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

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Massive relic neutrinos may contribute significantly to the density of the Universe, and if appropriately concentrated, could explain the 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 nonstandard particle models with an extended scalar sector or minimal supersymmetry have sufficient freedom to accommodate this scenario. We derive observational constraints on these neutrinos 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 do not rule out the Sciama scenario; they only rule out specific distributions of radiatively decaying neutrinos in clusters of galaxies.

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I. 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, 160 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

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

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 $\mathcal{E}_\gamma = 13.6$ 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

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

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 [4–7]. 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

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

This value for m_ν corresponds to a photon wavelength of 89–91 nm. Sciama assumes that the lighter species has a mass $m_1 < 3$ eV, thus justifying the approximation in Eq. (1).

This massive neutrino could well be the τ neutrino because particle physics provides a much larger upper limit to its mass [9]. There are very few alternative particles which could be used to explain the hydrogen ionization in our Galaxy or in the Universe at large by radiative decay. 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, such as 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 primordial neutrinos is very large, $N_\nu \simeq 112$ neutrinos per cm^3, the mass represented by Eq. (3) would give an appreciable contribution to the total mass density of the universe:

$$m_\nu = \Omega_\nu \frac{\rho_c}{N_\nu} \simeq 95\Omega_\nu h^2 \mathrm{eV}. \qquad (4)$$

Here Ω_ν is the ratio of the present-day neutrino density to the critical mass density $\rho_c = 3H_0^2/8\pi G = 10.6h^2$ keV/cm^3, and Hubble's constant is $H_0 = (100h)\mathrm{km\,s^{-1}\,Mpc^{-1}}$, where we take $0.4 < h < 0.8$. 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.

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. Sandage [12] review the many observational approaches to constraining the present-day density parameter $\Omega_0 = \rho_0/\rho_c$ 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. One is then led to the conclusion that 90% of the matter in the Universe is dark and its composition unknown. One of the attractive features of massive neutrinos as hot dark matter (HDM) is that they could provide a straightforward solution to the problem of galactic rotation curves.

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 Sec. II we describe the reasons for this lifetime in the standard model and discuss various nonstandard models in which the neutrinos could have other interactions that would speed up the radiative decay considerably. We show that models with properties required for the Sciama neutrino can be consistently constructed. We do not address the problem of the solar electron neutrino deficit or the atmospheric muon neutrino deficit, which require some other explanation.

In Sec. III 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 Sec. IV we summarize our results and discuss experiments that will be able to confirm or reject the neutrino hypothesized by Sciama.

II. RADIATIVE NEUTRINO DECAY IN NONSTANDARD MODELS

A. 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 electrons, but there would be no explanation for the smallness of their mass. The simplest explanation would be the seesaw scenario: the right-handed components have a very large Majorana mass term (a term connecting neutrinos and antineutrinos), as a result of which there will be two mass eigenstates (per family), 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 Sec. II E.

Radiative decay of massive neutrinos is caused by a transition magnetic moment μ_{ij}, antisymmetric for Majorana neutrinos [13], with a rate

$$\Gamma(\nu_i \to \nu_j \gamma) = \frac{|\mu_{ij}|^2}{8\pi} m_i^3 \left[1 - \left(\frac{m_j}{m_i}\right)^3\right]. \qquad (5)$$

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

$$\mu_{ij} \sim (0.5 - 1) \times 10^{-14} \mu_B, \qquad (6)$$

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.

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

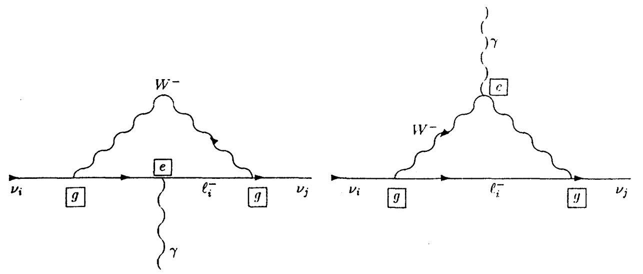

FIG. 1. The diagrams generating radiatively the magnetic moment in the standard model with massive Dirac neutrinos.

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. The magnetic moment of a massive neutrino in the standard model is, however, very small, $3\times10^{-19}\mu_B(m_\nu/\mathrm{eV})$, 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 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 $10^{-11}\mu_B$ [14,15].

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 [16,17] radiatively produces both the mass and the magnetic moment with a scalar exchange loop. In its simplest form this model cannot consistently satisfy Sciama's conditions. A hybrid model involving both the Zee model and the seesaw mechanism is much better, as shown later. 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 [17]. 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.

Models explaining even higher disparities with neutrino masses and magnetic couplings using specifically designed symmetries have been explored [18] in the context of the spin-flip solution to the solar neutrino problem. Given that the requirements for the solar neutrino rotation are magnitudes stronger than for Sciama's neutrinos, those models may be applied to our case without too much difficulty. Typically the custodial symmetry is arranged so that it allows the magnetic moment but forbids the mass terms in a perturbative picture causing the loops contributing to the mass to cancel. The small masses could then arise controllably by a small breaking of the symmetry. It has also been suggested that both the mass and the magnetic moments arise as a result of the symmetry breakings by nonzero vacuum expectation values of separate Higgs field representations, so that they are effectively independent [19]. Because of the large variety of possible models and submodels, and the number of degrees of freedom in each of them, we do not consider them here in more detail. They may, however, turn out to be important in building a model reconciling Sciama's scenario with other phenomena, such as solar or atmospheric neutrino deficits.

In the following subsections we will explore some models in more detail. We will show explictly that models that can incorporate the neutrino hypothesized by Sciama do exist which are consistent with all present observational constraints.

B. 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 nonzero 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

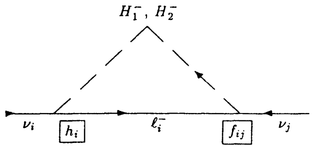

FIG. 2. The diagram generating neutrino mass radiatively in the Zee model.

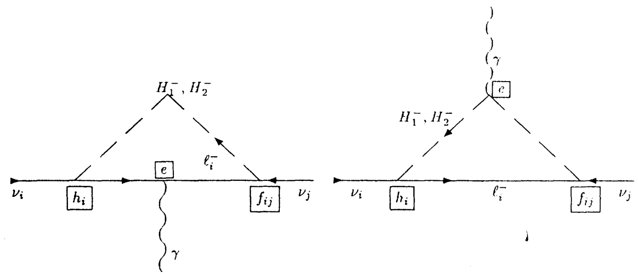

FIG. 3. The graphs generating the electromagnetic coupling to the neutrinos in the Zee model.

neutrino mass matrix in the flavor basis (i.e., the physical eigenbasis of the charged leptons) [20,21,15]

$$m_{ik} = \frac{f_{ik} g \sin 2\beta \cot\alpha}{32\pi^2 M_W}(M_2^2 - M_1^2)[m_i^2 I(m_i, M_1, M_2) \\ -m_k^2 I(m_k, M_1, M_2)], \tag{7}$$

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 ith family, M_1 and M_2 are the masses of the physical Higgs scalars, and

$$I(m, M_1, M_2) = \frac{1}{M_2^2 - M_1^2}\left(\frac{M_1^2}{M_1^2 - m^2}\ln\frac{m^2}{M_1^2} - \frac{M_2^2}{M_2^2 - m^2}\ln\frac{m^2}{M_2^2}\right) \\ \approx \frac{1}{M_2^2 - M_1^2}\ln\frac{M_2^2}{M_1^2}. \tag{8}$$

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. (4) if the heavy states have the Sciama mass [Eq. (3)]. (The general phenomenology of this model has been recently studied in Ref. [22].)

A magnetic moment arises at the one-loop level by coupling a photon to the internal lines as shown in Fig. 3. Now one has [15,20]

$$\mu_{ij} = \frac{g m_e f_{ij}}{32\pi^2 M_W}\sin 2\beta \cot\alpha \mu_B \\ \times[m_i^2\{I(M_1, m_i, m_i) - I(M_2, m_i, m_i)\} \\ +m_j^2\{I(M_1, m_j, m_j) - I(M_2, m_j, m_j)\}], \tag{9}$$

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

$$I(M, m, m) = \frac{M^2}{(M^2 - m^2)^2}\ln\frac{M^2}{m^2} - \frac{1}{M^2 - m^2} \\ \approx \frac{1}{M^2}\ln\frac{M^2}{m^2} - \frac{1}{M^2}. \tag{10}$$

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

$$\mu_{ij} = m_{ij} m_e \frac{I(M_1, m_i, m_i) - I(M_2, m_i, m_i)}{(M_2^2 - M_1^2) I(m_i, M_1, M_2)}\mu_B\,. \tag{11}$$

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

$$\mu_{\tau l} \simeq \left(\frac{m_{\tau l}}{1\ \text{eV}}\right)\left(\frac{45\ \text{GeV}}{M}\right)^2 10^{-15}\mu_B, \tag{12}$$

where $l = \mu, e$. 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 which conflicts 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 is the symmetric two-loop of Babu [23] which 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.

A better model may be obtained by using two different mechanisms, one to generate a diagonal mass element for the heaviest neutrino and the other to generate nondiagonal elements and transition moments. Combining the Zee and Babu models might achieve this, though it may seem odd (but not necessarily unnatural) to have the two-loop result exceed the one-loop result. Indeed, there is a variation of the Babu model having the particle content of the Zee model and containing sufficiently undetermined coupling constants to allow the tuning of an appropriate mass matrix. A more plausible alternative may be to combine the Zee model with seesaw generated masses.

Assume now that we have one heavy (28 eV) neutrino with its mass generated either by the seesaw 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 τ neutrino, the nondiagonal mass elements are bound by laboratory experiments in the following way [24]:

$$m_{e\tau} < 3.8 \text{ eV} \quad (\Leftarrow \sin^2 2\theta_{e\tau} < 0.07) \tag{13}$$

and

$$m_{\mu\tau} < 0.9 \text{ eV} \quad (\Leftarrow \sin^2 2\theta_{\mu\tau} < 0.004), \tag{14}$$

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

$$m_{e\mu} < 0.7 \text{ eV} \quad (\Leftarrow \sin^2 2\theta_{e\mu} < 0.0025). \tag{15}$$

Hence the only sizable transition moment can arise between τ and electron neutrinos. If we assume that there are no accidental cancellations between different components, Eq. (12) gives an upper limit for the transition moment:

$$\mu_{e\tau} < 4 \times 10^{-15} \mu_B. \tag{16}$$

This may be only marginally consistent with Sciama's requirement, although by allowing the nondiagonal mass terms from the Zee loop and the seesaw to partly cancel, one can obtain larger transition moments, up to one's tolerance for fine tuning. Consequently, one can predict that within this scheme the electron neutrino mass is about 0.5 eV, accessible for searches of neutrinoless double β decay. Future neutrino oscillation experiments will clarify whether these kinds of scenarios are viable.

C. Two-loop graphs

Going to higher orders in perturbation theory may improve the situation. Namely, within the plain Zee model there exist two-loop graphs contributing substantially to the magnetic moment. We have considered the graph depicted in Fig. 4 [25]; it includes a W gauge boson. This provides three benefits: (1) the coupling of the gauge boson to a fermion is much larger than the Yukawa couplings of the scalars; (2) the amplitude does not explicitly contain the lepton mass, since there is only one scalar coupling (the vertex with f_{ij}) causing the chirality flip; (3) there is no similar graph without the external photon that would contribute to the mass, because spin conservation forbids the effective coupling between two particles of different spin (here 1 and 0) [17]. Hence the higher order corrections do not essentially change the mass matrix given by the hybrid model.

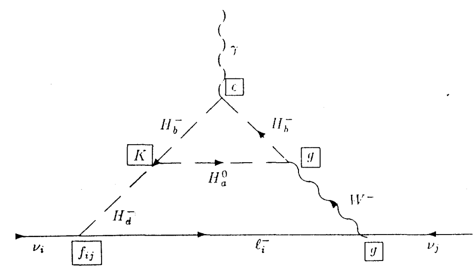

FIG. 4. The graph generating the electromagnetic coupling to the neutrinos in the Barr-Freire-Zee model.

The two-loop graph contributes to the magnetic moment as

$$\mu_{ij} \simeq \sum_{abcd} \frac{f_{ijd} U_{ab} K_{bcd} U^\dagger_{ca} g^2 m_e}{(16\pi^2)^2 M_W^2} \times I(M_W, M_a, M_b, M_d)\mu_B, \tag{17}$$

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 e-μ component of the transition moment can be enhanced by two 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 range of parameters in which the two-loop contribution will be the same or a few times larger than the one-loop contribution. Hence one can obtain transition moments

$$\mu_{\tau l} \lesssim \left(\frac{m_{\tau l}}{1 \text{ eV}}\right) 2 \times 10^{-15} \mu_B \tag{18}$$

or

$$\mu_{\mu e} \lesssim \left(\frac{m_{\mu e}}{0.3 \text{ eV}}\right) 10^{-14} \mu_B. \tag{19}$$

Thus we can now easily accommodate Sciama's value Eq. (6) even when the decaying particle is a muon. The expected mass scale for the electron neutrinos would now be 0.1–0.5 eV if the heavy state is the τ neutrino, or $(0.1 - 3) \times 10^{-2}$ eV for the channel $\nu_\mu \to \gamma\nu_e$.

Smaller neutrino mixings are possible in models that forbid the one-loop contribution to the mass. For instance, adding a third doublet to the above model, one can arrange a symmetry forbidding the direct self-coupling between the scalars that couple to fermions [17]. Hence the neutrino mass (mixing) can only appear at the two-loop level, being strongly suppressed. Since the magnetic moment arises also at the two-loop level, but is unsuppressed, one avoids all the constraints for neutrino mixing, and one gains more freedom for the theoretical parameters. Note that the 30 eV mass scale will have to be given by a different mechanism, like the seesaw.

D. Supersymmetric models

The minimal supersymmetric standard model contains, in itself, the necessary particles for radiative generation of masses. 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 use the respective doublet (right-handed) slepton. The singlet and doublet sleptons mix (with mixing angle ξ) because of soft (dimension less than or equal to 3) supersymmetry-breaking terms. We also have to break R parity (a discrete symmetry required to avoid fast proton decay in supersymmetric models) [26] 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. (7):

$$m_{ik} = \sum_{lj} \frac{\lambda_{ikl}\lambda_{jlk} + \lambda_{ilk}\lambda_{jkl}}{32\pi^2} m_l (M_2^2 - M_1^2) I(m_l, M_1, M_2) \sin 2\xi, \tag{20}$$

where λ_{ijk} is a trilinear coupling of the superfields, here acting as a Yukawa coupling. The matrix is antisymmetric in the exchange of the first two indices. In contrast with the plain Zee model, the diagonal elements can be nonzero since we do not impose any symmetry condition with respect to the last component in the coupling λ_{ijk}. Hence we can obtain a single 30 eV Majorana neutrino within this model, with no need for additional mass terms by other mechanisms. Moreover, the neutrino mass depends only linearly on the charged lepton mass.

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

$$\mu_{ik} = \sum_{lj} \frac{\lambda_{ikl}\lambda_{jlk} - \lambda_{ilk}\lambda_{jkl}}{32\pi^2} m_e m_l [I(M_1, m_l, m_l) - I(M_2, m_l, m_l)] \sin 2\xi \mu_B. \tag{21}$$

It has been shown [27–29] 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 τ neutrinos, but the τ neutrino would be disfavored since an ugly set of parameters would be required to satisfy all the necessary constraints. However, the scenario with the muon as the heaviest state is inevitably in conflict with the neutrino oscillation experiments.

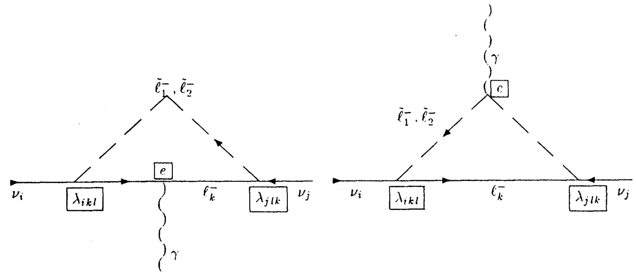

FIG. 5. The graphs generating the electromagnetic coupling of the neutrinos in the minimal supersymmetric standard model.

With the minimal particle content one cannot realize the Barr-Freire-Zee mechanism. To apply it one has to introduce two new doublets [30]. It has been shown [31] that in such a case one can choose the lepton number violating terms in the Lagrangian so that the mass matrix satisfies the oscillation experiments and the transition moments are sufficiently large. This is also consistent with the stringent constraints for supersymmetric couplings.

E. 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 (such as 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

$$m_{ij} = \sum_{l} \frac{f_{il}^L m_l f_{lj}^R}{64\pi^2} \sin 2\beta (M_1^2 - M_2^2) I(m_l, M_1, M_2) \tag{22}$$

and

$$\mu_{ij} = \sum_{l} \frac{m_e f_{il}^L m_l f_{lj}^R}{64\pi^2} \sin 2\beta [I(m_l, m_l, M_1) - I(m_l, m_l, M_2)] \mu_B, \tag{23}$$

where f are Yukawa couplings and β is the scalar mixing angle. In contrast with the Majorana case where the transition moment matrix was explictly antisymmetric, and hence involves only nondiagonal elements, the magnetic moment matrix of Dirac neutrinos is symmetric. The nondiagonal 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., when the mass terms by different loops cancel [32–34]. Apart from fine-tuning, this can be achieved by imposing an appropriate symmetry [e.g., SU(2) or SU(3)] between the left- and right-handed components. 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 nondiagonal 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.

F. Radiative neutrino decay in nonstandard models: Summary

We have studied the theoretical possibilities 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. The decaying neutrino can be either a muon or τ flavor, with only the electron neutrino being excluded. Consequently, the daughter neutrino can be of any flavor. Plausible models can be constructed for each case.

Although 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, the more elaborate models which we have investigated here are able to accommodate the required decay rates with a suitable choice of parameters. These parameters are not constrained by other physics, and thus there is no conflict with the Sciama scenario.

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 radiative model, the Zee model, leads most naturally to a neutrino mass spectrum consisting of one light state and two heavier states, almost degenerate with very large mixing. This, however, conflicts with the cosmology bound of Eq. (4) if the heavy states have the Sciama mass [Eq. (3)]. However, a hybrid model combining the seesaw mechanism with the Zee model provides a consistent solution. At the one-loop level the consistency of this model is still marginal, because it appears to conflict with the neutrino oscillation experiment. The model involves, however, two-loop graphs that can give a sufficiently large contribution to the magnetic moment to avoid the conflict, without fine-tuning.

We also demonstrate that more sophisticated models with expanded particle contents or specific symmetries can be designed to satisfy all the requirements of the Sciama neutrino. A supersymmetric extension can provide the necessary ingredients, but making the model plausible requires a few special assumptions. We note that in this scenario it is easier to accommodate the neutrinos as Majorana particles, since with Dirac neutrinos one needs more elaborate ingredients.

III. OBSERVATIONS

A. General approach

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 observations divide into two kinds: EUV ($\lambda < 91$ nm) observations that constrain the emissivity from neutrinos within our Milky Way Galaxy, and FUV ($\lambda > 91$ nm) observations that constrain extragalactic emissivity at redshifts ≥ 0.03.

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 cannot be solely the product of gravitation forces produced by the observed luminous matter in the Galaxy; a standard solution is to hypothesize an extended halo of dark matter. If the dark matter is the Sciama neutrino, the halo will, in fact, be luminous, but the decay luminosity of the neutrinos will be many orders of magnitude below the overall luminosity of the Galaxy since the neutrino lifetime is far greater than the age of the universe.

It is now known if large concentration factors (of the order of one million) of primordial neutrinos actually occur. Steigman *et al.* [35] 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 and Gunn [36]) 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

$$m_\nu > [0.5\rho(h/v)^3]^{1/4} \tag{24}$$

which is compatible with the Sciama neutrino.

For definiteness, we adopt here the detailed Galactic halo model of Overduin *et al.* [37], 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 5×10^7 cm^{-3}. With a decay lifetime of 10^{23} s these assumptions lead to a local volume emissivity E of 5×10^{-16} photon/cm^3/s.

Absorption by the interstellar medium in the EUV and FUV is crucial in understanding the observational constraints on a decaying neutrino population. The attenuation cross section of the neutral interstellar medium reaches its maximum of about 6×10^{-18} cm^2 per hydrogen atom in the EUV band [38].

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

$$I = \frac{1}{4\pi} \int \exp[-\tau(r)] E(r) dr. \tag{25}$$

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. (25) must be replaced by an integral over redshift for extragalactic distances, which in a standard Friedmann universe gives a continuum photon intensity of

$$\frac{dI}{d\lambda}(\lambda) = \frac{ER}{4\pi} \frac{\lambda_e^{3/2}}{\lambda^{5/2}} \frac{1}{\sqrt{1 + (2q_0 - 1)(1 - \frac{\lambda_e}{\lambda})}} \times \exp\left(-\frac{t}{t_\nu}\right), \tag{26}$$

where $I(\lambda)$ is the observed flux of photons per unit wavelength at $\lambda > \lambda_e$, λ_e is the emitted wavelength, E is the volume photon emissivity at a point along the line of sight whose redshift is $z = (\lambda/\lambda_e) - 1, R = c/H_0$ is the Hubble radius, t is the age of the universe at redshift z, and q_0 is the deceleration parameter of the Universe [39].

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.

B. Component 1: Emission from material in the local solar neighborhood

The spatial distributions of source and absorber jointly control the integral (25) 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} [40] over an extent of several parsecs (pc) at least in some directions. 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 (25) for a uniformly intermixed emitter and absorber and obtain

$$I = \frac{E}{4\pi n_1 \sigma}. \tag{27}$$

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

C. 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 interstellar medium (ISM) component with $T = (0.5 \text{ to } 1.0) \times 10^6$ K. In contrast with 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 [40–43], and in the FUV [44,45] 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 $l = 180°$ to $l = 270°$ from Orion to the Gum Nebula in Vela. The total column densities suggested including the local cloud are in some cases below 1×10^{18} 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 Vallerga *et al.* [46], and especially by Vennes *et al.* [47] 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 $r_1 = 1.7$ 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

$$I = \frac{\exp(-n_1 \sigma r_1)}{4\pi n_1 \sigma} E. \tag{28}$$

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 $E = 5 \times 10^{-16}$ photons cm^{-3} 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 [48]. A reanalysis of the Voyager data by Edelstein [49] has established an upper limit of 12 000 photons cm^{-2} s^{-1} 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.

D. 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 [41], at present the most extensive information comes from velocity-integrated 21 cm neutral hydrogen (HI) surveys [50,51]. 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 [50] for which the velocity-integrated Galactic HI column reaches an all-sky minimum of 4×10^{19} 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.

E. 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 [52,53,48] 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.

F. Component 5: Uniform extragalactic, $z > 0.03$

We know from spectroscopy of quasistellar objects (QSO's) 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 have been made with Hubble Space Telescope observations of nearby bright quasars [54].

Because of 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 (25) for this component must be extended over redshift as in Eq. (26) 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. (26), the flux intensity at 140 nm should be 5000 photons cm^{-2} s^{-1} sr^{-1} nm^{-1}, with only a slight dependence on cosmological model, for $t_\nu = 10^{23}$ 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.* [39] sets an observational upper limit to diffuse extragalactic light at the level of 3600 photons $\mathrm{cm}^{-2}\ \mathrm{s}^{-1}\ \mathrm{sr}^{-1}\ \mathrm{nm}^{-1}$, marginally consistent with the expected flux. The most sensitive available FUV diffuse spectroscopy is presently that of Martin *et al.* [55], which detected no positive extragalactic flux and yielded an upper limit of 2800 photons $\mathrm{cm}^{-2}\ \mathrm{s}^{-1}\ \mathrm{sr}^{-1}\ \mathrm{nm}^{-1}$ at 140 nm. This flux requires $t_\nu > 1.8 \times 10^{23}$ s.

G. Component 6: Clumped extragalactic, $z > 0.03$

Equation (26) 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 comparison 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 4×10^{56} 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 photons $\mathrm{cm}^{-2}\ \mathrm{s}^{-1}$. The observed emission line should have an observed wavelength $\lambda = \lambda_e(1+z)$ corresponding to the redshift of the cluster and have a breadth corresponding to the cluster's velocity dispersion.

Davidsen *et al.* [56] 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 $10^{-25}\mathrm{g\,cm}^{-3}$, 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 3×10^{24} s assuming that all the dark matter is neutrinos.

Fabian *et al.* [57] 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 $\mathrm{g\,cm}^{-2}$. No emission line was detected, and they were able to place a lower limit to the lifetime of the neutrino of 2×10^{23} 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 would obscure or degrade any decay radiation. X-ray spectroscopy of nearby clusters of galaxies has revealed that substantial amounts of neutral diffuse matter can be found in clusters [58]; even a small fraction of this material could destroy the emission signature [57].

IV. 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 their decay could explain the anomalously high degree of ionization of hydrogen within our Milky Way Galaxy and a variety of other effects.

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 radiative model, the Zee model, does not yet provide a consistent solution, but a hybrid model combining the seesaw mechanism with the Zee model will do so. At the one-loop level the consistency of that model is still marginal, with an apparent conflict with the neutrino oscillation experiment. The model involves, however, two loop graphs that can give a sufficiently large contribution to the magnetic moment to avoid the conflict, without fine-tuning. There are also more sophisticated models with expanded particle contents or specific symmetries that can be designed to satisfy all the requirements. A supersymmetric extension may provide the necessary ingredients, but making the model plausible requires a few special assumptions.

The particle physics allows the decaying neutrino to be either a muon or τ flavor, with only the electron neutrino being excluded. Consequently, the daughter neutrino can be of any flavor. Plausible models can be constructed for each case. It is easier to accommodate the neutrinos as Majorana particles, since with Dirac neutrinos one needs more elaborate ingredients.

The existence of these decaying neutrinos can be established astronomically by their decay photon signature. The expected EUV signature in our Galaxy is strongly dependent on absorption by the local interstellar medium and varies with view direction. The expected redshifted FUV signature is dependent on the degree of coalescence of neutrinos with ordinary matter in the present Universe.

If massive neutrino halos surround clusters of galaxies, and supply the missing mass needed to explain cluster velocity dispersions, the signature of their decay would 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 [56], which shows no such emission line, sets a lower limit to the decay lifetime of 3×10^{24} s. If these 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.

A variety of tests can be used, in principle, to test the Sciama hypothesis. The particle models discussed herein predict new phenomena that can be searched for experimentally. The possible signatures include the existence of new particles, neutrino oscillations, neutrinoless double β decay, and corrections to the properties of 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 Sciama particle will have direct astronomical signatures. The diffuse FUV effects are sufficiently subtle that they will be virtually impossible to disentangle from other, more conventional, astronomical backgrounds. FUV effects in clusters of galaxies are much more straightforward to detect, but are subject to substantial uncertainties concerning the amount of absorption in these clusters.

A direct observational approach is to 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. An instrument now under development at the University of California in collaboration with the Instituto Nacional de Tecnica Aerospacial of Spain employs a new spectrometer concept with a grating which both focuses and disperses EUV diffuse radiation with 0.5 nm resolution. This spectrometer, in combination with a newly developed very low background detector, has 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 1996.

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