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Theoretical possibilities and observational constraints for radiatively decaying neutrinos with mass near 30 eV

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Abstract

Massive relic neutrinos may contribute significantly to the density of the Universe, and if appropriately concentrated, could explain characteristics of luminous matter in galaxies. As pointed out by Sciama, the radiative decay of 30 eV neutrinos could explain several observational puzzles including the anomalous degree of ionization of interstellar matter within the Milky Way Galaxy. Although these neutrinos are impossible in the standard particle physics model, we show that various non-standard particle models with extended scalar sector or minimal supersymmetry have sufficient freedom to accommodate this scenario. We review the cosmological impact of these neutrinos, and discuss observational constraints on them in the immediate Solar neighborhood, in nearby regions of low interstellar absorption, in the Galactic halo, in clusters of galaxies, and in extragalactic space. Contrary to often-expressed opinions, present observational constraints only rule out models in which extreme concentrations of radiatively decaying neutrinos occur in clusters of galaxies.

Ref. S.I.S.S.A. 115/94/EP (November 1994)

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1 Introduction

In the minimal standard model of electroweak particle interactions, neutrinos are massless and, hence, absolutely stable against decay. Clearly, such neutrinos cannot contribute to the missing dark matter in the universe. Although the experimentally known properties of all the three species of neutrinos are in good agreement with the standard model, the observational mass limits of the ν_μ , 220 keV [1], and ν_τ , 31 MeV [2], are sufficiently high that they are of no relevance to the questions discussed here. Relic cosmic neutrinos must not have a mass larger than 30 eV or they would decelerate the universe in excess of observational limits.

A relic neutrino ν_2 could decay radiatively into a lighter neutrino ν_1 and a photon if it has a nonvanishing mass m_2 and an electromagnetic coupling. The decay photons would be monochromatic, with energy

$$\mathcal{E}_\gamma = \frac{m_2^2 - m_1^2}{2m_2} \approx \frac{1}{2}m_\nu . \quad (1)$$

Here we denote $m_2 = m_\nu$ when m_1 is negligible.

If the neutrino mass is in the range stated above, the photon wavelength λ is in the far ultraviolet (FUV) or the extreme ultraviolet (EUV) bands. Recall that the dividing line between FUV and EUV is the hydrogen ionizing Lyman limit at energy $E_\gamma = 13.6$ eV corresponding to the wavelength $\lambda = 91.1$ nm. Since astronomers are more accustomed to expressing wavelengths in Å units and physicists in energy units, it may be useful to cast Eq. (1) in the form,

$$\left(\frac{\lambda}{1000\text{Å}} \right) \cdot \left(\frac{m_\nu}{25\text{eV}} \right) \approx 1 . \quad (2)$$

The hypothesis of a radiatively decaying relic neutrino is attractive in that it establishes a plausible identity for dark matter manifested in clusters of galaxies and for the high degree of ionization of intergalactic gas as evidenced by the Gunn-Peterson effect [3]. If these neutrinos have an appropriate mean life of about 10^{24} s, the photons produced can furnish the ionizing energy needed to maintain the high ionization state of intergalactic hydrogen [4, 5, 6, 7]. Moreover, Sciama has shown [8] that the decay radiation is capable of explaining the anomalously high degree of hydrogen ionization within our Milky Way Galaxy if the radiatively decaying neutrino has a mass and mean life in the narrow window

$$m_\nu = (27.5 \pm 0.3) \text{ eV} \quad t_\nu = (2 \pm 1) \cdot 10^{23} \text{ s} . \quad (3)$$

This value for m_ν corresponds to a photon wavelength of 89–91 nm. This massive neutrino could well be the tau neutrino because particle physics provides a much larger upper limit to its mass, as noted above [9]. Sciama assumes that the lighter species has a mass $m_1 < 3$ eV, thus justifying the approximation in Eq. (1).

There are very few alternative particles to explain the hydrogen ionization in our Galaxy or in the universe at large. It would be very unnatural for ordinary bosons to have a mass as low as 28 eV because there is no natural way to protect the mass from

being extremely heavy. Although there are models with light massive bosons, like axions or pseudo-Goldstone bosons which get their masses through some gravity effect, these have not gained wide support. The required model would be far more complicated and more *ad hoc* than is the case for neutrinos.

Still another alternative would be to introduce a new, sterile neutrino. Such models have been introduced [10] to allow conventional particle physics solutions to the solar electron neutrino deficit and the atmospheric muon neutrino deficit, but their validity for unification remains to be proven.

Since the relic number density of neutrinos is very large, the mass represented by Eq. (3) would give an appreciable contribution to the total mass density of the universe, perhaps even closing the universe ($\Omega_0 > 1$). Neutrinos are hot dark matter (HDM), and an excessive contribution from them would have unwanted effects on galaxy formation. In Section 2 we discuss the implications of neutrino dark matter in general, and in particular its role in galaxies. We briefly discuss mechanisms of galaxy formation consistent with this role.

While physical theories admitting massive neutrinos can easily be formulated, the electromagnetic coupling of these neutral particles requires more elaboration. The standard model of electroweak interactions admits neutrino radiative decay only in second-order processes. This makes the radiative decay extremely slow, yielding a mean lifetime six orders of magnitude larger than the mean lifetime quoted in Eq. (3) for a 28 eV mass.

In Section 3 we describe the reasons for this lifetime in the standard model and discuss various non-standard models in which the neutrinos may have other interactions that could speed up the radiative decay considerably. We do not address the problem of the solar electron neutrino deficit or the atmospheric muon neutrino deficit, which require some other explanation.

In Section 4 we turn to the observational aspects of cosmic populations of neutrinos radiating in the EUV band in: the immediate Solar neighborhood, directions of low galactic absorption, the Galactic halo, nearby clusters of galaxies, and throughout extragalactic space. As we shall see, FUV and EUV observations rule out only those scenarios in which massive neutrinos concentrate to an extreme degree. Hence, the massive neutrino remains an extremely interesting possibility that deserves to be tested carefully.

In Section 5 we summarize our results and suggest an alternative scenario that satisfies observational constraints. We also discuss forthcoming experiments that will be able to detect the neutrino hypothesized by Sciama, and will constrain other models of massive neutrinos which decay radiatively.

2 Cosmological aspects

2.1 Neutrino dark matter

In the hot big bang model the present universe has emerged from a thermal bath in which nucleons, charged leptons, neutrinos, and photons were in thermal equilibrium. From the equilibrium distributions, the number density of neutrinos N_ν of each species is related

to the number density of photons N_γ at the same temperature T by

$$N_\nu = \frac{3}{4} N_\gamma . \quad (4)$$

When the universe cooled below the threshold for electron-positron pair production the temperature of the photons increased by a factor $(11/4)^{1/3} \simeq 1.40$. Noting that number densities are proportional to T^3 , and $N_\gamma = 410$ photons/cm³ at the present time the number of neutrinos is

$$N_\nu = \frac{3}{11} N_\gamma \simeq 112 \frac{\text{neutrinos}}{\text{cm}^3} , \quad (5)$$

and their temperature is $T_\nu \simeq T_\gamma/1.40 \simeq 1.946$ K. Thus the present-day momenta of primordial neutrinos should be about that of the cosmic microwave background photons, unless the neutrinos have been affected by gravitational processes.

There is solid observational evidence for nonluminous mass in all large scales in the universe [11]. Rotation curves of galaxies and virial theorem mass estimates of clusters of galaxies reveal mass-to-light ratios much greater than ratios expected from stellar populations. We parameterize the present-day average mass density of the universe ρ_0 by the ratio $\Omega_0 = \rho_0/\rho_{\text{crit}}$, where $\rho_{\text{crit}} = 3H_0^2/8\pi G$ is the critical density above which the universe is bound. Sandage [12] reviews the many observational approaches to constraining Ω_0 and concludes that the best current data require $0.5 < \Omega_0 < 2.8$. Only a tiny fraction of this mass (about 1%) is luminous matter within galaxies. Of the remainder, perhaps 10% is found in groups and clusters of galaxies.

It may be that the universe is critical with a present-day mass density ρ_0 equal to the critical mass density $\rho_c = 10.6h^2$ keV/cm³; this would be in agreement with the expectations of inflationary models. One is then led to the conclusion that 90% of the matter in the universe is dark and its composition unknown. It is also arguable that the universe may be purely baryonic. This scenario is of no interest in terms of this paper.

If dark matter exists and the particle is a neutrino, it would contribute a missing Ω_ν if its mass is

$$m_\nu = \Omega_\nu \frac{\rho_c}{N_\nu} \simeq 95\Omega_\nu h^2 \text{ eV} . \quad (6)$$

A neutrino mass of 28 eV combined with $0.5 < \Omega_\nu < 1.0$ gives a normalized Hubble constant of $h = 0.54$ to 0.77 .

The number of distinct neutrino species is known to be three from the measurements at LEP of the decay rate of Z bosons into invisible channels [9]. Thus the well known e , μ , and τ families are complete. Note that the result (6) applies to the total mass of neutrinos, so if only one of these species is this heavy, the two other species must contribute only negligibly to the mass density of the universe.

2.2 Neutrinos in galaxies

One of the attractive features of massive neutrinos is they could provide a straightforward solution to the problem of galactic rotation curves. The missing mass needed to explain

the Milky Way rotation curve has a localized density of $0.6 - 1.7 \text{ GeV cm}^{-3}$ depending upon the assumed halo geometry (spherical or ellipsoidal). If this component is 30 eV neutrinos, their maximum number density is $(2-6) \times 10^7 \text{ cm}^{-3}$.

In galaxy-binding models of this sort, the aggregate mass of the neutrino halo dominates the baryonic mass, but the decay luminosity of the neutrinos is many orders of magnitude below the luminosity of the host galaxy since the neutrino lifetime is far greater than the age of the universe. It is for this reason that the neutrinos behave as dark matter: they provide binding with very little emission. It is not presently known if such large concentration factors (of the order of one million) actually occur. Steigman et al. [13] furnish a plausible mechanism for binding cosmic neutrinos into large galaxies and clusters as part of their formation.

Large concentration factors are not ruled out by phase space occupation arguments (e.g., Tremaine & Gunn [14]) as long as the momentum distribution broadens as neutrinos accumulate. To achieve a millionfold increase in density, the neutrino momenta need to increase a hundredfold. Suppose bound fermions of mass m are momentum broadened to occupy a range of velocities $v \approx v_{\text{escape}}$. Then to exhibit a phase space density less than $2/h^3$ yet yield a mass density ρ , they must satisfy

$$m > (0.5\rho(h/v)^3)^{1/4} . \quad (7)$$

This fermion mass lower limit is $\sim 30 \text{ eV}$ for conditions typical of our Milky Way Galaxy. $\rho \approx 1 \text{ GeV cm}^{-3}$ and $v \approx 3 \times 10^7 \text{ cm s}^{-1}$ consistent with the massive neutrinos considered herein.

2.3 Galaxy formation models

The sizes and distributions of present-day galaxies and galaxy clusters reflect the spectrum of initial density fluctuations, seeded by random motion or perhaps by cosmic strings or topological defects. The most common assumption is that the power spectrum of the fluctuations is simply proportional to the wave number k of the fluctuations at the time when they crossed into our horizon, in which case, density variations of all wavelengths are generated with equal amplitudes. Inflationary models also predict this scale-invariant power spectrum (the Harrison-Zel'dovich spectrum).

In general, the power spectrum may be of the form

$$P(k) \propto k^n , \quad (8)$$

where n is called the spectral index. The Harrison-Zel'dovich spectrum corresponds to $n = 1$. If n were smaller than 1, large-scale density fluctuations would dominate and vice versa.

If the density fluctuations were carried by hot neutrinos and the spectrum were of scale-invariant form, free streaming neutrinos would have such large velocities and such weak interactions that they would only stop in overdensities of 40 Mpc (present-day) linear size and $10^{15} M_{\odot}$ mass. Thus, scale-invariant HDM models predict a 'top-down' scenario where pancake-shaped supergalaxies form first and galaxies form later from the crumbs of collapsing pancakes. But supergalaxies at typical distances $z \lesssim 0.5$ are not as

old as quasars with redshifts up to $z = 4-5$. There are also several examples of galaxies that are older than the groups in which they are found now. Moreover, galaxies in our neighborhood are generally falling in toward the Virgo cluster rather than streaming away from it, as one would expect in the ‘top-down’ scenario.

Computer simulations of pancake formation and collapse show that matter at the end of the collapse is so shocked that it does not condense but remains ionized, unable to form galaxies and attract neutrino halos. Moreover, large clusters have higher escape velocities so they should trap five times more neutrinos than large galaxies. This is not supported by observations, which show that the ratio of dynamic mass to luminous mass is actually greater in smaller objects; furthermore, in larger objects the dark matter is more concentrated near the center [15]. Thus galaxy formation arguments suggest that only a fraction of dark matter could be neutrinos, or else that the primordial spectrum was not of the scale-invariant form.

We now know that the very popular unbiased cold dark matter scenarios (CDM) (which depend on particles other than neutrinos) have difficulties in matching the relatively quiet velocity field of galaxies (they should fall in faster towards the deep CDM potential wells), and they lack power on large scales (they underproduce large structures).

Since no self-consistent set of simple initial conditions like those of CDM or HDM scenarios matches the observed universe, one may need some mixture of CDM, HDM and baryonic dark matter [16, 17]. In Section 5 we shall come back to a discussion of what kind of particles would be consistent with present observations in the ultraviolet.

3 Radiative neutrino decay in nonstandard models

3.1 General approach

In the minimal standard model the neutrinos are massless and left-handed (negative helicity states, spin and momentum vectors antiparallel), right-handed neutrinos being absent. Enlarging the model to also include right-handed states, all neutrinos would be naturally massive four-component Dirac particles like the electrons, but there would be no explanation for the smallness of their mass. The simplest explanation would be the see-saw scenario: the right-handed components have a very large Majorana mass (a term connecting neutrinos and antineutrinos), as a result of which there will be two mass eigenstates, one very heavy and one very light. These neutrinos would be two component Majorana particles that are their own antiparticles and do not have any conserved charge. Alternatively, one can generate masses through radiative mechanisms that will be discussed below. Although physicists favor Majorana neutrinos over Dirac neutrinos, we shall return to a discussion of the latter in Section 3.5.

In the standard model a small electromagnetic coupling arises because of radiative corrections involving the weak gauge boson W when right-handed neutrinos are included. In the perturbative picture of the quantum theory, the electromagnetic coupling is generated as a second-order perturbation at one-loop level, as is described by the Feynman graphs shown in Fig. 1. This loop now involves two charged states, the W^- and a charged lepton ℓ_i^- , to which the photon can couple. Radiative decay of massive neutrinos is caused by a

Figure 1: The diagrams generating radiatively the magnetic moment in the standard model with massive Dirac neutrinos.

transition magnetic moment μ_{ij} , antisymmetric for Majorana neutrinos [18], with a rate

$$\Gamma(\nu_i \rightarrow \nu_j \gamma) = \frac{|\mu_{ij}|^2}{8\pi} m_i^3 \left(1 - \left(\frac{m_j}{m_i} \right)^3 \right). \quad (9)$$

To obtain the decay rate preferred by Sciama, $\Gamma = (1 - 0.3) \cdot 10^{-23} \text{ s}^{-1}$ one needs a transition moment

$$\mu_{ij} \sim (0.5 - 1) \cdot 10^{-14} \mu_B. \quad (10)$$

where μ_B is the Bohr magneton. Note that the electromagnetic moment coupling always flips the helicity (spin) of the neutrino. For Dirac neutrinos this means that if a left-handed neutrino flips to a right-handed state, the resulting neutrino would be practically inert since the weak interactions treat different helicity states differently because of parity violation.

The magnetic moment of a massive neutrino in the standard model is, however, very small, $3 \cdot 10^{-19} \mu_B (m_\nu/\text{eV})$, since it is proportional to the neutrino mass. The dependence on the neutrino mass can be understood heuristically: the graph must explicitly cause a transition from the left-handed fermion state to a right-handed state. Such a transition can result from a vertex involving a scalar particle or a mass term. Here we have no scalars present, and the purely left-handed coupling to the W boson forbids any left to right transition in the internal fermion line. Hence, the only possible cause of the flip would be the external neutrino mass.

In order to obtain enhanced electromagnetic couplings, we have to utilize models that allow the flip to take place inside the loop. A necessary condition for this outcome is an explicit interaction, coupling to both the left- and right-handed neutrino states. In the simplest scenario, the gauge bosons (with spin 1) are replaced by one new scalar (spin 0) boson with electric charge -1 and a singlet under $\text{SU}(2)$. The flip can then take place in the internal charged lepton line, and the result is proportional to the charged lepton mass. With this kind of model one can obtain magnetic moments up to $10^{-11} \mu_B$ [19, 20].

Figure 2: The diagrams generating neutrino mass radiatively in the Zee model.

Although we have removed the explicit dependence of the neutrino mass, the result is still related to it because the same loops without the external photon will cause a correction to the mass term. Actually, this provides an alternative way to generate neutrino masses. We may assume that neutrinos are totally massless at the first order of perturbation theory because of the absence of the right-handed states. The mass will then arise at higher orders, and, in a natural manner, it can be much smaller than the masses of other particles.

The Zee model [21, 22] radiatively produces both the mass and the magnetic moment with a scalar exchange loop. In its simplest form this model cannot consistently satisfy S-ciani's conditions. It can be improved by combining it with the see-saw model. Scenarios similar to the Zee model appearing in supersymmetric models may also be plausible. All these models generate sizable mixings between neutrino flavors, and are currently close to being tested with neutrino oscillation experiments.

In the previous discussion, we considered the helicity flip to be caused by the mass in the internal line. There were, in fact, three flips, the other two flips occurred in the scalar vertices. One flip would be sufficient, and this scenario (the Barr-Freire-Zee mechanism) can be realized if the lepton line couples to a scalar at one vertex and to a gauge boson at the other [22]. In this way one can obtain magnetic moments independent of the charged lepton masses but at the price of going to higher orders in the perturbation theory. Despite the loop suppression one can obtain sufficient parameter space to escape the stringent laboratory constraints for neutrino oscillations.

3.2 Scalar exchange models

The Zee model expands the scalar sector of the standard model by adding two new scalar bosons—a singly charged $SU(2)$ singlet (η^-) and a second doublet (ϕ_2)—in addition to the Higgs doublet (ϕ_1) present in the standard model. Both doublets will develop non-zero vacuum expectation values, causing the scalar bosons to mix. To avoid flavor changing neutral currents, one requires that only one of the Higgs doublets couples to a given fermion, e.g., ϕ_1 to up-type quarks, and ϕ_2 to down-type quarks and charged leptons. This structure is achieved by imposing an appropriate symmetry on the theory, e.g., a discrete Z_2 symmetry.

In the Zee model the mass of the three neutrinos is generated by the graph in Fig. 2. This mass is readily calculated, and one obtains the elements of the 3×3 neutrino mass

matrix in the flavor basis (i.e., the physical eigenbasis of the charged leptons) [23, 24, 20].

$$m_{ik} = \frac{f_{ik}g \sin 2\beta \cot \alpha}{32\pi^2 M_W} (M_2^2 - M_1^2) [m_i^2 I(m_i, M_1, M_2) - m_k^2 I(m_k, M_1, M_2)] . \quad (11)$$

where f_{ik} is the dimensionless coupling constant of η to the lepton doublets i and j , and α and β are scalar mixing angles. m_i are the masses of the charged leptons of the i th family. M_1 and M_2 are the masses of the physical Higgs scalars, and

$$\begin{aligned} I(m, M_1, M_2) &= \frac{1}{M_2^2 - M_1^2} \left(\frac{M_1^2}{M_1^2 - m^2} \ln \frac{m^2}{M_1^2} - \frac{M_2^2}{M_2^2 - m^2} \ln \frac{m^2}{M_2^2} \right) \\ &\approx \frac{1}{M_2^2 - M_1^2} \ln \frac{M_2^2}{M_1^2} . \end{aligned} \quad (12)$$

Because of Fermi-Dirac statistics, the coupling constants f_{ij} are antisymmetric in the flavor indices i and j , and the diagonal elements of the mass matrix vanish. The neutrino mass matrix can be diagonalized by a unitary rotation, and the eigenstates of the mass matrix that represent the physical particles are linear combinations of the neutrino states in the flavor basis.

With standard leptons one can actually obtain neutrino masses as heavy as tens of keV within the phenomenological constraints of the couplings. Note that the mass formula depends on the square of the internal lepton mass, one power of mass because of the internal propagator term, and the other follows from the proportionality of the Higgs coupling to the mass of the fermion. The most natural neutrino mass spectrum consists of one light state and two heavier states, almost degenerate with very large mixing. This, however, conflicts with the cosmology bound Eq. (6) if the heavy states have the Sciama mass [Eq. (3)]. The general phenomenology of this model has been recently studied (in ref. [25]).

A magnetic moment arises at one loop level by coupling a photon to the internal lines as shown in Fig. 3. Now one has [20]

$$\begin{aligned} \mu_{ij} &= \frac{gm_\epsilon f_{ij}}{32\pi^2 M_W} \sin 2\beta \cot \alpha \mu_B \\ &[m_i^2 \{I(M_1, m_i, m_i) - I(M_2, m_i, m_i)\} + m_j^2 \{I(M_1, m_j, m_j) - I(M_2, m_j, m_j)\}] . \end{aligned} \quad (13)$$

where the function $I(M, m, m)$ can be expressed more simply

$$I(M, m, m) = \frac{M^2}{(M^2 - m^2)^2} \ln \frac{M^2}{m^2} - \frac{1}{M^2 - m^2} \approx \frac{1}{M^2} \ln \frac{M^2}{m^2} - \frac{1}{M^2} . \quad (14)$$

Here the magnetic moment depends on the squares of the lepton masses.

The magnetic moment and the mass are now closely related, so one can write the magnetic moment in terms of the corresponding mass element

$$\mu_{ij} = m_{ij} m_\epsilon \frac{I(M_1, m_i, m_i) - I(M_2, m_i, m_i)}{(M_2^2 - M_1^2) I(m_i, M_1, M_2)} \mu_B . \quad (15)$$

Figure 3: The graphs generating the electromagnetic coupling to the neutrinos in Zee model.

For the optimal case $M_1 \approx M_2 \equiv M$ this can be written more explicitly

$$\mu_{\tau\ell} \simeq \left(\frac{m_{\tau\ell}}{1\text{eV}}\right) \left(\frac{45\text{GeV}}{M}\right)^2 10^{-15} \mu_B, \quad (16)$$

where $\ell = \mu, \epsilon$. We see the generic phenomenon that the ratio of magnetic moment to mass is inversely proportional to the square of the mass of the scalar bosons. The latter is bounded by accelerator experiments that have sought but not found signatures of new particles up to about 45 GeV [9]. Note, however, that one cannot have $M_2 = M_1$ exactly; in that case both the mass and the magnetic moment would be zero.

Although the Zee model provides a sufficiently large transition moment, we have seen that it fails by generating a mass spectrum in conflict with cosmology. The absence of diagonal elements in the Zee model mass matrix resulted from the symmetry conditions in lepton flavor indices for the couplings to scalars: one coupling was antisymmetric because of Fermi-Dirac statistics, and the other was symmetric by definition. The first condition can be removed by adding more lepton doublets, while the second is relaxed if the scalar doublet does not contribute to the tree-level mass of the charged fermions, i.e., it has a zero vacuum expectation value.

The next simplest radiative model, the symmetric two-loop of Babu [26] is capable of producing an acceptable mass matrix but does not give radiative decays at observational levels. In this model the transition moment involves an additional suppression factor proportional to the squares of the charged lepton masses in the loop because of the symmetric structure of the loop.

We may get a better model if we would use two different mechanisms, one to generate a diagonal mass element for the heaviest neutrino and the other to generate non-diagonal elements and transition moments. Combining the Zee and Babu models might do, though it may seem odd (but not necessarily unnatural) to have the two loop result exceeding the one loop result. Indeed, there is a variation of the Babu model having the particle

content of the Zee model and containing sufficient undetermined coupling constants to allow the tuning of an appropriate mass matrix. A more plausible alternative may be to combine the Zee model with see-saw generated masses.

Assume now that we have one heavy (28eV) neutrino with its mass generated either by the see-saw mechanism or Babu's two-loop model (or a modified Zee model with an exotic fermion inside the loop). Suppose that the Zee loop generates both a transition moment between this neutrino and lighter states, and a mass mixing between them. If the heavy state is the tau neutrino, the non-diagonal mass elements are bound by laboratory experiments in the following way [27]

$$m_{e\tau} < 3.8 \text{ eV} \quad (\Leftarrow \sin^2 2\theta_{e\tau} < 0.07) \text{ and} \quad (17)$$

$$m_{\mu\tau} < 0.9 \text{ eV} \quad (\Leftarrow \sin^2 2\theta_{\mu\tau} < 0.004), \quad (18)$$

where $\theta_{i\tau}$ is the mixing angle. The corresponding limit for the possibility $\nu_\mu \rightarrow \nu_e + \gamma$ would be

$$m_{e\mu} < 0.7 \text{ eV} \quad (\Leftarrow \sin^2 2\theta_{e\mu} < 0.0025). \quad (19)$$

Hence the only sizable transition moment can arise between tau and electron neutrinos. If we assume that there are no accidental cancellations between different components, the Eq. (16) gives an upper limit for the transition moment.

$$\mu_{e\tau} < 4 \cdot 10^{-15} \mu_B. \quad (20)$$

This may be only marginally consistent with Sciama's requirement, but it could be improved by allowing some mild cancellations. Consequently, one can predict that the electron neutrino mass is about 0.5 eV, accessible for searches of neutrinoless double beta-decay. Future neutrino oscillation experiments will clarify whether these kinds of scenarios are viable.

3.3 Two-loop graphs

So far we have only considered one-loop transition moments. Inevitably, however, there will also be two-loop contributions, and in certain cases the two-loop contributions may dominate over one-loop. Consider the two-loop diagram in Fig. 4 [28] that is always present in the general Zee model. The crucial trick is to now use a W gauge boson in the loop. This gives two benefits: the coupling of the gauge boson to a fermion is much larger than the Yukawa couplings of the scalars, and the amplitude does not explicitly contain the lepton mass, since the flip occurs only once, because of the presence of only one scalar coupling. One can estimate the magnetic moment to be

$$\mu_{ij} \simeq \sum_{abc} \frac{f_{ij} U_{ab} K_{bc} U_{ca}^\dagger g^2 m_\epsilon}{(16\pi^2)^2 M_W^2} I(M_W, M_a, M_b, M_\eta) \mu_B. \quad (21)$$

where I is a very complicated function of the internal masses. K is the scalar self-coupling of the mass dimension, and U is the scalar mixing matrix. The scalar mixing is non-zero, as is evident from the unequal number of charged and neutral states.

Figure 4: The graphs generating the electromagnetic coupling to the neutrinos in the Barr-Freire-Zee model.

Because of the absence of the lepton mass factor, the two-loop contribution may be substantial and can be compared with the one-loop contribution. The ϵ - μ component of the transition moment can be enhanced by 2 orders of magnitude with respect to the one-loop term, which depends on m_μ^2 . The enhancement is smaller for the other components with the one-loop term proportional to m_τ^2 , and typically the one-loop contribution wins over the two-loop. There is, however, a certain range of parameters that may also lead to the amplification of these terms by about a factor of 3. Hence, one can obtain transition moments

$$\mu_{\tau\ell} \lesssim \left(\frac{m_{\tau\ell}}{1\text{eV}} \right) 2 \cdot 10^{-15} \mu_B \quad (22)$$

or

$$\mu_{\mu\epsilon} \lesssim \left(\frac{m_{\mu\epsilon}}{0.3\text{eV}} \right) 10^{-14} \mu_B. \quad (23)$$

Thus, we can now easily accommodate Sciama's value Eq. (10) even when the decaying particle is a muon.

Note that no two-loop contribution to the mass exists because of spin conservation [22]. Nevertheless, the two-loop transition moment does depend on the one-loop mass because of common parameters. The expected mass scale for the electron neutrinos would now be 0.1–0.5 eV if the heavy state is the tau neutrino, or $(0.1\text{--}3) \cdot 10^{-2}$ eV for the channel $\nu_\mu \rightarrow \gamma \nu_\epsilon$.

3.4 Supersymmetric models

The minimal supersymmetric standard model contains, in itself, the necessary particles to realize the Zee mechanism. In supersymmetric theories each ordinary particle has a superpartner; the companion of the fermions is a boson, and the companion of the bosons is a fermion. The role of the charged singlet scalar is taken by the singlet sleptons, the bosonic superpartners of the right-handed charged leptons. As a doublet scalar one can

Figure 5: The graphs generating the electromagnetic coupling of the neutrinos in the minimal supersymmetric standard model

use the respective doublet (right-handed) slepton. The singlet and doublet sleptons mix (with mixing angle ξ) because of soft supersymmetry breaking terms. We also have to break R-parity [29] since it forbids Majorana masses and magnetic moments. We do this explicitly by appropriately chosen trilinear coupling constants. One can write the mass matrix elements in analogy to Eq. (11),

$$m_{ik} = \frac{\lambda_{ikl}\lambda_{jlk} + \lambda_{ilk}\lambda_{jkl}}{32\pi^2} m_l (M_2^2 - M_1^2) I(m_l, M_1, M_2) \sin 2\xi. \quad (24)$$

where λ_{ijk} is a trilinear coupling of the superfields, here acting as a Yukawa coupling. The matrix is antisymmetric for the exchange of the first two indices. In contrast to the plain Zee model, the diagonal elements can be non-zero since we do not impose any symmetry condition with respect to the last component in the coupling λ_{ijk} . Moreover, the neutrino mass depends only linearly on the charged lepton mass.

The transition magnetic moment (Fig. 5) is given by

$$\mu_{ik} = \frac{\lambda_{ikl}\lambda_{jlk} - \lambda_{ilk}\lambda_{jkl}}{32\pi^2} m_\epsilon m_l [I(M_1, m_l, m_l) - I(M_2, m_l, m_l)] \sin 2\xi \mu_B. \quad (25)$$

It has been shown [30, 31, 32] that in this model one can simultaneously have one 29 eV Majorana neutrino and two lighter states, with substantial transition moments between them. In principle the heavy state can be made (dominantly) of either muon or tau neutrinos, but the tau neutrino would be disfavored since an ugly set of parameters would be required to satisfy all the necessary constraints.

The above model is still in danger of conflicting with the neutrino oscillation experiments. The theoretically more plausible scenario with the muon neutrino as the heaviest state is unacceptable. With the minimal particle contents one cannot realize mixing without extending the Barr-Freire-Zee mechanism.

The Barr-Freire-Zee mechanism can be applied, though, by introducing two new doublets [33]. By an appropriate choice of the terms violating lepton number it is possible to reproduce a convenient mass matrix satisfying the oscillation experiments and having the desired transition moments [34]. This is also consistent with the stringent constraints for the supersymmetric couplings.

3.5 Dirac neutrinos

Although a strong prejudice favors neutrinos to be Majorana particles, there are at least some reasons to assume the contrary. In several superstring inspired models (like E_6) there is no way to impose Majorana masses, so neutrinos are necessarily Dirac particles if they are massive.

The only attractive way to have light Dirac masses is to forbid the tree-level mass and generate the masses radiatively. As we have already seen, one can then also obtain a magnetic moment from the same loops. Assuming two mixing charged singlet scalars, typically we have

$$m_{ij} = \sum_l \frac{f_{il}^L m_\ell f_{lj}^R}{64\pi^2} \sin 2\beta (M_1^2 - M_2^2) I(m_\ell, M_1, M_2) \quad (26)$$

and

$$\mu_{ij} = \sum_l \frac{m_\ell f_{il}^L m_\ell f_{lj}^R}{64\pi^2} \sin 2\beta [I(m_\ell, m_\ell, M_1) - I(m_\ell, m_\ell, M_2)] \mu_B. \quad (27)$$

where f are Yukawa couplings and β is the scalar mixing angle. In contrast to the Majorana case where the transition moment matrix was explicitly antisymmetric, and hence involves only non-diagonal elements, the magnetic moment matrix of Dirac neutrinos is symmetric. The non-diagonal elements are very small, hence, the transition moments are additionally suppressed as compared to the diagonal moments. Because of this phenomenon, the radiative decay amplitudes of Dirac neutrinos would be very small, even though they would have sizable diagonal magnetic moments. Thus, a trivial one-loop alone cannot yield any observable decay rate.

The suppression of the transition moment is reduced if there are several different loops contributing differently to the mass and to the magnetic moment. Of particular interest are models in which the mass is suppressed with respect to the magnetic moment, e.g., by the cancellation of terms [35, 36, 37, 38]. In that case the transition moment can be sizable enough, and in addition to causing the radiative decay, there might also be magnetic moments causing the neutrino spin-flip in the Sun.

So far the models presented have been designed with the solar neutrino problem in mind, and the non-diagonal elements have been ignored or even forbidden to avoid flavor changing effects. There is not, however, any fundamental reason to impose any flavor dependent symmetry, and with a careful construction of the model one should be able to obtain a transition moment of $10^{-14} \mu_B$ without giving rise to intolerably large flavor violating effects.

4 Observations

4.1 General approach

The Milky Way Galaxy is a highly flattened ellipsoid whose dimensions and rotational characteristics are well determined by radio and optical studies. The rotational characteristics can not be the product of gravitation forces due to the luminous matter in the Galaxy; a standard solution is to hypothesize an extended halo of dark matter. For definiteness, we adopt here the detailed Galactic halo model of Overduin et al. [39], which if populated by neutrinos consistent with the Sciama hypothesis, would produce the observed rotation curve. This model is a flattened ellipsoid having a scale radius of 8 kpc, a scale height of 3 kpc, and a local neutrino density of $5 \times 10^7 \text{ cm}^{-3}$. With a decay lifetime of 10^{23} s these assumptions lead to a local volume emissivity E of $5 \times 10^{-16} \text{ photon/cm}^3 \text{ s}$.

Absorption by the interstellar medium in the EUV and FUV is crucial in understanding the observational constraints on decaying neutrino populations. The attenuation cross section of the neutral interstellar medium reaches its maximum of about $6 \times 10^{-18} \text{ cm}^2$ per hydrogen atom in the EUV band [40]. For this reason any confrontation of theory with observations demands a careful treatment of combined attenuation and emission. Our approach is to quantify the expected appearance of the sky assuming the existence of Sciama type neutrinos, and to compare these results with observational limits.

The spatial and spectral signature of any volume emission process can be characterized in terms of the distribution of absorbing and emitting material along the line of sight. The photon emissivity per unit volume FOR one photon emitted per neutrino decay is equal to the ratio of the neutrino number density to the neutrino decay lifetime t_ν . We let n denote the number of absorbing atoms per unit volume and σ be the total attenuation cross section. Then $A = n\sigma$ represents the linear attenuation coefficient, disposed in some manner along the line of sight, so that the optical depth to any given point at distance r is

$$\tau(R) = \int_0^R A dr.$$

The observed total photon intensity in the line is given by

$$I = \frac{1}{4\pi} \int \exp(-\tau(r)) E(r) dr. \quad (28)$$

Extragalactic decaying neutrinos may also be detectable. The absorption cross section of neutral hydrogen becomes enormously smaller for wavelengths greater than the Lyman edge at 91.1 nm and the highest terms of the Lyman series lines 91.1–91.5 nm. These wavelengths correspond to a redshift of the order of $z = 0.03$ for a 28 eV neutrino and a 14 eV decay photon. This fact permits populations with these, or larger, redshifts to be viewed from our vantage point within our Galaxy in spite of our Galaxy's interstellar absorbing material. The local opacity-dominated integral of Eq. (28) must be replaced by an integral over redshift for extragalactic distances, which in a standard Friedman

universe gives a continuum photon intensity of

$$\frac{dI}{d\lambda}(\lambda) = \frac{ER\lambda_\epsilon^{3/2}}{4\pi\lambda^{5/2}} \frac{1}{\sqrt{1 + (2q_0 - 1)(1 - \frac{\lambda_\epsilon}{\lambda})}} \exp\left(-\frac{t}{t_\nu}\right) \quad (29)$$

where $I(\lambda)$ is the observed flux of photons per unit wavelength at $\lambda > \lambda_\epsilon$, λ_ϵ is the emitted wavelength, E is the volume photon emissivity at a point along the line of sight whose redshift is $z = (\lambda/\lambda_\epsilon) - 1$, $R = c/H_0$ is the Hubble radius, t is the age of the universe at redshift z , and q_0 is the deceleration parameter of the universe [41]. Here we take Hubble's constant $H_0 = (100 h) \text{ km s}^{-1} \text{ Mpc}^{-1}$, where $0.4 < h < 0.8$.

If neutrinos concentrate in large clusters of galaxies at redshifts beyond about 0.1, their spatial nonuniformity would lead to a characteristic spectral emission signature at the redshift of the cluster.

Because of the very different distance scales involved, the observations divide into two kinds: EUV ($\lambda < 91 \text{ nm}$) observations that constrain the emissivity from neutrinos within our Milky Way Galaxy, and FUV ($\lambda > 91 \text{ nm}$) observations that constrain extragalactic emissivity at redshifts ≥ 0.03 .

4.2 Component 1: emission from material in the local Solar neighborhood

The spatial distributions of source and absorber jointly control the integral (28) along each line of sight. A picture of our neighborhood of the Galaxy has emerged from detailed studies of the distribution of interstellar matter. From solar backscatter studies of neutral helium flowing through the solar system, it is certain that the solar neighborhood has an appreciable fraction of neutral hydrogen, whose density is approximately 0.1 cm^{-3} [42] over an extent of several parsecs (pc). We refer to this material as the local solar-neighborhood cloud and denote its number density n_1 .

To quantify the expected EUV line intensity from local neutrino decay, we evaluate THE integral (28) for a uniformly intermixed emitter and absorber and obtain

$$I = \frac{E}{4\pi n_1 \sigma}. \quad (30)$$

For the density of the local medium we adopt 0.1 cm^{-3} , giving a linear attenuation length of $1.7 \times 10^{18} \text{ cm} = 0.6 \text{ pc}$ at 90 nm. These figures combine to give a predicted local diffuse EUV line intensity of $60 \text{ photons cm}^{-2} \text{ s}^{-1} \text{ sr}^{-1}$.

4.3 Component 2: the hot interstellar medium

A variety of evidence including studies of the EUV and soft X-ray background and O VI measurements suggest the presence of a pervasive hot ISM component with $T = (0.5 \text{ to } 1.0) \times 10^6 \text{ K}$. In contrast to a neutral ISM, this plasma provides no absorption for 90 nm photons produced by the decay of 28 eV neutrinos. Attenuation would only occur in localized wisps of cooler matter containing neutral or partly ionized atoms. Accordingly, corridors

of high EUV transmission define long lines of sight along which neutrino decay emissivity could in principle integrate to significant values.

Studies in the visible [42, 43, 44, 45], and in the FUV [46, 47] show that some lines of sight to stars 50–200 pc distant have little additional absorbing material beyond that expected from the local solar neighborhood cloud. This is particularly true in the third quadrant of the Galactic plane, spanning from $l = 180^\circ$ to $l = 270^\circ$ from Orion to the Gum Nebula in Vela. The total column densities suggested including the local cloud are in some cases below $1 \times 10^{18} \text{ cm}^{-2}$, setting upper limits on the average neutral hydrogen density as low as 0.002 cm^{-3} over these distance scales.

Further evidence for the low attenuation in these directions comes from recent neutral hydrogen determinations along lines of sight to stars in the EUV by Vallerger et al. [48], and especially by Vennes et al. [49] who show from observations with the Extreme Ultraviolet Explorer satellite that many third-quadrant stars between 33 and 206 pc distant exhibit similar values of ISM absorption column density, and that, therefore, in these direction, only the local solar-neighborhood cloud contributes to this absorption.

To model the observed intensity expected in the view directions corresponding to hot ISM corridors, we add to the local intensity a line-of-sight contribution beginning at $r_1 = 1.7 \text{ pc}$ extending to a much larger distance r_2 characteristic of the depth of the corridor, for example 300 pc. In this region, the attenuation is determined by an average smoothed-out neutral atom-wisp density of 10^{-3} cm^{-3} . For the realistic situation $r_1 n_1 \sigma \gg 1$ and $r_2 n_2 \sigma \gg 1$, this contributes

$$I = \frac{\exp(-n_1 \sigma r_1)}{4\pi n_1 \sigma} E \quad (31)$$

This predicted flux in the direction of the hot ISM corridors is substantially greater than the flux in the case of the local medium alone, in spite of attenuation by the local solar-neighborhood cloud, because of the factor of 100 greater photon mean free path in the hot ISM corridors. This model gives 300 photon units for a volume emissivity of $E = 5 \times 10^{-16} \text{ photons cm}^{-3} \text{ s}^{-1}$ in view directions aligned with the hot ISM corridors.

The most sensitive upper limits that have been placed on diffuse EUV line emission in the 80–90 nm band are those from Voyager [50]. A reanalysis of the Voyager data by Edelstein & Bowyer [51] has established an upper limit of $12000 \text{ photons cm}^{-2} \text{ s}^{-1} \text{ sr}^{-1}$, in this band, about a factor of 40 higher than the best-case intensity expected from component 2 and a factor of 200 above the more typical component 1 expectations.

4.4 Component 3: Galactic halo

In this section we consider whether decay photons from our Galaxy's neutrino halo population would be observable from the vantage point of the Sun located at our Galaxy's midplane. To address this question we must examine the total column of absorbing matter along lines of sight that extend to the outer edge of our Galaxy. Although EUV spectroscopy of extragalactic objects is providing data on this point, [43], at present the most extensive information comes from velocity-integrated 21 cm neutral hydrogen (H I) surveys [52, 53]. These reveal that the lowest observed Galactic hydrogen columns are found in Canis Major, Cancer, Lynx, and Ursa Major. Ursa Major includes the specific

region identified by Dickey and Lockman [52] for which the velocity-integrated Galactic H I column reaches an all-sky minimum of $4 \times 10^{19} \text{ cm}^{-2}$, a column density that constitutes an absorption of $\exp -200$ to the decay photons: this eliminates the possibility of directly observing halo decay photons in the 80–91 nm band. Thus, regardless of the extent of the neutrino halo of the Galaxy, no direct observational limit on the halo neutrino decay luminosity can be set by EUV observations. However, the escaping halo radiation may well be directly visible from other galaxies and (especially) clusters of galaxies at distances beyond $z = 0.03$, as discussed below.

4.5 Component 4: galaxies and clusters of galaxies, $z < 0.03$

Since the gravitational potential of large clusters may significantly concentrate a population of massive neutrinos, spectroscopy or photometry of diffuse emission from these clusters provides upper limits to neutrino decay yielding photons with wavelengths longer than 91 nm. However, any 80–91 nm emission is unobservable by direct EUV observations of clusters with this redshift owing to absorption by Galactic hydrogen, even in the most favorable scenario of the source neutrino halo of the cluster being fully exterior to the hydrogen content of its member galaxies.

The most massive of the nearby clusters of galaxies is the Coma cluster whose virial mass is $10^{15} M_{\odot}$. Its redshift of 0.022 is, however, insufficient to bring the putative decay photons into the FUV where they would be observable through the Galactic hydrogen. For this reason, searches for diffuse light from Coma [54, 55, 50] in the FUV and the visible do not bear on the issue of EUV radiative neutrino decay; although they certainly do bear on the issue of a baryonic component to the missing mass in Coma.

4.6 Component 5: uniform extragalactic, $z > 0.03$

We know from spectroscopy of quasi-stellar objects (QSOs) beginning with Gunn and Peterson [3] that intergalactic space is almost free of neutral hydrogen. Absorption features caused by neutral material are confined to discrete redshift systems along the line of sight and are not distributed continuously with respect to redshift. The most sensitive such determinations to date are the Hubble Space Telescope observations of nearby bright quasars [56].

Due to this clumping, decaying neutrino populations beyond $z = 0.03$ benefit from the transparency of the ISM at wavelengths longer than the Lyman edge. Therefore the source function integral (28) for this component must be extended over redshift as in Eq. (29) above.

The immediate consequence is to extend a diffuse decay emission line into a diffuse continuum whose power law coefficient is determined by the volume emissivity and the Hubble parameter, and whose power law index is close to -2.5 in photon units. Such a flux would appear as a spatially uniform, spectrally smooth, ultraviolet continuum rising steadily toward shorter wavelengths. From Eq. (29), the flux intensity at 140 nm should be $5000 \text{ photons cm}^{-2} \text{ s}^{-1} \text{ sr}^{-1} \text{ nm}^{-1}$, with only a slight dependence on cosmological model, for $t_{\nu}=10^{23} \text{ s}$.

Detecting an emission source that is spatially uniform and spectrally smooth is a considerable challenge, considering that observations must be conducted from within our Galaxy. Galaxies harbor stars, emission nebulae, and interstellar clouds that efficiently scatter ultraviolet photons. The review by Kimble et al. [41] sets an observational upper limit to diffuse extragalactic light at the level of $3600 \text{ photons cm}^{-2} \text{ s}^{-1} \text{ sr}^{-1} \text{ nm}^{-1}$, marginally consistent with the expected flux. The most sensitive available FUV diffuse spectroscopy is presently that of Martin et al. [57], which detected no positive extragalactic flux and yielded an upper limit of $2800 \text{ photons cm}^{-2} \text{ s}^{-1} \text{ sr}^{-1} \text{ nm}^{-1}$ at 140 nm. This flux requires $t_\nu > 1.8 \times 10^{23} \text{ s}$.

4.7 Component 6: clumped extragalactic, $z > 0.03$

Equation (29) expresses the observed intensity of ultraviolet light expected from a population of decaying neutrinos uniform over cosmological scales. Since a principal use of these neutrinos is to furnish the missing mass of clusters of galaxies, their decay emission must reveal them to be spatially nonuniform and to have a spectral (velocity) distribution of the missing mass of the particular cluster. Signatures of this kind are much more detectable than uniform fluxes since spatial and spectral on-off comparisons become possible.

The amount of emission expected depends on the degree of clumping. For a rich cluster of galaxies with mass $10^{15} M_\odot$, an emissivity equal to $4 \times 10^{56} \text{ photons s}^{-1}$ is expected. At a redshift $z = 0.2$ this would yield an observable flux in the direction of the cluster with an intensity of $2 \text{ photons cm}^{-2} \text{ s}^{-1}$. The observed emission line should have an observed wavelength $\lambda = \lambda_\epsilon(1 + z)$ corresponding to the redshift of the cluster and have a breadth corresponding to the cluster's velocity dispersion.

Davidson et al. [58] used the Hopkins Ultraviolet Telescope (HUT) to observe the rich cluster of galaxies Abell 665. This cluster has a redshift of 0.181, bringing decay photons to a local wavelength of 106 nm where they are not strongly attenuated by Galactic gas. The velocity dispersion of the cluster, 1200 km s^{-1} , yields a central mass density of the order of $10^{-25} \text{ g cm}^{-3}$, most of which must be dark matter. They found no sign of an emission feature and set a minimum lifetime for cluster-binding neutrinos of $3 \times 10^{24} \text{ s}$, assuming that all the dark matter is neutrinos.

Fabian et al. [59] examined International Ultraviolet Explorer (IUE) spectra of the cluster of galaxies surrounding the quasar 3C263 for evidence of emission. At a redshift of $z = 0.646$ this cluster does not have a well-determined velocity dispersion or total mass. On the basis of well-studied clusters, they adopted a surface mass density of 1 g cm^{-2} . No emission line was detected, and they were able to place a lower limit to the lifetime of the neutrino of $2 \times 10^{23} \text{ s}$ if the cluster mass is dominated by neutrinos.

Both the IUE and the HUT workers point out that these limits are weakened if absorbing material intervenes. Neutral hydrogen is a strong absorber of 15 eV photons, and dust is a strong absorber of hydrogen recombination radiation. Consequently the presence of small amounts of neutral gas and dust could serve to obscure any decay radiation. Moreover, X-ray spectroscopy of nearby clusters of galaxies has revealed that substantial amounts of neutral diffuse matter can be found in clusters [61]; even a small fraction of this material could destroy the emission signature [59].

5 Conclusions

We have studied the theoretical possibilities and observational constraints for the existence of a neutrino with the properties conjectured by Sciama [8]: a mass near 28 eV and a decay mechanism yielding a massless neutrino and a photon in the 89-91 nm extreme-ultraviolet wavelength range. If such a neutrino type exists, it would have been produced copiously in the early universe and would play a cosmologically significant role as diffuse hot dark matter. If these neutrinos are able to coalesce within galaxies or clusters of galaxies, they could constitute a portion of the widely-observed missing mass in these systems and their decay could explain the anomalously high degree of ionization of hydrogen within our Milky Way Galaxy.

Neither the standard model of electroweak particle interactions nor any trivial extensions can lead to radiative neutrino decays with sufficiently short time scales to accommodate the Sciama scenario. We have therefore investigated a variety of more elaborate models; these are able to accommodate the required decay rates. The proposed models are based on radiative generation of the electromagnetic coupling and mass matrices, and require new scalar particles and in some cases new symmetries. The simplest viable scenario is a hybrid model combining the see-saw mechanism with the Zee model. The consistency of that model is marginal, but other models with more elaborated ingredients can improve the situation. A supersymmetric extension may provide the necessary ingredients, but making the model plausible requires a few special assumptions.

The models predict new phenomena that can be searched for experimentally. The possible signatures include the existence of new particles, neutrino oscillations, neutrinoless double beta decay, and some model-dependent corrections to the properties of the charged fermions. Many of the simplest models explored here can be confirmed or rejected by forthcoming experiments, but some will be difficult to verify experimentally in the foreseeable future.

The existence of these decaying neutrinos could be established astronomically if their decay photon signature could be unambiguously detected. The expected signature is strongly dependent on the degree of coalescence of neutrinos with ordinary matter in the present Universe. The possibility of such capture hinges on the dynamics of galaxy formation [13], which is not well understood.

We have examined three scenarios which can accommodate the neutrinos postulated by Sciama which are consistent with existing observational constraints:

- 1) If massive neutrino halos surround galaxies and clusters of galaxies, they could supply the missing mass needed to explain galactic rotation curves and cluster velocity dispersions. The emission signature of their decay would then be an emission line at the redshifted wavelength of the decay photon, with a spatial profile similar to the gravitational potential of the cluster. In this scenario, the ultraviolet spectrum of the massive cluster of galaxies Abell 665 [58], which shows no emission line, sets a lower limit to the decay lifetime of 3×10^{24} s.

- 2) If massive neutrinos do not appreciably coalesce, their decay signature is spatially uniform but smeared out by redshift into a smooth spectrum that declines rapidly with increasing wavelength. Existing spectroscopic measurements of the diffuse cosmic ultraviolet background [57] presently set a lower limit to the decay lifetime of 1.8×10^{23}

s.

3) If massive neutrinos coalesce into galaxies and share a common spatial distribution with ordinary baryonic matter, then attenuation by photoabsorption by dispersed material within the galaxies would weaken or eliminate the Abell 665 lower limit. In this picture the only direct observational approach would be a search for nonredshifted EUV line emission from material in the local solar neighborhood. Existing upper limits, obtained from a spectrometer with 3 nm resolution, are more than one to two orders of magnitude larger than the predicted flux. A spectrometer now under development at the University of California in collaboration with the Instituto Nacional de Tecnica Aeroespacial employs a new spectrometer concept with a grating which both focuses and disperses EUV diffuse radiation. This spectrometer, in combination with a newly developed very low background detector, has a spectral resolution of ~ 0.2 nm and a minimum flux detection capability well below that predicted by the Sciama scenario. It will be flown as part of the scientific payload of the Spanish Minisat satellite to be launched in early 1996.

ACKNOWLEDGMENTS

The authors gratefully acknowledge helpful comments from Juan Perez-Mercader, Paolo Salucci, David Schramm, Dennis Sciama, Goran Senjanovic, Alexei Smirnov, and Gary Steigman. Stuart Bowyer acknowledges the support of NASA grant NGR05-003-450. Mike Lampton acknowledges the support of NASA grant NAGW-3823. Juha Peltoniemi acknowledges the support of a postdoctoral fellowship from Istituto Nazionale di Fisica Nucleare.

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