Production of single plasmons and photons by neutrinos in a medium

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(Received 7 October 1991)

A neutrino passing through a dielectric medium with an index of refraction greater than unity can emit a single photon into the medium, and, in a related process, a neutrino passing through a conductor or semiconductor can emit a single plasmon into the medium. The rates for other processes are calculated, taking a simple electron gas as the medium for the plasmon calculation. Two forms of coupling of the neutrino to the electrons in the medium are considered: (1) the left-handed current coupling of the standard model; (2) a neutrino magnetic moment coupling. For the case of the largest magnetic moments allowed by other considerations, the plasmon process could be on the verge of rendering solar "plasmon" electron neutrinos (with energies ≈ 2 keV) observable. The plasmon process also has some potential in the detection and identification of neutrinos of mass > 10 keV, and large magnetic moments.

PACS number(s): 14.60.Gh, 13.15.$-$f, 29.40.$-$n

I. INTRODUCTION

In the absence of interactions so exotic that they have not yet even been proposed in the literature, it would seem that direct experimental detection of neutrinos of energy, say, much less than 0.1 MeV may be impossible, at least in the foreseeable future. Yet there exist a number of motivations, in addition to that of simple completeness, for trying to penetrate this domain experimentally, for example, the investigation of the possibility of cosmological dark matter consisting of neutrinos in the eV region, the possibility of an abundant flux of solar (plasma) neutrinos in the keV region (particularly if a low-mass neutrino has a large magnetic moment), or the possibility of data giving direct evidence on neutrino masses.

There have been a number of attempts to find a way in which the properties of atomic or solid matter could enhance the cross sections for something observable to be produced by an incoming neutrino, a natural direction to look as one considers neutrino wavelengths on the order of interatomic spacings or larger [1]. To the best of our knowledge, there has been no great success in this direction, however.

The present work reports calculations to see whether or not the neutrino-stimulated production of a single photon or plasmon in a material, with no other excitation of the material, could have, *under any currently favored assumptions as to neutrino properties and for some neutrino energy,* one or more of the following characteristics: (a) a significantly greater reaction rate than that estimated from the neutrino-electron cross section multiplied by the density of electrons in the material, (b) a greater reaction rate than that resulting from the neutrino-electron interaction in the particular material medium under consideration [this would generally be a weaker demand than (a), above, since at low neutrino energies the effect of electron binding will be to reduce greatly the rate over that of (a)], and (c) a particularly distinctive signal, so that there is an opportunity to effectively discriminate against backgrounds.

Of course, to each item of this list should be added the phrase "and some reasonable possibility for observation, given expected fluxes from various sources." We shall find no great cause for optimism in our results. Yet the results will display phenomena more or less on the fringe of observability, in some cases, for the most favorable choice of parameters, and there could be improvements that would emerge from more sophisticated considerations.

We treat here the excitation, by neutrinos, of two familiar coherent excitations in matter, namely, photons and plasmons. We consider conditions under which a single excitation is emitted by the neutrino into the medium, a phenomenon which is just Cherenkov radiation in the case of the photon process, except that the photon is not produced directly from an electric charge of the radiating particle. There is no special name for the plasmon process, but it is completely analogous to the Cherenkov process for photons (as is the emission of plasmons from fast electrons passing through conductors or semiconductors).

The program described above is very similar to that undertaken in the work reported in Ref. [2], although we address somewhat different issues. Moreover, we are in significant disagreement with the results of Ref. [2] as they apply to the common questions raised in the two papers.

We shall principally consider situations in which the momentum transfer from the neutrino of the excitation is small compared with the momentum of the neutrino. In this case the energy loss of the neutrino is given by

$$\Delta E_\nu = \omega \approx \mathbf{q}\cdot\mathbf{v} \ , \tag{1.1}$$

where $(\omega,\mathbf{q})$ is the four-momentum transferred and $\mathbf{v}$ is the velocity $(|\mathbf{v}|\approx c=1)$ of the neutrino, which, unless stated otherwise, will be taken as an electron neutrino of very small mass. The condition (1.1) can be satisfied only when the excitation is spacelike, $|\mathbf{q}|>\omega$. Thus a photon can be emitted in an energy region in which the index of refraction is greater than unity, for the case of a dielectric medium. For the case of an electron plasma, it is a longitudinal plasmon, rather than the transverse excitation (photon), which has a spacelike dispersion relation for sufficiently large momentum. Its dispersion curve $\omega(\mathbf{q})$ begins, for $\mathbf{q}=0$, at the plasma frequency ω_p and, for the case of a nonrelativistic electron gas, only moves slightly upward as $|\mathbf{q}|$ increases past the value $|\mathbf{q}|=\omega$, above which the phase velocity of the excitation is less than c. At much larger $|\mathbf{v}|$, when the phase velocity decreases to the Fermi velocity, the plasmon curve enters the region of strong damping from particle-hole production. In between the two thresholds lies the kinematical region on which well-defined plasmons may be radiated from a particle moving at velocity c.

We require a production mechanism, as well as kinematics, which conserves energy. We shall consider two possibilities: (a) production through a neutrino magnetic moment and (b) production through the conventional coupling of electrons to left-handed neutrinos. Although the theory behind such calculations is well developed, we present it in the next section in a way which unifies the four calculations: production of photons or longitudinal plasmons by a magnetic moment or neutrino-electron coupling.

II. EMISSION VERTICES

Choosing units such that $\hbar=c=1$, we take two forms of coupling of the electron neutrino to a medium, ignoring all neutrino-nucleon couplings.

(A) The neutrino-electron interaction, which is written, in charge-retention form, keeping only the vector part of the electron current, as

$$\mathcal{L}_I=-(G_1/2)[\bar{\nu}_e\gamma^\mu(1-\gamma_5)\nu_e]\bar{\psi}_e\gamma_\mu\psi_e\ , \tag{2.1}$$

where, in the standard model, $G_1=2^{1/2}G(\frac{1}{2}+2\sin^2\theta_W)$ and $G=1.16\times10^{-5}m_p^{-2}$. The axial-vector part does not contribute at all to the production of a longitudinal plasmon, and an estimate indicates that its contribution to photon production is smaller than that of the vector part, itself very small, as we shall see.

(B) A neutrino magnetic-moment term

$$\mathcal{L}_I=10^{-10}\mu_{10}e(4m_e)^{-1}\bar{\nu}_e\sigma_{\alpha\beta}\nu_eF^{\alpha\beta}\ . \tag{2.2}$$

We calculate the matrix element of the electromagnetic potential operator $A_\mu(x)$ between initial and final neutrino states:

$$\begin{aligned}W_\mu(p,q)&=\langle p-q|A_\mu(0)|p\rangle\\&=[E_\mathbf{p}E_{\mathbf{p}-\mathbf{q}}]^{-1/2}D_{\mu\nu}(q)R^\nu(p,q)\ ,\end{aligned} \tag{2.3}$$

where $D_{\mu\nu}$ is the photon propagator in the medium and $q_\mu=(\omega,\mathbf{q})$ the four-momentum transfer to the medium. The vector R^ν is given by a lowest-order perturbation calculation, for the two cases (A) and (B) above, as

$$R_\mu^A=G_12^{-1}e^{-1}\Pi_{\mu\lambda}(q)\bar{u}_f(\mathbf{p}-\mathbf{q})\gamma^\lambda(1-\gamma_5)u_i(\mathbf{p}) \tag{2.4}$$

and

$$R_\mu^B(p,q)=10^{-10}\mu_{10}e(2m_e)^{-1}\bar{u}_f(\mathbf{p}-\mathbf{q})\sigma_{\mu\nu}q^\nu u_i(\mathbf{p})\ . \tag{2.5}$$

Here $\Pi_{\mu\nu}(q)$ is the proper polarization part for the electromagnetic potentials:

$$D_{\mu\nu}^{-1}(q)=q^2g_{\mu\nu}+\Pi_{\mu\nu}(q)+\zeta q_\mu q_\nu\ , \tag{2.6}$$

where the parameter ζ has no physical effect. We write

$$D_{\mu\nu}(q)=-\sum_\alpha\frac{r_\mu^{(\alpha)}(\mathbf{q})r_\nu^{(\alpha)}(\mathbf{q})}{\omega-\omega^{(\alpha)}(\mathbf{q})}+\cdots\ , \tag{2.7}$$

where $\omega=\omega^{(a)}(\mathbf{q})$ is the energy-momentum relation for species of excitation (a). We have taken the energies $\omega^{(a)}(\mathbf{q})$ to be real, thereby neglecting the absorption of the excitations by the medium. The rate of production of an excitation of type (a) is given by a golden rule in which the residue of the pole at $\omega=\omega^{(a)}(\mathbf{q})$ [in expression (2.3) for W_{μ}, with one factor $r^{(a)}(\mathbf{q})$ truncated] serves as the matrix element:

$$\begin{aligned}\Gamma^{(a)}=\frac{1}{4\pi^2}\int d^3q(E_\mathbf{p}E_{\mathbf{p}-\mathbf{q}})^{-1}\\\times\delta[\omega^{(a)}(\mathbf{q})-\mathbf{q}\cdot\mathbf{v}]|R^\mu r_\mu^{(a)}|^2\ .\end{aligned} \tag{2.8}$$

The argument in the δ function follows from assuming $|\mathbf{q}|\ll|\mathbf{p}|$ as in (1.1). Equation (2.8) can be verified by calculating the imaginary part of the neutrino-neutrino scattering amplitude that would result from attaching another factor of $(E_{p'}E_{p'-q})^{-1/2}R^\mu$ to expression (2.3). Note that in case (A) we can simplify (2.4) by using the relation

$$\Pi_\alpha^\beta r_\beta^{(a)}=-q^\mu q_\mu r_\alpha^{(a)}\ , \tag{2.9}$$

which holds when $\omega=\omega^{(a)}(\mathbf{q})$. We obtain

$$\begin{aligned}\Gamma^A=\frac{1}{16\pi^2e^2}G_1^2\int d^3q(E_\mathbf{p}E_{\mathbf{p}-\mathbf{q}})^{-1}\\\times\{\mathbf{q}^2-[\omega(\mathbf{q})]^2\}^2\delta[\omega(\mathbf{q})-\mathbf{q}\cdot\mathbf{v}]\\\times\sum_{\text{polar}}|\bar{u}_f\gamma^\mu(1-\gamma_5)r_\mu^{(a)}u_i|^2\end{aligned} \tag{2.10}$$

and

$$\begin{aligned}\Gamma^B=\frac{1}{16\pi^2m_e^2}\mu_{10}^2e^210^{-20}\\\times\int d^3q(E_\mathbf{p}E_{\mathbf{p}-\mathbf{q}})^{-1}\delta[\omega(\mathbf{q})-\mathbf{q}\cdot\bar{\mathbf{v}}]\\\times\sum_{\text{polar}}|\bar{u}_f\sigma^{\alpha\beta}q_\alpha r_\beta^{(a)}u_i|^2\ ,\end{aligned} \tag{2.11}$$

where the sum is over the possible polarization state [type (a)] of the excitation (we assume that the incoming neutrinos are left handed).

While the above development, supplemented by the equations for the residue functions $r_\mu(\mathbf{q})$ given later by (3.3) and (4.4), is theoretically complete, a more formal approach can be followed that expresses the reaction rates in terms of integrals over the correlation functions of the electric and magnetic fields in the medium. Our results are recaptured when the correlation functions are replaced by single-excitation pole terms. Reference [3] gives a complete treatment, along the above-indicated lines, of the related problem of plasmon decay into $\nu\bar{\nu}$.

We write the general form for the polarization part which obeys the constraint of current conservation, $q_\mu \Pi^{\mu\nu}=0$, as [4]

$$\Pi_{ij}=\left[\delta_{ij}-\frac{\mathbf{q}_i\mathbf{q}_j}{|\mathbf{q}|^2}\right]B(\mathbf{q},\omega)+\frac{\mathbf{q}_i\mathbf{q}_j}{|\mathbf{q}|^2}C(\mathbf{q},\omega)\ ,$$

$$\Pi_{i0}=\frac{\mathbf{q}_i}{\omega}C(\mathbf{q},\omega),\quad \Pi_{00}=\frac{|\mathbf{q}|^2}{\omega^2}C(\mathbf{q},\omega)\ . \tag{2.12}$$

The above development, along with the identifications [4] $\Pi_L=(1-|\mathbf{q}|^2/\omega^2)C$, $\Pi_T=B$, should be equivalent to the general results presented in Ref. [3]. However, it differs in a number of important respects that are critical to the present work [5].

III. PRODUCTION OF PHOTONS IN A DIELECTRIC MEDIUM

In this case the function $B(\mathbf{q},\omega)$, which describes the transverse excitations, is of order ω^2 in the limit $\mathbf{q}=0$, $\omega\rightarrow 0$:

$$B(0,\omega)=C(0,\omega)=[1-\epsilon(\omega)]\omega^2\ , \tag{3.1}$$

where $\epsilon(\omega)$ is the dielectric constant. The dispersion relation for the transverse excitations is

$$\omega^2=\epsilon^{-1}(\omega)|\mathbf{q}|^2\ . \tag{3.2}$$

The residue functions $r_\mu^{(T)}$ are

$$r_\mu^{(T)}(\mathbf{q})=\left[\frac{\partial}{\partial\omega}[\omega^2\epsilon(\omega)]_{\omega=\omega(\mathbf{q})}\right]^{-1/2}\xi_\mu^{(T)}$$
$$\equiv r^T(\mathbf{q})\xi_\mu^{(T)}\ , \tag{3.3}$$

where $\xi_\mu^{(T)}$ are the transverse-photon-polarization vectors, normalized to unity. Putting (3.3) into the expressions for the production rate [(2.10) and (2.11), respectively] for the two excitation mechanisms, we perform the spin sums and angular integrations, obtaining

$$\Gamma^A=G_1^2(2e^2\pi)^{-1}\int_{\underline{q}}^{\bar{q}}dq\, q^5[\epsilon(q)-1]^3\epsilon^{-3}(q)[r^T(q)]^2 \tag{3.4}$$

and

$$\Gamma^B=(8\pi m_e^2)^{-1}10^{-20}\mu_{10}^2e^2$$
$$\times\int_{\underline{q}}^{\bar{q}}dq\, q^3[\epsilon(q)-1]^2\epsilon^{-2}(q)[r^T(q)]^2\ , \tag{3.5}$$

for the case of constant ϵ, using (3.3). The upper and lower limits bound the region of q for which $n(q)>1$. For the case of a constant (or slowly varying) dielectric constant, the result (18) agrees exactly with the result of the classical calculation of Ginzburg [6].

IV. PRODUCTION OF PLASMONS

In the case of propagation in a plasma, the functions $B(0,\omega)$ and $C(0,\omega)$ are nonvanishing for $\omega=0$, and we have $B(0,\omega)=C(0,\omega)$. At $\mathbf{q}=0$ both the longitudinal and transverse modes have $\omega^2=\omega_p^2$. We limit our consideration to the longitudinal case; only the longitudinal dispersion relation enters the region $\mathbf{q}>\omega$ in which the emission of the excitation into the medium from a source moving at the speed of light is allowed. To find the dispersion relation and the residue function, we solve

$$D_{\mu\nu}^{-1}r^\nu(\mathbf{q},\omega)=0\ , \tag{4.1}$$

obtaining

$$\omega^2=C(\mathbf{q},\omega)\ , \tag{4.2}$$

as the equation which determines the dispersion relation, $\omega=\omega^L(\mathbf{q})$, for the longitudinal excitations. The eigenvector, for the case of propagation in the 3 direction, is

$$r_1^L=r_2^L=0,\quad r_3^L=\omega f^L(\mathbf{q}),\quad r_0^L=|\mathbf{q}|f^L(\mathbf{q})\ . \tag{4.3}$$

The normalization function $f^L(\mathbf{q})$, which determines the momentum-dependent factor in the coupling strengths of the neutrino to the plasmon, can be found by differentiating $D_{\mu\nu}$ as given by (2.6) with respect to ω, substituting (2.7), and equating coefficients of the double-pole term:

$$-\frac{r_\mu^{(a)}r_\nu^{(a)}}{[\omega-\omega^a(\mathbf{q})]^2}$$
$$=\frac{r_\mu^{(a)}r_\nu^{(a)}r_\alpha^{(a)}\left[-2\omega g^{\alpha\beta}+\frac{\partial}{\partial\omega}\Pi^{\alpha\beta}\right]r_\beta^{(a)}}{[\omega-\omega^{(a)}(\mathbf{q})]^2}+\cdots\ , \tag{4.4}$$

where the omitted terms are less singular at $\omega=\omega^{(a)}(\mathbf{q})$. Using (2.12), we obtain

$$(f^L)^{-2}=\omega^{-1}(\omega^2-\mathbf{q}^2)^2[2-\omega^{-1}(\partial C/\partial\omega)]|_{\omega=\omega(\mathbf{q})}$$
$$=-(\omega^2-\mathbf{q}^2)^2\frac{\partial}{\partial\omega}\left[\frac{C}{\omega^2}\right]\Bigg|_{\omega=\omega(\mathbf{q})}\ , \tag{4.5}$$

where we have used (4.2) to obtain the second form. All of the above is applicable both to nonrelativistic plasmas in metals and to relativistic plasmas in astrophysical situations. In the estimates that follow, we shall use standard results for a one-component, nonrelativistic electron plasma, at zero temperature and for small $\mathbf{q}$ [4,7]:

$$C(\mathbf{q},\omega)=\omega_p^2\left[1+\frac{3}{5}\left[\frac{k_f|\mathbf{q}|}{m_e\omega}\right]^2+\cdots\right]\ , \tag{4.6}$$

where

$$\omega_p=\left[\frac{n_ee^2}{m_e}\right]^{1/2},\quad k_f=(3\pi^2n_e)^{1/3}\ , \tag{4.7}$$

and n_3 is the electron-number density. Using (4.6), we can calculate the coupling function $r^L(q)$ from (4.5) and then the production rates from (2.10) and (2.11). In this case the large momentum cutoff is most reasonably taken to be that momentum for which the phase velocity equals the Fermi velocity v_F, the point beyond which strong damping sets in and the plasmon loses its identity. In principle, we need to keep more terms in the expansion (4.6) in this region, since the expansion parameter is of order unity at the cutoff point. However, for our order-of-magnitude estimates, it should suffice to keep only the first two terms in (4.6) in evaluating the upper cutoff momentum, giving $\mathbf{q}\approx 1.3(v_F)^{-1}\omega_p$. The energy corresponding to this momentum is given by $\bar{\omega}\approx 1.3\omega_p$. For the coupling function we obtain

$$f^L(\mathbf{q})=2^{-1/2}|\mathbf{q}|^{-2}\omega_p^{1/2}\left[1-0.48\frac{q^2}{\bar{q}^2}+O(q^4)\right], \tag{4.8}$$

from (4.5) and (4.6), making an expansion in powers of $(q/\bar{q})^2$ and setting $q^2-\omega^2\approx q^2$, the last approximation justified by the fact that almost all of the phase space for production is in a region in which q is many times its threshold value, defined by the solution to $|\mathbf{q}|=\omega(\mathbf{q})$. We obtain the respective expressions for the transition rate:

$$\Gamma^A\approx 1.4\times 10^{-2}G_1^2e^{-2}\omega_p\bar{q}^4 \tag{4.9}$$

and

$$\Gamma^B\approx 7.56\times 10^{-23}\mu_{10}^2 m_e^{-2}e^2\omega_p\bar{q}^2 . \tag{4.10}$$

V. DISCUSSION

In materials in which the valence electrons are tightly bound, the region in which the index of refraction is greater than unity can extend well into the ultraviolet [8]. To estimate the order of magnitude of Cherenkov rates in such materials, we consider a material with an index of refraction of 2 up to a cutoff of 10 eV, followed by a sharp break to $n<1$ for $\omega\geq 10$ eV. In this case the cutoff $\mathbf{q}$ is at 20 eV/c in (3.4) and (3.5). We express the answers in terms of $(L)^{-1}=$ transitions per cm, which we define as the "rate":

$$\begin{aligned} L_A^{-1}&\approx 10^{-33}\ \mathrm{cm}^{-1} ,\\ L_B^{-1}&\approx 10^{-25}\mu_{10}^2\ \mathrm{cm}^{-1} . \end{aligned} \tag{5.1}$$

The results are probably too small to be interesting. For case (B), where the coupling is through a neutrino magnetic moment, the Cherenkov rate is about one-ten-thousandth the rate of electron-neutrino scattering arising from photon exchange in the presence of the same magnetic-moment interaction [9], for an electron density of 5×10^{23} cm^{-3}. Both of these rates are essentially independent of the energy of the neutrino, as long as the neutrino energy is much larger than the scale of atomic energy levels. At still lower neutrino energies, the direct magnetic term will die off rapidly, but the Cherenkov scattering will be constant down to about 20 eV. It does not seem worth pursuing this quantitatively at this time, since there is little likelihood of space being filled with, e.g., 20-eV *relativistic,* exotic neutrinos, with magnetic moments of $10^{-7}\mu_e$, up to the limits allowed by cosmology. Indeed, this is one of the few possibilities which is currently not being considered in the game of guessing as to what may be lurking around us, undetected. Note, nonetheless, that the direct V-A interactions, both with nuclei and electrons, would produce negligible scattering at energies below 1 keV, even on the scales we are considering.

The photon-production rate [Eq. (5.1)] coming from the ordinary left-handed current interactions, in the absence of a magnetic moment, is far too small to be of further interest.

Turning to the plasmon rates (4.9) and (4.10), we make an estimate on the basis of an electron density of $\rho=5\times 10^{23}$ cm^{-3}, whence $\omega_p=3.4\times 10^{-2}$ keV, $k_F=5.7$ keV/c, and $\bar{q}=4$ keV/c. We obtain

$$\begin{aligned} L_A^{-1}&\approx 1.3\times 10^{-24}\ \mathrm{cm}^{-1} ,\\ L_B^{-1}&\approx 0.75\times 10^{-21}\mu_{10}^2\ \mathrm{cm}^{-1} , \end{aligned} \tag{5.2}$$

for all incident neutrino energies greater than 2 keV. Note that because of the smallness of ω/q on the dispersion curve, we can use the simplified kinematical condition (1.1) for any case in which the maximum-allowed momentum transfer to the neutrino $2E_\nu$ is greater than the cutoff momentum; the reaction rate will be independent of energy throughout this region. At values of the ν energy below $\bar{q}$, the rate will be less. We compare (4.9) with the upper bound on the rate of energy deposition $<Q$ through the vector part of the neutrino-electron interaction established by Kirzhnits, Losyakov, and Chechin and valid in any cold medium [Eq. (8.6) of Ref. [10], for $E_\nu\gg\omega_p$]:

$$Q\leq(20/3\pi)G_1^2e^{-2}\omega_p^2E_\nu^4 .$$

Multiplying (4.10) by the plasmon energy at cutoff $1.3\omega_p$, we find that for $E_\nu=2$ keV the bound is approximately one-half saturated.

The results [Eq. (5.2)] are worthy of some consideration. While the result for case (A) (vanishing or small neutrino magnetic moment) is apparently still about two orders of magnitude too small for detection, it does have some interesting properties: At energies below about 30 keV, it is larger than the rate given by the neutrino cross section on free electrons times the density of 5×10^{23} cm^{-3}. Below about 10 keV the competing cross section from noncollective effects (e.g., ionization) become further damped by the effects of binding. Finally, the emission is dominantly at nearly 90° to the neutrino beam, as a result of the condition $\omega/q\ll 1$ and the kinematical relation (1.1).

For case (B), in which the process is dominated by the neutrino magnetic moment, the reaction rate (5.2) can be compared with a rate estimated on the basis of free electrons, with the e-ν cross section from single-photon exchange, the photon coupled to the neutrino via the magnetic-moment interaction. The cross section is given in the laboratory system by [9]

$$\begin{aligned}\sigma=(16\pi)^{-1}e^4\mu^2\{&\ln[E_\nu(\underline{E}-m_e)^{-1}(1+m_e/2E_\nu)^{-1}]\\ &\times(1+m_e/2E_\nu)^{-1}+(\underline{E}-m_e)/E_\nu\} .\end{aligned} \tag{5.3}$$

Here $\underline{E}$ is a lower cutoff in the kinetic energy of the final electron. Again, taking the electron density as 5×10^{23} cm^{-3}, we find a rate L^{-1}, which is a constant in energy, 1.2×10^{-21} cm^{-1}, times a logarithmic factor that is of order unity. Therefore the noncollective effects of the magnetic-moment coupling can be expected to be about an order of magnitude larger than the plasmon production, for the case in which neutrino energies are high enough that the energy transfer to the electron is larger than typical binding energies in the atom or solid. When the neutrino energy is reduced to the range in which the main range of energy transfers begins to fall *below* the binding-energy range, for neutrino energies below 10 keV or so, the plasmon rate will remain constant, as long as E_ν is greater than $\underline{E}$, but the noncollective background should fall off rapidly. At, say, $E_\nu=2$ keV, where the maximum energy transfer to a *free* electron would be ≈10 eV, we expect the attenuation of the background due to binding to be very large. Unfortunately, there is no extensive body of literature known to us on the passage of a neutral particle with a magnetic moment, and no nuclear interactions, through matter; therefore, at the moment, our conclusions are qualitative at best.

The result (5.2), for case (B), is in a region of possible interest for values of the neutrino magnetic at or near the maximum allowed by the constraints imposed by stellar evolution. The flux, at earth, of plasma neutrinos from the Sun is estimated as $2\times10^{13}\mu_{10}^2\ \text{cm}^{-2}\,\text{s}^{-1}$, with a spectrum peaked at about 2 keV. Thus the reaction rate in a detector would be of the order of $10^{-8}\mu_{10}^4\ \text{cm}^{-3}\,\text{s}^{-1}$. Since the maximum-allowed magnetic moment of a light neutrino (mass less than 10 keV) is, very conservatively, limited by stellar-evolution considerations to the range $\mu_{10}<0.1$ [11], the observation again seems to be somewhat out of range, according to the above estimate.

Another possibility is that of detecting some species of nonelectron neutrino of mass more than 10 keV and magnetic moment μ_{10} *greater* than unity (the mass in this case providing the defense against too large a stellar cooling rate through the plasmon mechanism). For example, the magnetic moment of the muon neutrino is currently bounded, from electron-scattering data, at about $\mu_{10}<10$ [12]. Since accelerators and reactors would provide the source of such neutrinos, their energies would now be high (on the scales that we have been considering) energies of 0.1 MeV at minimum. In this case, for example, in the case of a neutrino energy of 10 MeV the logarithm in expression (5.3) would be of the order of magnitude of 10, if the low-energy electron cutoff is chosen as, say, 10 eV. Thus the direct magnetic-moment reactions would dominate the magnetic production of a plasmon by about a factor of 10. However, a distinctive signal, one that provides strong discrimination against backgrounds, might well override a factor of 10 in providing information.

We note that the plasmon mechanism is inapplicable on kinematical grounds [Eq. (1.1)] to the detection of neutrinos of velocity less than v^F ($\approx0.01c$ under the conditions of our example), as in scenarios of neutrinos as dark matter.

Of course, up until now, we have avoided the truly difficult question: What do we have when we have a plasmon? It is an excitation that is never detected directly, in, e.g., experiments on electron scattering from solids. Here we offer a suggestion of an efficient transformation mechanism of plasmon to photon, with the caution that it requires much more study. Data indicate significant anisotropies in the plasmon-dispersion relations for certain crystalline substances, with up to a 10% dependence on direction of the value of ω for values of q in the region of our upper limit (of course, the long-wavelength limit does not detect crystal structure) [13]. Let us parametrize the inverse propagator (or dielectric tensor) in a minimal way that achieves the splitting, by adding, say, off-diagonal terms in the form of 1,3 components proportional to q^2 to the structure given in (2.12), together with the requisite 1,4 terms for consistency with current conservation. We can then determine the new polarization vectors for the new (predominantly) longitudinal and transverse modes in the presence of the mixing term. Next, we consider the longitudinal wave normally incident on the boundary of the electron gas. In the absence of mixing, we obtain the expected result: The continuity of the electric potential across the boundary dictates the perfect reflection of the wave. With the mixing included, however, there is a calculable additional effect at the boundary in the form of the generation both of a reflected transverse wave in the medium and a photon proceeding into free space. The efficiencies of transmutation can be greater than 50% in the estimate sketched above.

For the purpose of further speculation, then, let us assume that materials can be found that accomplish the efficient transmutation of (predominantly) longitudinal waves to photons at an interface. The mechanism would only be effective, however, when the plasmon is produced within a plasmon-absorption length of the surface, so that films (or filaments) would be required, with a thickness in-between the plasmon wavelengths (10^{-8} cm or so for our dominant modes) and the absorption length (perhaps 10^{-5} cm). Recalling that the plasmons, and the converted photons (for the case of normal incidence on the interface), are emitted at nearly 90° to the neutrino direction, we might envision a bundle of filaments pointed in the direction of the neutrino source as a useful geometry.

We have done estimates as to whether pure transition radiation, in the form of transverse plasmons emitted directly into the medium, arising from the effect of the plasma on the fields attached to the entering neutrino, would provide a more effective way of producing photons than the conversion process conjectured above. It is amusing to do the calculation, particularly in the case of the radiation generated when the coupling is through the electron-neutrino interaction (a novel result, we believe). But the results are somewhat too small to be of interest.

It is also worth noting that the plasmon-production mechanism, which our estimates show could be the dominant mode of neutrino-energy deposition for the case of magnetic-moment interactions at low energies, coexists with the cryogenic detection possibilities that have been discussed in the literature [14]. In this case one would look for materials choices and geometrical configurations that would make the rate of plasmon *generation* depend on the orientation of the detector, since one detects only the heat produced.

There has been no mention, in this paper, of the neutrino's interactions with nuclei. The nuclear cross sections fall like E_ν^2, for low-energy neutrinos, as do the neutrino-electron cross sections, and the nuclear cross-section gains from the coherence effect. However, a single nuclear scattering deposits a very small amount of energy in the medium.

Again, we emphasize that the results of these speculations are not positive enough to point the way for a present effort to design new detectors. However, the demonstrated possibility that the effects of the medium can transmute much of the reaction amplitude for low-energy neutrinos into a channel with a distinctive signal, so that backgrounds could potentially be minimized, and the estimated cross sections themselves, which are higher than one might have supposed, both argue for fuller and more authoritative treatment of the plasmon possibility and for a fuller survey of the role of other energetic electronic collective excitations in materials.

ACKNOWLEDGMENTS

This work was supported by NSF Grant No. PHY 86-14185.

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[5] In particular, Ref. [3] finds *no* contribution of a magnetic-moment coupling to the production of longitudinal plasmons, a contribution that provides the most promising application of the present work. The discrepancy appears to stem from the construction given in Ref. [3] of the photon Green's function in the medium in terms of the longitudinal and transverse dielectric constants [Eq. (18)], a construction that we believe is incorrect. The same construction also leads to results different from ours on the coupling of the longitudinal plasmon through the ordinary vector interaction. On a different issue, note that the erratum cited in Ref. [2] withdraws the result that plasmon effects greatly enhance the electron scattering, for neutrinos of a few keV energy. The authors also did not consider the energy deposition from plasmons with momentum below the strong-absorption threshold (since it does not contribute to their effective neutrino cross section on single electrons; that is, the plasmons do not have a single-particle-hole decay mode in this region).

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