude diagram. Calculations (Seaton 1966) based on observations by O'Dell (1963) indicate that the time scale for their contraction to the limiting or Chandrasekhar radius lies between 3×10^4 and 5×10^4 years. Hayashi, Hōshi, and Sugimoto (1962) have calculated that model central stars of 0·6 and 0·4 solar masses go through the contraction in 10^6 years if neutrino processes are not taken into account. From other model calculations, Chin, Chiu, and Stothers (1966) find that this time is reduced to about 10^4 years if the effect of the photoneutrino process is included.

An approximate analysis based on equations (15) and (16) suggests that their estimate is not reduced by more than 5% for central stars with masses less than one solar mass, but that the difference may be increased in heavier stars. It is therefore desirable to have much more astronomical information on the heavier central stars of planetary nebulae to clarify this point.

The most promising situation for discussion concerns the evolution of stars through the stages of helium and carbon burning in their cores after they have left the main sequence on a colour–magnitude diagram. It should be possible to make observations on many more stars in this situation than in any other of the cases mentioned above. With the help of a model for the evolution of stars of $15 \cdot 6$ solar masses, Hayashi and Cameron (1962) have analysed a diagram prepared from observations by Johnson and Morgan (1955) on the star clusters h and χ Persei. Time scales for evolution through various regions of the diagram are computed from the model and compared with "experimental" times obtained from equation (14). The scales disagree only for the region occupied by stars at the stage of carbon burning in their cores. With limited data, the measured time for this stage is between 7×10^5 and 10^6 years, while the theoretical time is $2 \cdot 5 \times 10^5$ years when neutrino processes are neglected. However, the disagreement becomes worse if neutrino production is also considered. During carbon burning, cores are nondegenerate and at temperatures where the photoneutrino process (2) is important. With the luminosity (Chiu and Stabler 1961) for this process included in equation (15), the time becomes 2×10^4 years (Chiu 1966).

Let $\tau_{12 \ldots n}$ be the time scale calculated through equation (15) when luminosities for n different reactions are considered. Let $L_1(T)$ be the luminosity obtained from the burning of ^{12}C in nuclear reactions, $L_2(T)$ be the luminosity accompanying the burning of the oxygen isotope ^{16}O (which is also present in stellar cores after the end of the helium-burning stage), $L_3(T)$ be the luminosity for the photoneutrino process, and $L_4(T)$ be the luminosity obtained from equations (12) and (16) for the two-photon reaction (6). Then $\tau_{12} = 2 \cdot 5 \times 10^5$ years and $\tau_{123} = 2 \times 10^4$ years.

It has been suggested that the mismatch between τ_{12} and the observed result is due to uncertainties about the relative concentrations of ^{12}C and ^{16}O present at the beginning of the carbon-burning stage. The concentration of carbon (and therefore the overall luminosity) decreases as the rate for the reaction

$$^{12}C + \alpha \rightarrow {}^{16}O + \gamma \tag{17}$$

increases, which allows τ_{12} to increase. In the present work a calculation of Hayashi, Hōshi, and Sugimoto (1962) on the rates of reactions contributing to $L_1(T)$ and $L_2(T)$ has been repeated with different reduced widths for the nuclear reaction (17), and a match between τ_{12} and the result of observations has been found for reduced widths in the range from $0 \cdot 23$ to $0 \cdot 32$. This adjustment is nevertheless not very plausible, because the most recent experimental information about the reduced width (Loebenstein *et al.* 1967) is that it lies between $0 \cdot 06$ and $0 \cdot 14$, in agreement with a detailed theoretical estimate by Stephenson (1966). It may be possible to obtain values inside the experimental range if the square-well potential used by Hayashi, Hōshi, and Sugimoto is replaced by a more realistic optical potential, as is implied by Stephenson and by Vogt, Michaud, and Reeves (1965). However, $L_3(T)$ is notably larger than either $L_1(T)$ or $L_2(T)$ at the temperatures of interest, so that no new description of the reaction (17) is likely to alter greatly the result that $\tau_{123} = 2 \times 10^4$ years.

When $L_4(T)$ is included in equation (15), the calculated time scale for carbon burning in stars in the clusters h and χ Persei becomes $\tau_{1234} = 1 \cdot 75 \times 10^4$ years. Therefore at present both τ_{123} and τ_{1234} disagree with observations. It is necessary to have a much greater quantity of astronomical data in the form of colour–magnitude diagrams for star clusters to determine whether the disagreement for h and χ Persei is merely accidental, or whether the coupling constant for the local $(e\nu_e)(e\nu_e)$ interaction is up to an order of magnitude smaller than the coupling constant for other vertices in the current–current theory of weak interactions. The question is further complicated by Iben's (1967) summary of objections to the identification of carbon-burning stars (and therefore to the estimate of the carbon-burning lifetime) in h and χ Persei. Iben has suggested that the identified stars may be recently formed objects in the process of gravitational contraction onto the main sequence, that it is difficult to choose one definite direction of evolution in the colour–magnitude diagram for stars in relatively young clusters, and that the position of helium-burning stars on the main sequence can vary considerably according to their chemical composition. Therefore, in order to isolate the effect of two-photon annihilation as well as possible, calculations bearing on all of these points are needed. To assist the calculations, observational colour–magnitude diagrams for a large number of clusters like h and χ Persei are also required.

V. Conclusions

In equation (12) is presented an expression for the rate of loss of energy in stars from the process of two-photon annihilation into neutrino–antineutrino pairs. The circumstances under which this reaction is competitive with other neutrino-producing processes are indicated in Section IV. The inclusion of equation (12) in astrophysical models, followed by comparisons of the models with astronomical observations, constitutes a test of the existence of charged W mesons. The present observational evidence quoted in Section IV is questionable because it seems to exclude not only the annihilation process but also the existence of other processes of loss of energy in neutrino reactions which occur either with or without mediation by charged W mesons. Although it is unlikely that the effects due to two-photon annihilation can be separated beyond doubt from the other effects mentioned in the last paragraph of Section IV, calculations concerning these effects are necessary for an understanding of the role of neutrino processes in general in stellar evolution. It is desirable to have more than the present limited amount of observational data on evolution in stellar clusters in support of these calculations.

VI. Acknowledgments

The author wishes to thank Dr. M. Nakagawa for a discussion, and Professor F. C. Barker, Professor R. H. Dalitz, and Professor I. Iben for drawing various references to his attention.

VII. References

Boccaletti, C., de Sabbata, V., and Gualdi, C. (1964).—*Nuovo Cim.* **33**, 520.
Burns, R., *et al.* (1965).—*Phys. Rev. Lett.* **15**, 42.
Campbell, J. A. (1967).—*Nucl. Phys.* B **1**, 283.
Chin, C-W., Chiu, H-Y., and Stothers, R. (1966).—*Ann. Phys.* **39**, 280.
Chiu, H-Y. (1966).—*Ann. Rev. nucl. Sci.* **16**, 591.
Chiu, H-Y., and Morrison, P. (1960).—*Phys. Rev. Lett.* **5**, 573.
Chiu, H-Y., and Stabler, R. (1961).—*Phys. Rev.* **122**, 1317.
Feynman, R. P., and Gell-Mann, M. (1958).—*Phys. Rev.* **109**, 193.
Gell-Mann, M. (1961).—*Phys. Rev. Lett.* **6**, 70.
Glashow, S. L., Schnitzer, H. J., and Weinberg, S. (1967).—*Phys. Rev. Lett.* **19**, 205.
de Graaf, T., and Tolhoek, H. A. (1966).—*Nucl. Phys.* **81**, 596.
Hayashi, C., and Cameron, R. C. (1962).—*Astrophys. J.* **136**, 166.
Hayashi, C., Hōshi, R., and Sugimoto, D. (1962).—*Prog. theor. Phys., Kyoto* Suppl. No. 22.
Hearn, A. C. (1966).—*Communs Ass. comput. Mach.* **9**, 573.
Iben, I. (1967).—*Ann. Rev. Astr. Astrophys.* **5**, 571.
Johnson, H. L., and Morgan, W. W. (1955).—*Astrophys. J.* **122**, 429.
Lamb, R. C., *et al.* (1965).—Proc. Int. Conf. on Weak Interactions, Argonne National Laboratory (Report nr. ANL-7130), Argonne, Illinois, p. 247.
Lee, T. D., and Wu, C. S. (1965).—*Ann. Rev. nucl. Sci.* **15**, 381.
Lee, T. D., and Yang, C. N. (1960*a*).—*Phys. Rev. Lett.* **4**, 307.
Lee, T. D., and Yang, C. N. (1960*b*).—*Phys. Rev.* **119**, 1410.
Loebenstein, H. M., Mingay, D. W., Winkler, H. C., and Zaidins, C. S. (1967).—*Nucl. Phys.* A **91**, 481.
Marshak, R. E. (1966).—"High-energy Physics." (Ed. G. Takeda.) p. 60. (W. A. Benjamin Inc. and Syokabō Publishing Co.: Tokyo.)
O'Dell, C. R. (1963).—*Astrophys. J.* **138**, 67.
Seaton, M. J. (1966).—*Mon. Not. R. astr. Soc.* **132**, 113.
Stephenson, G. J. (1966).—*Astrophys. J.* **146**, 950.
Vogt, E., Michaud, G., and Reeves, H. (1965).—*Phys. Lett.* **19**, 570.
Zaidi, M. H. (1965).—*Nuovo Cim.* A **40**, 502.