ЖЭТФ, 2004, том 126, вып. 3 (9), стр. 533–541

EXCITATION OF STRONG WAKEFIELDS BY INTENSE NEUTRINO BURSTS IN A MAGNETIZED ELECTRON–POSITRON PLASMA

A. Serbeto [a], ***L. A. Rios*** [a], ***J. T. Mendonça*** [b], ***P. K. Shukla*** [c], ***R. Bingham*** [d]

[a] *Instituto de Física, Universidade Federal Fluminense*
24210-340, Rio de Janeiro, Brazil

[b] *GoLP/Instituto Superior Técnico, Universidade Técnica de Lisboa*
1049-001, Lisbon, Portugal

[c] *Institut für Theoretische Physik IV, Ruhr-Universität Bochum*
D-44780, Bochum, Germany

[d] *Rutherford Appleton Laboratory, Chilton, Didcot*
Oxon OX11 0QX, England

Submitted 12 April 2004

A classical fluid description is used to investigate nonlinear interactions between an electron-type neutrino burst and a collisionless magnetized electron–positron plasma. It is found that the symmetry between the electron and positron dynamics is broken due to the presence of intense neutrino bursts. The latter can excite strong upper-hybrid wakefields, which can produce unlimited acceleration of pairs across the external magnetic field direction via a surfatron mechanism. Implications of our results to the production of high-energy electrons and positrons in astrophysical environments are discussed.

PACS: 13.15.+g, 52.25.Dg, 97.10.Cv

1. INTRODUCTION

It is well established that processes involving electron–positron plasmas are of significant importance in a variety of astrophysical phenomena [1–4] that presumably took place in the early Universe, during the period $10^{-4} < t < 1$ s after the Big Bang, and in ultra-relativistic pair plasma jets. Electron–positron plasmas are also found in active galactic nuclei [5, 6] and in the pulsar magnetosphere [7, 8]. Another important phenomenon is the γ-ray burst [9, 10], where many of its characteristics are explained through the relativistic expansion of the electron–positron plasma («the fireball model») [10]. It is well known that such a plasma has a very peculiar nature compared with the traditional electron–ion plasma, because electrons and positrons have the same mass and represent a symmetry due to their opposite electric charges. These unique properties are responsible for many linear and nonlinear wave phenomena that are different from those occurring in the electron–ion plasma. This difference is due to the electron–ion mass ratio, which gives rise to different timescales associated with the electron and ion dynamics in the plasma. A systematic study (see, e.g., Ref. [11] and references therein) of nonlinear interactions between intense electromagnetic waves and relativistic electron–positron plasmas has been carried out, especially in the pulsar pair plasma environment for understanding the origin of the pulsar radio emission.

Another very important material elements in the astrophysical settings are the neutrinos. They are produced by the core of stars and in very high-explosive astrophysical situations such as those in supernova explosions and in ultrarelativistic pair plasma jets producing γ-ray bursts. Recently, many authors [12–15] have proposed that the shock expansion mechanism in the supernova could be due to the energy–momentum transfer from the neutrino bursts to the magnetized plasma cloud that surrounds the core of stars. Here, we are interested in nonlinear interactions of intense neutrino bursts with a relativistic magnetized pair plasma. It is well known [16–19] that neutrinos propagating through the plasma can acquire an effective (induced) electric charge due to charged and neutral currents associated with the weak nuclear force causing exchange of $W^{\pm}$ and Z^0 bosons. There also appears a nonlinear coupling between the neutrinos and the plasma through the weak Fermi nuclear interaction and the ponderomotive force of the neutrino beam [20–25].

In this paper, we study the generation of large-amplitude upper-hybrid waves in a magnetized electron–positron plasma interacting with an electron-type neutrino burst. The neutrino dynamics can be considered semiclassical by assuming that the neutrino de Broglie wavelength is much shorter than the typical scalelength of the perturbation in the effective neutrino weak interaction potential. We ignore all quantum-mechanical effects (e.g., the neutrino magnetic moment) caused by external magnetic fields because

$$\frac{B_0}{B_{QM}} \equiv \frac{\hbar/\omega_{ce}}{m_e c^2} \ll 1,$$

where $\hbar$ is the Planck constant divided by 2π,

$$\omega_{ce} = \frac{eB_0}{m_e c}$$

is the electron gyrofrequency, e is the magnitude of the electron charge, m_e is the electron mass, c is the speed of light in the vacuum, and $B_{QM} \approx 4 \cdot 10^{13}$ Gauss is the Landau–Schwinger critical field. Neutral current forward scattering of neutrinos off neutrinos in the background gives a contribution proportional to the density matrix, which has no effect on the flavor evolution [26].

This paper is organized as follows. In Sec. 2, we present the governing equations for neutrinos and electrostatic waves. Section 3 contains an analytic description for the upper-hybrid wakefield in the presence of neutrino fluxes. Numerical results for the wakefield are presented in Sec. 4. Finally, Sec. 5 highlights our results and contains possible applications of our work to acceleration of pairs by large-amplitude upper-hybrid waves in astrophysical environments.

2. GOVERNING EQUATIONS

We consider an electron-type neutrino burst in a magnetized electron–positron plasma. The external magnetic field $B_0\hat{\mathbf{z}}$ is along the z direction. The dynamics of an ensemble of the neutrinos can be described by the equations [22]

$$\frac{\partial N_\nu}{\partial t} + \nabla \mathbf{J}_\nu = 0 \qquad (1a)$$

and

$$\frac{\partial \mathbf{p}_\nu}{\partial t} + (\mathbf{v}_\nu \nabla)\mathbf{p}_\nu = -\frac{1}{N_\nu}\nabla P_\nu + \sum_\sigma G_{\sigma\nu}\left(\mathbf{E}_\sigma + \frac{\mathbf{v}_\nu}{c} \times \mathbf{B}_\sigma\right), \qquad (1b)$$

which couple the neutrino density N_ν and the neutrino momentum $\mathbf{p}_\nu$. Here,

$$G_{\sigma\nu} = \sqrt{2} G_F [\delta_{\sigma e}\delta_{\nu\nu_e} + (I_\sigma - 2Q_\sigma \sin^2\theta_w)]$$

is the «bare» charge, with $G_{\sigma\nu} = -G_{\bar{\sigma}\nu}$, which allows the neutrinos to couple to the plasma fluid. Furthermore, σ denotes the electron (e^-) and positron (e^+) species of the plasma, $G_F (\approx 9 \cdot 10^{-30}$ eV $\cdot$ cm$^{-3})$ is the Fermi weak-interaction coupling constant, θ_W is the Weinberg mixing angle ($\sin^2\theta_W \approx 0.23$), I_σ is the weak isotopic spin of the particle of the species σ (equal to $-1/2$ and $1/2$ for the electrons and positrons, respectively), and $Q_\sigma = q_\sigma/e$ is the particle normalized electric charge. It should be noted that the first term in the neutrino «bare» charge is due to charged currents, which is valid only for electrons (positrons) and the electron-type neutrinos, and other terms come from the neutral currents and are valid for all particle species.

In Eq. (1b), $P_\nu = N_\nu T_\nu$ is the neutrino kinetic pressure and the second term in Eq. (1b) is the weak force, $\mathbf{F}_\nu$, on a single neutrino due to the plasma. Furthermore,

$$\mathbf{E}_\sigma = -\nabla N_\sigma - \frac{1}{c^2}\frac{\partial \mathbf{J}_\sigma}{\partial t},$$
$$\mathbf{B}_\sigma = c^{-1}\nabla \times \mathbf{J}_\sigma$$

are the effective electric and magnetic fields [22], respectively, and

$$\mathbf{J}_\nu = \mathbf{v}_\nu N_\nu, \quad \mathbf{J}_\sigma = \mathbf{v}_\sigma N_\sigma$$

are the neutrino and σ-species currents, respectively. The linear momentum of the neutrino is

$$\mathbf{p}_\nu = \frac{\mathbf{v}_\nu}{c^2} E_\nu,$$

with E_ν being the neutrino total energy. The term $\partial \mathbf{J}_\sigma / \partial t$ is the neutrino-plasma (treated covariantly) analogue of the electromagnetic-plasma energy transfer, as described in Ref. [22]. Furthermore, strong dc magnetic fields can create magnetic-field-aligned motion of the neutrinos and plasma particles.

The plasma particle dynamics is governed by the continuity and momentum equations, which are, respectively,

$$\frac{\partial N_\sigma}{\partial t} + \nabla \mathbf{J}_\sigma = 0, \qquad (2a)$$

and

$$\frac{\partial \mathbf{P}_\sigma}{\partial t} + (\mathbf{v}_\sigma \nabla)\mathbf{P}_\sigma = q_\sigma \mathbf{E} + q_\sigma \frac{\mathbf{v}_\sigma}{c} \times \mathbf{B} + \\ + \sum_\nu G_{\sigma\nu} \left(\mathbf{E}_\nu + \frac{\mathbf{v}_\sigma}{c} \times \mathbf{B}_\nu \right), \quad (2b)$$

where

$$\mathbf{P}_\sigma = \gamma_\sigma m_\sigma \mathbf{v}_\sigma,$$

is the momentum of the particle species σ (electrons and positrons),

$$\gamma_\sigma = \frac{1}{\sqrt{1 - v_\sigma^2/c^2}}$$

is the gamma factor, and $\mathbf{B} = B_0 \hat{\mathbf{z}}$ is the external uniform magnetic field in the z direction. The right-hand side in Eq. (2b) is the total force acting on the plasma due to all types of the neutrinos, and

$$\mathbf{E}_\nu = -\nabla N_\nu - c^{-2} \frac{\partial \mathbf{J}_\nu}{\partial t},$$

$$\mathbf{B}_\nu = c^{-1} \nabla \times \mathbf{J}_\nu$$

are the «weak-electromagnetic» fields. Furthermore, N_σ is the number density of the species σ. Because we focus on the wakefield generation on timescales that are either comparable with or shorter than the electron plasma period, collisions between pairs do not play an essential role, for instance, in the supernova and pair plasma jets [24, 25], where the plasma number density and the pair temperature are 10^{30} cm^{-3} and 10^5–10^6 eV, respectively.

To simplify our model, we consider only a cold electron-type neutrino streaming along the x direction with the velocity v_ν close to c, interacting nonlinearly with a collisionless cold magnetized electron–positron plasma. Hence, antineutrinos are neglected. It is well known that the interaction of the electron neutrinos with the plasma does not change their local energy and density significantly. For instance, in type-II supernova explosions, only 1 % of the neutrino energy [14] is supposed to be transferred to the plasma that surrounds the core of the star. Hence, without the loss of generality, we can assume that the electron-type neutrino flux only transfers a very small part of its energy E_ν to the plasma and keeps its density N_ν nearly constant. Accordingly, the electron-type neutrino fluid dynamical equations can be rewritten as

$$\frac{\partial E_\nu}{\partial t} + c\frac{\partial E_\nu}{\partial x} \approx \\ \approx -\sqrt{2} G_F c \left(\frac{\partial}{\partial x}(N_{\bar{e}} - N_e) + \frac{1}{c^2}\frac{\partial}{\partial t}(J_{\bar{e}} - J_e) \right), \quad (3a)$$

and

$$\frac{\partial N_\nu}{\partial t} + c\frac{\partial N_\nu}{\partial x} \approx 0. \quad (3b)$$

With the longitudinal plasma waves propagating across the external magnetic field direction, with an associated electric field $\mathbf{E} = E\hat{\mathbf{x}}$ and the wavenumber $\mathbf{k} = k\hat{\mathbf{x}}$, the electron–positron plasma fluid equations are

$$\frac{\partial N_e}{\partial t} + \frac{\partial J_e}{\partial x} = 0, \quad (4a)$$

$$\frac{\partial P_{ex}}{\partial t} + \frac{c P_{ex}}{\sqrt{1+P_e^2}} \frac{\partial P_{ex}}{\partial x} = -\frac{eE}{m_e c} - \frac{\omega_c P_{ey}}{\sqrt{1+P_e^2}} - \\ - \sqrt{2} G_F \left(\frac{\partial N_\nu}{\partial x} + \frac{1}{c^2}\frac{\partial J_\nu}{\partial t} \right), \quad (4b)$$

$$\frac{\partial P_{ey}}{\partial t} + \frac{c P_{ey}}{\sqrt{1+P_e^2}} \frac{\partial P_{ey}}{\partial x} = \frac{\omega_c P_{ex}}{\sqrt{1+P_e^2}}, \quad (4c)$$

and

$$\frac{\partial N_{\bar{e}}}{\partial t} + \frac{\partial J_{\bar{e}}}{\partial x} = 0, \quad (5a)$$

$$\frac{\partial P_{\bar{e}x}}{\partial t} + \frac{c P_{\bar{e}x}}{\sqrt{1+P_{\bar{e}}^2}} \frac{\partial P_{\bar{e}x}}{\partial x} = \frac{eE}{m_e c} + \frac{\omega_c P_{\bar{e}x}}{\sqrt{1+P_{\bar{e}}^2}} + \\ + \sqrt{2} G_F \left(\frac{\partial N_\nu}{\partial x} + \frac{1}{c^2}\frac{\partial J_\nu}{\partial t} \right), \quad (5b)$$

$$\frac{\partial P_{\bar{e}y}}{\partial t} + \frac{c P_{\bar{e}y}}{\sqrt{1+P_e^2}} \frac{\partial P_{\bar{e}y}}{\partial x} = -\frac{\omega_c P_{\bar{e}x}}{\sqrt{1+P_{\bar{e}}^2}}, \quad (5c)$$

for the electron and positron plasma species, with P_{ex}, $P_{\bar{e}x}$, P_{ey}, and $P_{\bar{e}y}$ being the x and y components of the electron and positron momenta, respectively.

Equations (3), (4), and (5) form a set for studying the generation of large-amplitude plasma waves. To eliminate the electron and positron fluid currents in Eq. (3a), we supplement our system of equations with

$$\mathbf{J}_{\bar{e}} - \mathbf{J}_e = \frac{1}{4\pi e}\frac{\partial \mathbf{E}}{\partial t}, \quad (6)$$

because we consider the generation of longitudinal (electrostatic) waves.

3. UPPER-HYBRID WAKEFIELD

It is convenient to introduce a new independent variable $\xi = (x - v_\phi t)$, where v_ϕ is the plasma wave phase speed. Hence, Eq. (3a) can be rewritten as

$$\frac{dE}{d\xi} = 4\pi e \frac{c^2}{v_\phi^2} \times \\ \times \left((N_{\bar{e}} - N_e) + \frac{1 - \beta_\phi}{\sqrt{2} G_F}(E_0 - E_\nu(\xi)) \right), \quad (7)$$

where E_0 is the neutrino initial energy. Using the definition $E = -d\Phi/d\xi$, where Φ is the electric potential associated with the wakefield, we obtain from Eq. (7) that

$$\frac{d^2\Psi}{d\xi^2} = -\frac{\omega_p^2}{v_\phi^2}\left(\frac{N_{\bar{e}}}{N_0} - \frac{N_e}{N_0} + S_\nu\right), \tag{8}$$

where

$$\Psi = \frac{e\Phi}{m_e c^2}$$

is the normalized plasma potential,

$$\beta_\phi = \frac{v_\phi}{c}$$

is the normalized phase speed,

$$\omega_p^2 = \frac{4\pi e^2 N_0}{m_e}$$

is the squared electron plasma frequency, and N_0 is the equilibrium electron (positron) number density. Furthermore,

$$S_\nu = \frac{E_0(1-\beta_\phi)}{\sqrt{2}G_F N_0}\,\frac{\Delta E_\nu}{E_0}$$

represents the neutrino-driven term, with

$$\frac{\Delta E_\nu}{E_0} = \frac{\Delta\omega_\nu}{\omega_\nu}$$

being considered the amount of the neutrino energy transferred to the plasma. Here, $\Delta\omega_\nu$ is the spectral width of the neutrino spectrum. Assuming that the magnitude

$$\frac{\Delta E_\nu}{E_0} = \frac{\Delta\omega_\nu}{\omega_\nu} \ll 1$$

(actually, from the observational fact in the supernova SN1987A, the «visible» energy of the supernova is a very small part of the neutrino energy [14]), we can consider the neutrino flux as an external action into the plasma such that the amount of the neutrino energy deposited in the plasma can be taken as a constant input in Eq. (8). We note that the latter is the Poisson equation written in a moving frame, where the total charge density includes the neutrino effective charge density represented by the term S_ν. Transforming Eqs. (4) and (5) and substituting the electron (positron) density response in Eq. (8), we obtain

$$\frac{d^2\Psi}{d\chi^2} = \frac{\Gamma_e\sqrt{1+P_e^2}}{P_{ex} - \beta_\phi\sqrt{1+P_e^2}} - \frac{\Gamma_{\bar{e}}\sqrt{1+P_{\bar{e}}^2}}{P_{\bar{e}x} - \beta_\phi\sqrt{1+P_{\bar{e}}^2}} - S_\nu, \tag{9}$$

$$\frac{dP_{ex}}{d\chi} = \frac{\sqrt{1+P_e^2}}{P_{ex} - \beta_\phi\sqrt{1+P_e^2}} \times \left(\frac{d\Psi}{d\chi} - \beta_\phi\Omega_c\frac{P_{ey}}{\sqrt{1+P_e^2}}\right), \tag{10}$$

$$\frac{dP_{ey}}{d\chi} = \beta_\phi\Omega_c\frac{P_{ex}}{(P_{ex} - \beta_\phi\sqrt{1+P_e^2})}, \tag{11}$$

$$\frac{dP_{\bar{e}x}}{d\chi} = \frac{\sqrt{1+P_{\bar{e}}^2}}{P_{\bar{e}x} - \beta_\phi\sqrt{1+P_{\bar{e}}^2}} \times \left(-\frac{d\Psi}{d\chi} + \beta_\phi\Omega_c\frac{P_{\bar{e}y}}{\sqrt{1+P_{\bar{e}}^2}}\right), \tag{12}$$

and

$$\frac{dP_{\bar{e}y}}{d\chi} = -\beta_\phi\Omega_c\frac{P_{\bar{e}x}}{(P_{\bar{e}x} - \beta_\phi\sqrt{1+P_e^2})}, \tag{13}$$

where

$$P_{e,\bar{e}} = p_{e,\bar{e}}/m_e c$$

is the normalized electron (positron) momentum,

$$\chi = \omega_p/v_\phi(x - v_\phi t)$$

is the normalized distance (phase), and

$$\Omega_c = \omega_c/\omega_p$$

is the normalized electron (positron) gyrofrequency. Here,

$$\Gamma_{e,\bar{e}} = \frac{P_0 - \beta_\phi\sqrt{1+P_0^2}}{\sqrt{1+P_0^2}}$$

is a constant that depends on the initial value of the electron and positron momenta P_0. We note that the coupled equations (9)–(13) depend on the sign of the linear momentum of the plasma, i.e., the plasma can move either in the positive or in the negative χ-direction. It should be stressed that these equations describe the excitation of nonlinear relativistic upper-hybrid waves by neutrino beams in a magnetized plasma.

The nonrelativistic linear dynamics occurs for $P_{e,\bar{e}} \ll 1$, $\Psi \ll 1$, and $S_\nu \ll 1$. Subsequently, Eqs. (9)–(13) yield

$$\frac{d^2\Psi}{d\chi^2} = P_{ex} - P_{\bar{e}x} - S_\nu, \tag{14}$$

$$\frac{dP_{ey}}{d\chi} = -\Omega_c P_{ex}, \tag{15}$$

$$\frac{dP_{\bar{e}y}}{d\chi} = \Omega_c P_{\bar{e}x}, \tag{16}$$

$$\frac{d^2P_{ex}}{d\chi^2} + \Omega_{uh}^2 P_{ex} = P_{\bar{e}x} + S_\nu, \tag{17}$$

$$\frac{d^2 P_{\bar{e}x}}{d\chi^2} + \Omega_{uh}^2 P_{\bar{e}x} = P_{ex} - S_\nu, \quad (18)$$

where

$$\Omega_{uh}^2 = \omega^2/\omega_p^2 = 1 + \Omega_c^2$$

is the normalized squared upper-hybrid wave frequency. Equations (13)–(16) have been used to derive the coupled equations (17) and (18) for the x-components of the pair-plasma momentum. These linear equations show that the pair-plasma dynamics is directly forced by the presence of the neutrino fluxes, which can generate strong electrostatic waves with an oscillating frequency ω, propagating perpendicularly to the external magnetic field direction.

Solving Eqs. (17) and (18) and substituting the results in Eq. (14), we obtain the following normalized electric field ($E = -d\Psi/d\chi$) associated with the upper-hybrid wakefield:

$$E(x,t) = \frac{2\omega_p^3}{\omega^3} S_\nu \left(\frac{\Omega_c^2}{2}(kx - \omega t) + \sin(kx - \omega t) \right). \quad (19)$$

Here, the driven term S_ν is assumed constant, according to our initial assumption, and the initial conditions $P_{e\bar{e}} = 0$, $\Psi = 0$, and $E = 0$, for $\chi = 0$, are imposed. As we can see, this field has a large amplitude, clearly dependent on the amount of neutrino energy deposited into the plasma, and is independent of the neutrino density, i.e., the plasma dynamics is triggered by the neutrino energy deposit into the plasma cloud. This strong electric field can accelerate the plasma particles in the transverse direction, similarly to the surfatron acceleration mechanism [28], leading to the formation of a transverse pair-plasma jet.

4. NUMERICAL RESULTS

Figures 1a and 1b show the numerical solutions of Eqs. (9)–(13) for the neutrino-driven term $S_\nu = 3\cdot 10^{-3}$, $\Omega_c = 0.1$, and $\gamma_\phi = 350$. In Fig. 1a, we display the normalized electric field associated with the excited upper-hybrid waves. This field reaches a maximum whose value is close to the normalized gyrofrequency, Ω_c. This occurs because the plasma particle momentum in the y direction is much larger than the momentum in the x direction due to the weakness of the neutrino term S_ν. We can also see that the nonlinear regime only starts for large phase numbers ($\chi > 750$). During the linear regime, the dynamics of the electron and positron particles of the plasma is symmetric, as expected. We note that this linear regime corresponds to the analytic solution given by Eq. (19). However, for long neutrino–plasma interactions, this symmetry is broken owing to the presence of the neutrino flux, see Fig. 1b.

Figure 2a shows the normalized wake electric field for $\Omega_c = 0.1$, as in Fig. 1a but with a more intense neutrino-driven term, $S_\nu = 20$. We observe that as the neutrino-driven term increases, the x component of the particle momentum becomes much larger than the y component, which leads the electric field to a saturation level much larger than the value shown in Fig. 1a. We also note that in this case, the electric field presents a high intensity at the beginning of the neutrino–plasma interaction, and after a while starts to decay and to transfer its energy to the plasma particles. After that, the field reaches a saturation value $E \approx 0.5$, which is maintained during the interaction. Therefore, the plasma particles feel a constant electric force, which leads to an unlimited transverse acceleration, as we can see from Fig. 2b. It should also be pointed out that in this case, the asymmetry between the electron and positron dynamics is very intense, due to the high value of the induced negative charge that the neutrinos acquire during the interaction with the pair plasma. This process leads to a charge separation, which in turn creates the finite wakefields.

Until now, we have assumed that the gyrofrequency ω_c is smaller than the plasma frequency ω_p, which can occur in many astrophysical scenarios. Because our basic equations are normalized, we can also assume some astrophysical scenarios in which $\omega_c \gg \omega_p$, for example, a neutron star whose pair-plasma cloud surrounding the core of the star has a mean plasma number density close to $N_0 = 10^{30}$ cm^{-3} at 300 km away from the center of the star. This plasma density corresponds to the plasma frequency $\omega_p \approx 5.64 \cdot 10^{19}$ s^{-1}, which can be smaller than the gyrofrequency in some regions of the magnetized star.

Considering this scenario, we can assume the normalized gyrofrequency $\Omega_c = 10$, for instance. Figure 3a gives the normalized electric field associated with the upper-hybrid wakefield for the neutrino-driven term $S_\nu = 20$. We can see in this case that the electric field quickly reaches a huge saturation value, $E \approx 35$ (in terms of physical quantities, $E \approx 3.4 \cdot 10^{22}$ V/cm; we note that this value of the electric field is larger than the Schwinger limit for pair creation in the vacuum) for a small value of the phase χ, showing how intense the nonlinear behavior of the neutrino–plasma interaction is. In this extreme situation, the plasma particles are unlimitedly accelerated to very high energy, as we can see from Fig. 3b. Of course, the symmetry between the electron and positron dynamics is

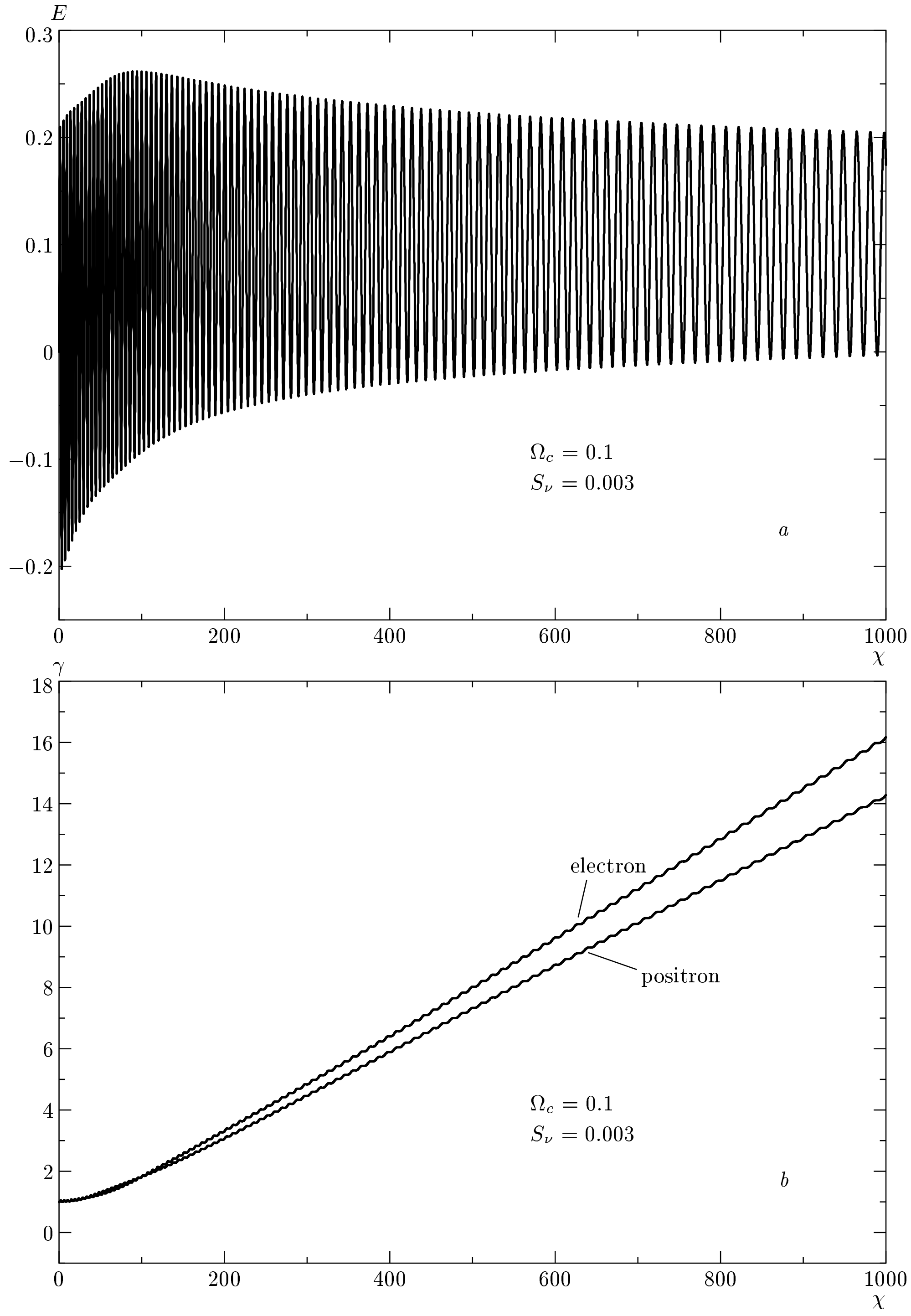

Fig. 1. Normalized electric field E (a), and normalized energy $\gamma = \sqrt{1 + P_x^2 + P_y^2}$ (b) for pair plasmas, versus the normalized distance (phase) χ for the normalized gyrofrequency $\Omega_c = 0.1$ and the neutrino-driven term $S_\nu = 0.003$

completely destroyed (see, e.g., Figure 3b) in the presence of these intense neutrino bursts, which means that the electron–plasmon cross section is now much larger than the positron–plasmon one.

5. CONCLUSIONS

In conclusion, we have presented a hydrodynamic description of large-amplitude upper-hybrid waves ex-

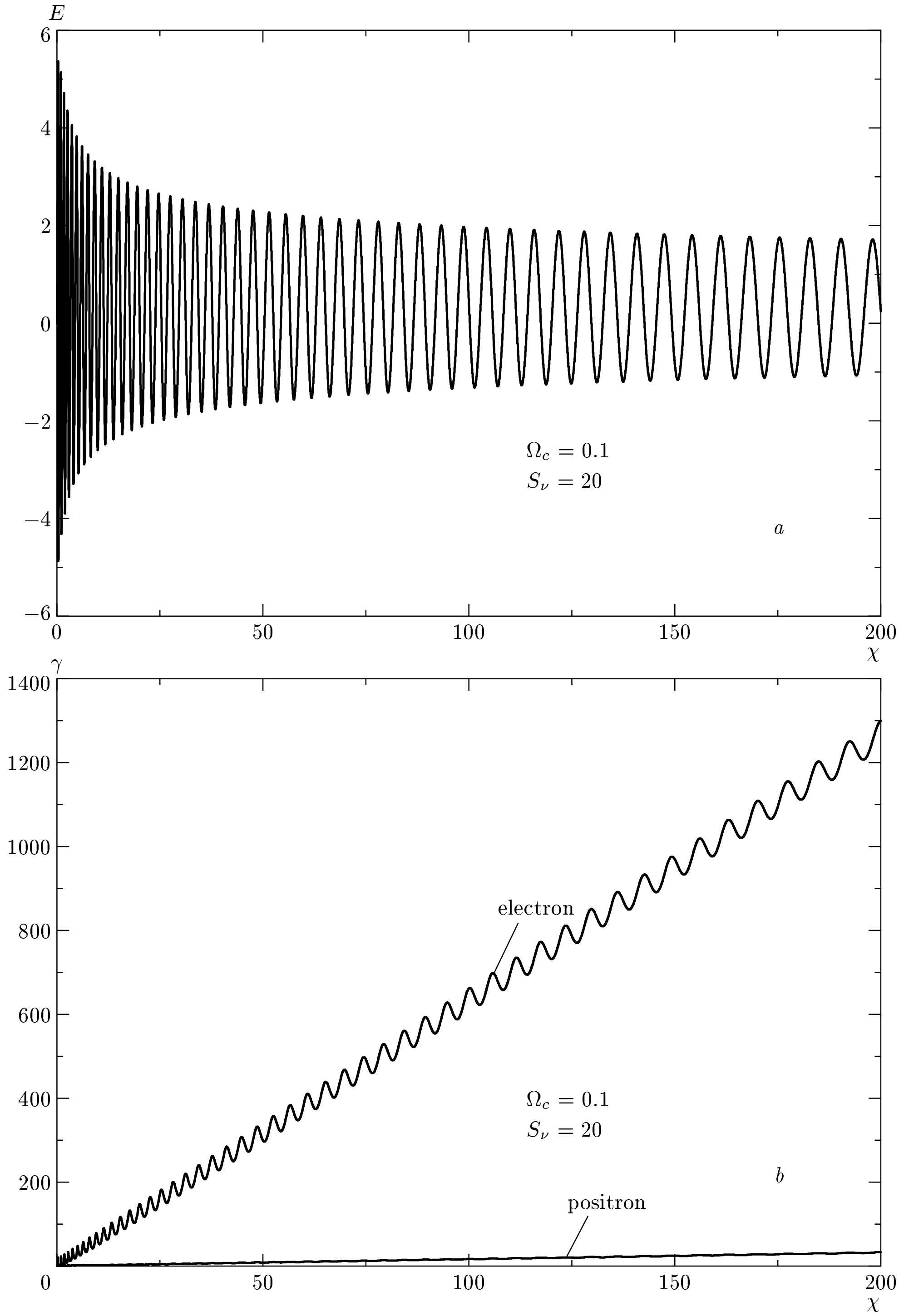

Fig. 2. Normalized electric field E (a), and normalized energy $\gamma = \sqrt{1 + P_x^2 + P_y^2}$ (b) for a pair plasma versus the normalized distance (phase) χ for the normalized gyrofrequency $\Omega_c = 0.1$ and the neutrino-driven term $S_\nu = 20$

cited by intense neutrino beams in a dense magnetized pair plasma. The excitation of wakefields in a pair plasma is possible because of the spontaneous breakdown of symmetry between the electron and positron dynamics due to the driving force of intense neutrino bursts. Physically, the symmetry breaking is attributed

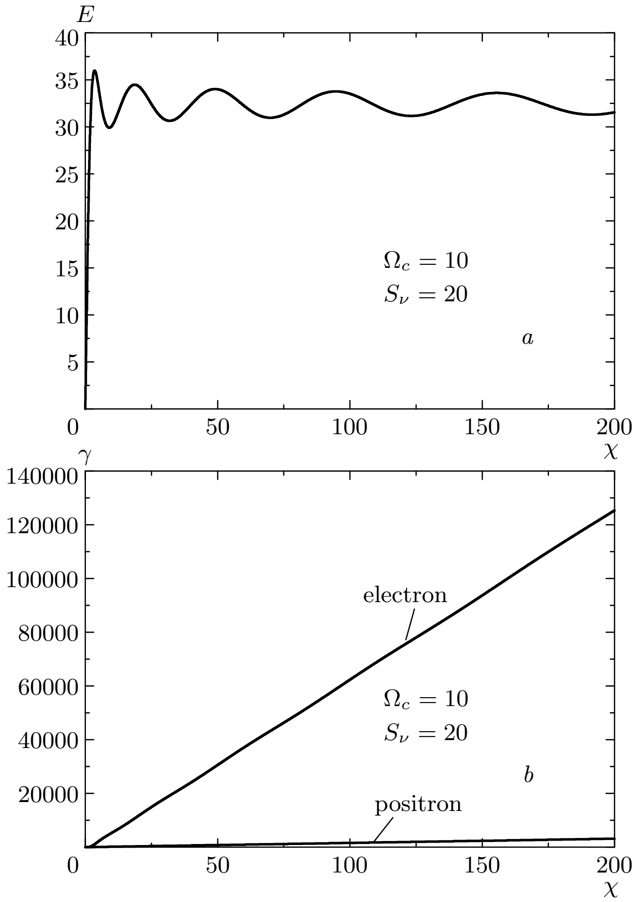

Fig. 3. Normalized electric field E (a) and normalized energy $\gamma = \sqrt{1 + P_x^2 + P_y^2}$ (b) for a pair plasma versus the normalized distance (phase) χ for the normalized gyrofrequency $\Omega_c = 10$ and the neutrino-driven term $S_\nu = 20$

to the induced negative charge that neutrinos acquire as they travel through the plasma. The induced negative charge pushes the electrons such that the electron momentum increases and the positrons are attracted by the effectively charged neutrinos. The resulting charge imbalance due to the charge separation, in turn, creates finite-amplitude wakefields. Our results, which are independent of the neutrino density but dependent on the neutrino energy deposited into the plasma, are valid for any astrophysical scenarios. They should be applied to understand the acceleration of pairs in stars as well as in ultrarelativistic pair-plasma jets that produce γ-ray bursts. Furthermore, the presence of an external magnetic field shows that we have the possibility of generating upper-hybrid wakefields that can transversely accelerate the pair plasma to ultrarelativistic energies by the surfatron mechanism [28].

By means of numerical computation, we have observed that the effects of the finite plasma and neutrino temperatures do not significantly affect the pair-plasma dynamics and the amplitude intensity of the electric field generated by the neutrino bursts. It should be pointed out that any amount of neutrino energy transferred to the plasma suffices to drive the wakefields. Of course, the amplitude of the wakefield depends on the amount of neutrino energy deposited into the plasma. Hence, the intensity of the generated electric field can reach values that allow generating pair creation by the wakefields. We are presently investigating the pair creation in magnetized plasmas.

This research was partially supported by the CNPq (Conselho Nacional de Desenvolvimento Cientìfico e Tecnológico) as CAPES (Coordenadoria de Apoio ao Pessoal do Ensino Superior) as well as by the European Commission (Brussels, Belgium) through contract № HPRN-CT-2001-00314 and the Sonderforschungsbereich 591 of the Deutsche Forschungsgemeinschaft (Bonn, Germany). We also acknowledge the support of Centre for Fundamental Physics at Rutherford Appleton Laboratory at Chilton, Didcot, UK.

REFERENCES

1. G. W. Gibbons, S. W. Hawking, and S. Siklos, *The Very early Universe*, Cambridge University Press, Cambridge (1983).
2. W. Misner, K. S. Thorne, and J. A. Wheeler, *Gravitation*, Freeman, San Francisco (1973), p. 763.
3. F. C. Michel, Rev. Mod. Phys. **54**, 1 (1982).
4. T. Tajima and T. Taniuti, Phys. Rev. A **42**, 3587 (1990).
5. N. J. Shaviv and A. Dar, Astrophys. J. **447**, 863 (1995); A. Dar, *ibid.* **500**, L93 (1998).
6. H. R. Miller and P. J. Witta, *Active Galactic Nuclei*, Springer-Verlag, Berlin (1987).
7. F. T. Gratton, G. Gnavi, R. M. O. Galvâ, and L. Gamberoff, Phys. Rev. E **55**, 3381 (1997).
8. J. G. Lomonadze, L. Stenflo, V. N. Tsytovich, and H. Wilhelmsson, Physica Scripta **26**, 455 (1982); A. C. L. Chian and C. F. Kennel, Astrophys. Space Sci. **97**, 9 (1983).
9. G. J. Fishman and C. A. Meegan, Annu. Rev. Astron. Astrophys. **33**, 415 (1995).

10. S. S. Gershtein, Private Communication; IHEP 2000-15.

11. P. K. Shukla, N. N. Rao, M. Y. Yu, and N. L. Tsintsadze, Phys. Rep. **138**, 1 (1986).

12. J. R. Wilson, in *Numerical Astrophysics*, ed. by J. M. Centrella, J. M. LeBlanc, and R. L. Bowers, Bartlett & Jones, Boston (1985), p. 422; H. A. Bethe and J. R. Wilson, Astrophys. J. **295**, 14 (1985).

13. J. Cooperstein, Phys. Rep. **163**, 95 (1988).

14. H. A. Bethe, Rev. Mod. Phys. **62**, 801 (1990).

15. S. Bludman, Da Hsuan Feng, Th. Gaisser, and S. Pittel (eds), Phys. Rep. **256**, 1 (1995).

16. V. N. Oraevsky, and V. B. Semikoz, Physica A **142**, 135 (1987).

17. J. F. Nieves, and P. B. Pal, Phys. Rev. D **49**, 1398 (1994).

18. N. L. Tsintsadze, J. T. Mendonça, and P. K. Shukla, Phys. Lett. A **249**, 110 (1998).

19. J. T. Mendonça, A. Serbeto, P. K. Shukla, and L. O. Silva, Phys. Lett. B **548**, 63 (2002).

20. L. O. Silva, R. Bingham, J. M. Dawson, J. T. Mendonça, and P. K. Shukla, Phys. Rev. Lett. **83**, 2703 (1999).

21. A. Serbeto, Phys. Plasmas **6**, 2943 (1999).

22. A. J. Bizard and J. S. Wurtele, Phys. Plasmas **6**, 1324 (1999).

23. P. Meszaros, *High Energy Radiation From Magnetized Neutron Stars*, University of Chicago Press, Chicago (1992), Ch. 2.

24. G. Raffelt, *Stars as Laboratories for Fundamental Physics*, University of Chicago Press, Chicago (1996).

25. A. A. Gvozdev and I. S. Ognev, Pisma Zh. Eksp. Teor. Fiz. **74**, 298 (2001).

26. J. Pantalone, Phys. Lett. B **342**, 250 (1995).

27. A. Serbeto, L. A. Rios, and P. K. Shukla, Phys. Plasmas **9**, 4406 (2002).

28. T. Katsouleas and J. M. Dawson, Phys. Rev. Lett. **51**, 392 (1983).