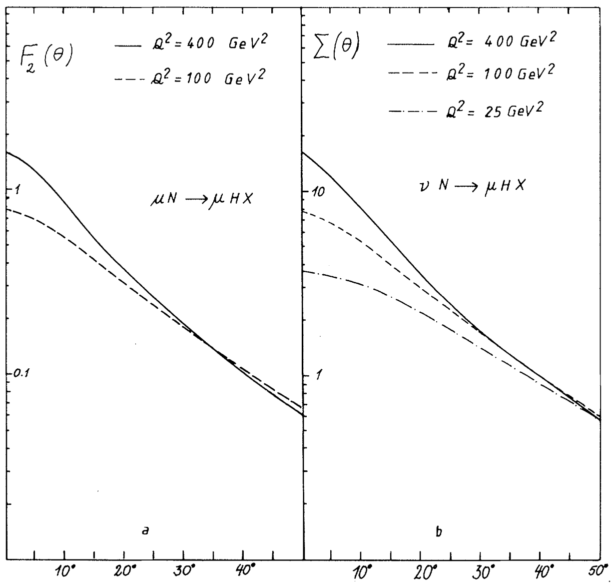

FIG. 4. Development of the acollinearity spectra with increasing Q^2 for muon and neutrino scattering of an isoscalar target.

hadron scattering.

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[1]R. Brandelik *et al.* (TASSO Collaboration), Phys. Lett. 86B, 243 (1979); D. P. Barber *et al.* (MARK-J Collaboration), Phys. Rev. Lett. 43, 830 (1979); Ch. Berger *et al.* (PLUTO Collaboration), Phys. Lett. 86B, 418 (1979); W. Bartel *et al.* (JADE Collaboration), Phys. Lett. 89B, 136 (1979).

[2]H. J. Lubatti, in Proceedings of the Eleventh International Symposium on Multiparticle Dynamics, Brugge, Belgium, June, 1980 (to be published).

[3]See e.g., F. Halzen and D. M. Scott, in Proceedings of the Eleventh International Symposium on Multiparticle Dynamics, Brugge, Belgium, June, 1980 (to be published).

[4]Y. L. Dokshitzer, D. I. D'Yakonov, and S. I. Troyan, Phys. Rep. 58, 269 (1980), and references therein.

[5]C. Berger *et al.* (PLUTO Collaboration), Phys. Lett. 90B, 312 (1980).

[6]R. Baier and K. Fey, to be published, and University of Bielefeld Report No. BI-TP-80/10 (to be published).

[7]F. Halzen and D. M. Scott, University of Wisconsin Report No. COO-881-141, 1980 (to be published).

[8]W. Marquardt and F. Steiner, DESY Report No. 80/24, 1980 (to be published).

[9]K. Kajantie and E. Pietarinen, DESY Report No. 80/19, 1980 (to be published).

[10]See, e.g., F. Janata, in *Proceedings of the 1979 International EPS Conference on High Energy Physics, Geneva, 1979* (CERN, Geneva, 1979).

[11]For the fragmentation functions, this equation was first derived by A. Bassetto, M. Ciafaloni, and G. Marchesini, Nucl. Phys. B163, 477 (1980).

[12]G. Altarelli and G. Parisi, Nucl. Phys. B126, 298 (1977).

[13]Our calculations were done with use of the parametrization of R. Baier, J. Engels, and B. Pettersson, Z. Phys. C 2, 265 (1979), for the structure functions. For the fragmentation functions we used the parametrization given in Ref. 6 independently of quark flavor. We assume $N_F = 4$ and $\Lambda = 0.5$ for α_s.

[14]J. Cleymans and M. Kuroda, to be published.

Spontaneously Broken Lepton Number and Cosmological Constraints on the Neutrino Mass Spectrum

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The cosmological constraints on neutrino masses can be altered if lepton number is broken globally giving rise to a very weakly coupled Goldstone boson—the Majoron. Then heavy neutrinos can decay sufficiently rapidly by Majoron emission, thereby giving negligible contributions to the mass density of the universe. Specifically, if M is the mass scale associated with lepton number breakdown, for $M \lesssim 10^6$ GeV there are no constraints on neutrino masses while for M high enough ($M \gtrsim 10^9$–10^{10} GeV) the standard bounds remain.

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A topic of great current interest is the spectrum of neutrino masses and its implication for the nature of the weak interactions. The most stringent limits on this spectrum come from cosmological and astrophysical arguments. If neutrinos were stable, then considerations relating to the observed mass density in the present universe indicate that the sum of neutrino masses should either be less than[1,2] about 50 eV or, alternatively, that these masses should be heavier than a few gigaelectronvolts.[3] This forbidden gap can in principle be breached by neutrinos which can decay sufficiently rapidly. Dicus, Kolb, and Teplitz[4] obtained limits on the lifetime of the process $\nu_H \to \nu_L + \gamma$ which would allow for heavy-neutrino masses to be in the forbidden zone. However, the required lifetimes found in Ref. 4 are, in general, too short to be obtained in realistic weak-interaction models. Furthermore, if one studies the effects which the photons produced in these decays have on the element abundance of the present universe,[5] one finds that the lifetimes required for heavy-neutrino decay are even shorter than those obtained in Ref. 4, being typically of the order of hours. Hence, if these cosmological arguments are correct, neutrinos with conventional weak interactions cannot have masses in the 50 eV to the gigaelectronvolt range.

In this Letter we would like to discuss how the above limits can be obviated if lepton number is a spontaneously broken *global* symmetry. We indicate here briefly the physics scenario and proceed, further on, to the details. If lepton number is indeed spontaneously broken, there is necessarily a zero-mass Goldstone boson in the theory. It turns out that, in a large class of realistic models studied by us recently,[6] this Goldstone boson (the Majoron) couples dominantly, but very weakly, to neutrinos and essentially negligibly to matter. Hence the Majoron's existence is not ruled out by experiment. Heavy neutrinos thus can decay via Majoron emission to light neutrinos, with a rate which depends both on the mass of the neutrinos as well as on the scale which characterizes the spontaneous breakdown of lepton number. The typical lifetime of neutrinos with masses in the cosmological forbidden range, and for values of the lepton-number–nonconserving scale which are sensible, turn out to be short compared with the lifetime of the universe. The heavy nuetrinos, because they have mostly decayed, will only make a negligible contribution to the mass density of the universe. Therefore this removes the need to restrict their masses. However, since the lifetime of the heavy neutrinos scales with M—the lepton-number–nonconserving scale—eventually for large enough M one has stable (with respect to the lifetime of the universe) neutrinos and one recovers the old cosmological bounds.

We shall assume that the weak interactions of the neutrinos are governed by the standard SU(2) $\otimes$U(1) theory.[7] In addition, however, we shall suppose that SU(2)$\otimes$U(1)-singlet, right-handed, neutrino fields exist. Neutrinos of a given generation can then have both lepton-number–conserving Dirac masses m, as well as lepton-number–nonconserving Majorana masses M. We shall suppose that only the right-handed neutrinos have a Majorana mass term. Then if $M \gg m$ the physical neutrinos, which are Majorana fields, have masses M and $m_\nu = m(m/M)$. The observed neutrinos are the light neutrinos of mass m_ν, while the superheavy neutrinos are at scales beyond the range of present experimentation. If M arises from the vacuum expectation value of a singlet Higgs field Φ, which carries lepton number, then $\langle\Phi\rangle \neq 0$ implies the spontaneous breakdown of lepton number. Unless lepton number is gauged, there will appear in general a Goldstone boson associated with this spontaneous breakdown—the Majoron.

In a recent note[6] we discussed the properties of the Majoron—for a one-generation model—and we summarize the salient features of that analysis here. Let ν be the light neutrino of mass m_ν and η be the superheavy neutrino of mass M. Then the effective coupling of the Majoron field χ to neutrino and matter fields (quarks and leptons) is given by

$$\mathcal{L} = -\frac{ih}{\sqrt{2}}\,\chi\left\{\bar{\eta}\gamma_5\eta - \left(\frac{m_\nu}{M}\right)^{1/2}(\bar{\eta}\gamma_5\nu + \bar{\nu}\gamma_5\eta) + \frac{m_\nu}{M}\,\bar{\nu}\gamma_5\nu + \frac{G_{\mathrm{F}}}{8\sqrt{2}\,\pi^2}\,m_\nu m_f g_f(\bar{f}\gamma_5 f)\right\}. \tag{1}$$

Here $g_f = +1$ for $f = e$ or u quarks, or $g_f = -1$ for d quarks, and h is a Higgs-fermion coupling constant. As can be seen from (1) the strength of the coupling of the Majoron to matter is extremely weak. Its coupling to neutrinos is also small, being suppressed by the factor m_ν/M. The Majoron only couples strongly to the superheavy neutrinos η. But these particles are unstable, decaying very rapidly into the light neutrinos via Majoron emission ($\tau_\eta \simeq 10^{-10}$ sec for $m_\nu \simeq 1$ eV, $h \simeq 10^{-2}$).

Because the Majoron's coupling to matter is so weak, terrestrial experiments cannot be used to rule out its existence. This point is discussed in some detail in Ref. 6, but it might be useful to summarize here the principal results of that analysis. The pseudoscalar nature of the coupling of the Majoron to matter yields a spin-dependent $1/r^3$ potential between two fermions (quarks or leptons). The strength of this potential is characterized by a coupling constant $\lambda_f = (4\pi)^{-1}(hG_F m_\nu/16\pi^2)^2$ which, taking as typical parameters $h \simeq 10^{-2}$, $m_\nu \simeq 1$ eV, is of order 10^{-65} cm^2. This number is so small that the Majoron exchange force is only comparable to gravity at typical nuclear distances. Hence Eötvos-type experiments[8] are totally insensitive to Majorons. The analysis of Feinberg and Sucher,[9] of possible nonmagnetic spin-dependent forces, again is of little use for Majorons since the bounds found for λ_f by these authors are many orders of magnitude larger than 10^{-65} cm^2 (typically $\lambda_f \leqslant 10^{-32}$ cm^2). Similarly axion searches could not have uncovered the Majoron because its coupling to matter is so weak.[10]

The analysis of Ref. 6 is easily generalized to the case of many generations of neutrinos. In this case again an equation like (1) ensues, but now, in general, the Majoron has off-diagonal couplings with neutrinos of different generations. Again, the superheavy neutrinos decay rapidly by Majoron emission. But now it is also possible for the heavier of the light neutrinos ν_H to decay to the lightest state ν_L by Majoron emission: $\nu_H \to \nu_L + \chi$. The lifetime for such a process is easily estimated

$$\tau(\nu_H \to \nu_L + \chi) = \frac{32\pi}{h^2 \sin^2\theta}\left(\frac{M}{m_{\nu_H}}\right)^2 \frac{1}{m_{\nu_H}}, \qquad (2)$$

where $\sin\theta$ is an intrageneration mixing angle.

The lifetime τ can be sufficiently short on a cosmological scale to remove effectively the ν_H contribution to the mass density of the universe, provided that M is not too large. If we take as typical parameters $h \simeq 10^{-2}$, $\sin\theta \simeq 10^{-1}$, and $M \simeq 10^5$ GeV, we find that $\tau \simeq 2\times 10^7$ yr for $m_{\nu H} = 100$ eV. Thus we see that for "sensible" scales M, where one might expect new physics to arise, neutrino decay by Majoron emission can play an effective role in removing the heavy-neutrino contribution to the mass of the universe. Actually the above discussion can be sharpened. One can determine a range of values of M for which heavy neutrinos of mass $m_{\nu H}$ can exist without violating any cosmological bounds. We proceed to do this below.

Our discussion here will be analogous to the analysis of Dicus, Kolb, and Teplitz.[4] First we note that the dominant interaction of ν_H and ν_L is through the standard weak-interaction neutral-current process $\nu_H + \bar{\nu}_H \to \nu_L + \bar{\nu}_L$. Thus the calculation of when the heavy neutrinos decouple remains the same as that of Ref. 4 and we will adopt their values for the decoupling temperature T_D, decoupling time t_D, and the density $n_H(T_D)$ at that epoch. One can then calculate the contribution to the energy density of the Majorons coming from ν_H decay. [There is, of course, a remnant Majoron black-body energy density coming from Majorons which decoupled at temperatures of the order of M, but this component—like that of the photons—makes a totally negligible contribution to the present-day energy density.] The energy density of the Majorons which are decay by-products is given by[4]

$$\rho = 2n_H(T_D)\left(\frac{1.9\,^\circ\mathrm{K}}{T_D}\right)^3 \int_{t_D}^{t_U} dt \left(\frac{m_{\nu H}}{2}\right)\left(\frac{t}{t_U}\right)^{1/2} \frac{1}{\tau}\exp\left(-\frac{t-t_D}{\tau}\right). \qquad (3)$$

Here t_U is the universe lifetime and t_D is the decoupling time for ν_H. The remaining factors in Eq. (3) are easily understood. The overall factor of 2 accounts for both heavy neutrino and antineutrino decays. The factor $(1.9\,^\circ\mathrm{K}/T_D)^3$ takes into account of the volume expansion from the decoupling of the neutrinos to the present epoch, while $n_H(T_D)\exp[-(t-t_D)/\tau]$ is the ν_H density at the moment of decay. Finally the factor $\frac{1}{2}m_{\nu H}(t/t_U)^{1/2}$ accounts for the red shift of the Majoron's energy from its decay value ($\frac{1}{2}m_{\nu H}$), while τ^{-1} in Eq. (3) is just the probability that ν_H decay occurred.

Since $t_d \simeq 10^{-1}$ sec we have, in general, $\tau \gg t_d$. Furthermore, if $\tau \ll t_U$ we can approximate (3) by

$$\rho \simeq n_H(T_D)\left(\frac{1.9\,^\circ\mathrm{K}}{T_D}\right)^3\left(\frac{\sqrt{\pi}}{2}\right) m_{\nu H}\left(\frac{\tau}{t_U}\right)^{1/2}. \qquad (4)$$

So as to avoid conflicts with astrophysical bounds we must require that the energy density of the Majorons from ν_H decay be less than the critical density[5] $\rho_c \simeq 5\times 10^{-3}$ MeV/cm^3. This requirement for stable neutrinos [essentially $\tau = t_U$ in Eq. (4)] gives the cosmological forbidden zone for neutrino masses[1-3]: 50 eV $\lesssim m_\nu \lesssim$ 2 GeV. In our case, since we have a free parameter M—the scale of the lepton-number breakdown—the constraint $\rho \leqslant \rho_c$ will give for a given M allowed values for m_{ν_H}. Using again as typical parameters $h \simeq 10^{-2}$ and $\sin\theta \simeq 10^{-1}$, one obtains the graph of Fig. 1 for the allowed zone of neutrino masses. We note the important result than if $M \lesssim 10^6$ GeV there are *no* cosmological constraints on neutrino masses.

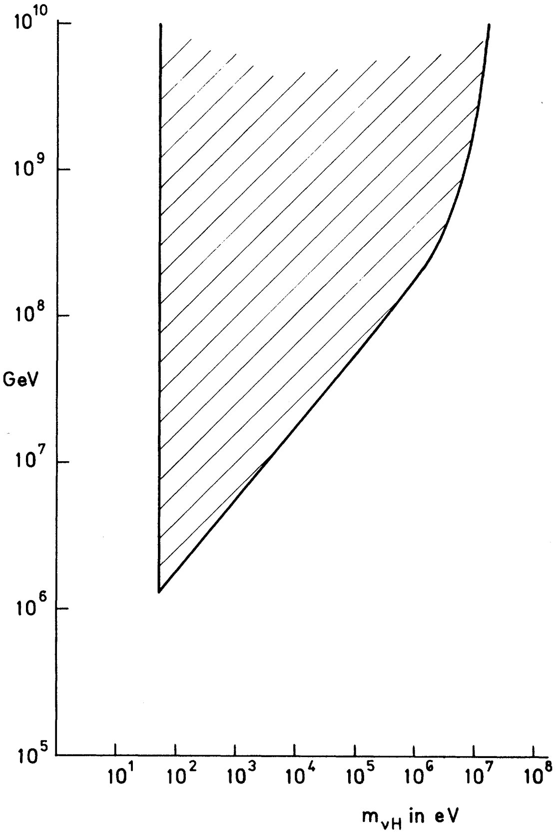

FIG. 1. The shaded area represents the forbidden domain of neutrino masses for a given range of the heavy Majorana lepton mass M.

On the other hand, if $M \gtrsim 10^9$–10^{10} GeV, the forbidden zone for neutrino masses essentially remains that of the standard analysis.[1-3] We remark that the straight-line portion of the boundary in Fig. 1 for $m_{\nu_H} \lesssim 1$ MeV arises because there $n_H(T_D)$ and T_D are fixed [see Ref. 4: $n_H(T_D) = 6 \times 10^{32}$ cm^{-3} and $T_D = 3.4\times10^{10}$ °K]. Then $\rho \simeq (M^2/m_{\nu_H})^{1/2}$. Above $m_{\nu_H} \simeq 1$ MeV the density factor $n_H(T_D)(1.9\ °\mathrm{K}/T_D)^3$ decreases rapidly and the bound $\rho \lesssim \rho_c$ begins to be ineffective.

We have shown that, if lepton number is a spontaneously broken global symmetry accompanied by a very weakly coupled Goldstone boson, it is possible to avoid cosmological constraints on the neutrino spectrum provided the scale $M \lesssim 10^6$ GeV. One may ask whether this is a reasonable scale for lepton-number breakdown. In terms of the vacuum expectation value of the Higgs field Φ which gives rise to the breakdown one has that $\langle\Phi\rangle \simeq M/h \lesssim 10^8$ GeV. Breakdown of lepton number at scales below these have been invoked previously in the context of left-right-symmetric models[11] and in connection with horizontal symmetries.[12] Thus in this sense the bound on $\langle\Phi\rangle$ seems an eminently reasonable one. However, we should point out that if the lepton-number breakdown occurs at a grand unified scale, one expects $\langle\Phi\rangle \simeq 10^{13}$–$10^{15}$ GeV and in that case even if Majorons exist, no lifting of the cosmological constraints is possible.

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[1]R. Cowsik and J. McClelland, Phys. Rev. Lett. 29, 669 (1972).

[2]D. N. Schram and G. Steigman, Gen. Relat. Grav. (to be published).

[3]B. W. Lee and S. Weinberg, Phys. Rev. Lett. 39, 165 (1977).

[4]D. Dicus, E. Kolb, and V. Teplitz, Phys. Rev. Lett. 39, 168, 973(E) (1977).

[5]K. Sato and M. Kobayashi, Prog. Theor. Phys. 58, 1775 (1977); D. Dicus, E. Kolb, V. Teplitz, and R. Wagoner, Phys. Rev. D 17, 1529 (1978).

[6]Y. Chikashige, R. N. Mohapatra, and R. D. Peccei, Max Planck Institute Report No. MPI-PAE/PTh 36/80, 1980 (to be published).

[7]S. Weinberg, Phys. Rev. Lett. 19, 1264 (1967); A. Salam and J. C. Ward, Phys. Lett. 13, 168 (1964); S. L. Glashow, J. Iliopoulos, and L. Maiani, Phys. Rev. D 2, 1285 (1970).

[8]P. G. Roll *et al.*, Ann. Phys. (N.Y.) 26, 442 (1964); R. Spero *et al.*, Phys. Rev. Lett. 44, 1645 (1980).

[9]G. Feinberg and J. Sucher, Phys. Rev. D 20, 1717 (1979).

[10]R. D. Peccei, in *Proceedings of the Nineteenth International Conference on High Energy Physics, Tokyo, Japan, 1978*, eddited by S. Homma, M. Kawaguchi, and H. Miyazawa (Physical Society of Japan, Tokyo, 1979); see also D. Dicus *et al.*, Phys. Rev. D 22, 839 (1980).

[11]R. N. Mohapatra and G. Senjanović, Phys. Rev. Lett. 44, 912 (1980), and Fermilab Report No. 80/61 THY (to be published); R. N. Mohapatra and R. E. Marshak, Phys. Rev. Lett. 44, 1316 (1980).

[12]Y. Chikashige, G. Gelmini, R. D. Peccei, and M. Roncadelli, Phys. Lett. 94B, 499 (1980).