relative to a change of the magnetic field, as is indeed observed in oriented polyethylene.

Simple calculations were made of the intensities, in accordance with [4], for all 13 x 6 = 78 lines, and the plotted envelope of these lines is in good agreement with the shape of the observed EPR spectrum.

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PRODUCTION OF HIGH-ENERGY PARTICLES BY ACCELERATION OF A PLASMA SCATTERING A STRONG BEAM OF FAST ELECTRONS

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The development of powerful pulsed electron accelerators with currents $10^5 - 10^6$ A and voltages 0.3 - 10 MV, with pulse durations of several dozen nanoseconds [2-4] has made it possible to employ now the previously proposed [1] method of accelerating a plasma with a frozen-in magnetic field, scattering an electron current, to accelerate the plasma to relativistic velocities and to obtain high-energy particles (in [1] we considered the nonrelativistic case).

Let us consider a beam of relativistic electrons incident on a plasmoid of sufficiently high density, whose scattering ability is enhanced by a frozen-in magnetic field. When the electrons are scattered by the magnetic field, the plasma receives the momentum of the electron current, and the plasma will be accelerated as a whole, provided a sufficiently strong coupling obtains between the electronic and ionic fractions of the plasma, i.e., if the resultant accelerating field E_{acc} does not exceed the maximum Coulomb field that can arise as a result of the separation of the plasma charges, $E_p \approx 4\pi n_e ed$, where n_e is the plasma density and d is the length of the plasmoid. (For example, if $E_{acc} = 10^9$ V/m we need $n_e \gtrsim 10^{12}$ cm^{-3} at a plasmoid length d = 5 cm.)

The magnetic field frozen in the plasma should have an intensity sufficient for strong scattering or reflection of the electron current: from the condition $eH\rho \approx mvc \approx \mathcal{E}$ we find the required $H \approx \mathcal{E}/ed \approx 0.6 \times 10^3 \mathcal{E}_{MeV} \approx 10^3$ Oe at an electron energy $\mathcal{E} = 3$ MeV. If the force lines of the frozen-in magnetic field emerge from the plasmoid, then the cross section of the plasmoid can be much smaller than the cross section of the electron beam and the efficiency of momentum transfer can still remain adequate. The duration of the acceleration should not exceed noticeably the skin-effect damping time of the currents in the plasma ($t \approx 4\pi\sigma d^2/c^2$, where σ is the electric conductivity of the plasma), if no special measures are taken to maintain the currents (say by means of a varying external magnetic field or by deformation of the plasmoid), or else should not exceed the time of spreading of the plasma under the influence of magnetic or thermal pressure. The pressure of the scattered electron stream can limit the transverse spreading and ensure inertial longitudinal condensation of the plasmoid, forming a "bowl" of accelerating forces.

Let us estimate the dynamics of the plasmoid acceleration. Assuming that in a coordinate frame traveling with the plasmoid the electron momentum p_x' changes strongly upon collision with the plasmoid (for example, it is reversed), we obtain for the change of the electron momentum in the laboratory frame

$$\Delta p_x = 2p_x'/\sqrt{1-\beta^2},$$

but

$$p_x' = \frac{p_x - v\mathcal{E}/c}{\sqrt{1-\beta^2}} = \frac{mc(\beta_e - \beta)}{\sqrt{1-\beta_e^2}\sqrt{1-\beta^2}}$$

therefore

$$\Delta p_x = \frac{2}{1-\beta^2}\frac{mc}{\sqrt{1-\beta_e^2}}(\beta_e - \beta).$$

Here $\beta_e = v_e/c$ is the velocity of the stream electrons and $\beta = v/c$ is the velocity of the plasmoid in the laboratory frame.

The change of the energy of plasmoid motion per unit path is

$$\frac{d}{dL}\frac{Mc^2}{\sqrt{1-\beta^2}} = Mc^2\frac{\beta}{(1-\beta^2)^{3/2}}\frac{d\beta}{dL} = \Delta p_x \nu = \Delta p_x N_1(\beta_e - \beta)c = \frac{2mc^2 N_1(\beta_e - \beta)^2}{(1-\beta^2)\sqrt{1-\beta_e^2}},$$

where the frequency of the impacts of the electrons against the plasmoid is $\nu = N_1(v_e - v)$, N_1 is the running number of electrons, and M is the total rest mass of the plasmoid, or

$$\frac{\beta\, d\beta}{(1-\beta^2)^{1/2}(\beta_e - \beta)^2} = A\,dL,$$

where

$$A = \frac{2mN_1}{M\sqrt{1-\beta_e^2}}.$$

The integration is carried out by making the substitution $\beta = \sin\phi$, and yields

$$\frac{1}{(1-\beta_e^2)}\left\{\frac{\beta_e\sqrt{1-\beta^2}}{(\beta_e - \beta)} - 1 - \frac{1}{(1-\beta_e^2)^{1/2}} \times \times \ln\left[\frac{1-\beta_e(1-\sqrt{1-\beta_e^2})/\beta(1+\sqrt{1-\beta_e^2})}{1-\beta_e(1-\sqrt{1-\beta_e^2})/\beta(1-\sqrt{1-\beta_e^2})}\right]\right\} = AL.$$

We introduce $\gamma = 1/\sqrt{1-\beta^2}$ (i.e., $\beta = \sqrt{\gamma^2 - 1}/\gamma$); then

$$\ln[\ ] = \ln\frac{1 - \left(\frac{1-1/\gamma}{1+1/\gamma}\right)^{1/2}\left(\frac{1-1/\gamma_e}{1+1/\gamma_e}\right)^{1/2}}{1 - \left(\frac{1-1/\gamma}{1+1/\gamma}\right)^{1/2}\left(\frac{1+1/\gamma_e}{1-1/\gamma_e}\right)^{1/2}}.$$

$\ln[\ ] = \ln[(\gamma_e + \gamma)/(\gamma_e - \gamma)]$ when $\gamma \gg 1$ and $\ln[\ ] \to 2\gamma/\gamma_e$ when $\gamma_e \gg \gamma$. In this case we get

$$2\gamma^3 \asymp AL = \frac{2mN_1 L}{M\sqrt{1-\beta_e^2}},$$

i.e.,

$$\gamma = \frac{1}{\sqrt{1-\beta^2}} = \left\{\frac{mN_1 L}{M\sqrt{1-\beta_e^2}}\right\}^{1/3}.$$

For example, when

$$\mathcal{E} \asymp mc^2/\sqrt{1-\beta_e^2} \asymp 10 \text{ MeV},$$

a running electron density in the beam $N_1 \approx I/ec \approx 2 \times 10^{14}$ el/cm (current $I = 10^6$ A), a plasma rest mass $M \approx d^3 n_e m_p/2 \approx 3 \times 10^{-10}$ g, and an acceleration path $L = 10$ m (this path may exceed by many times the length of the electron stream when $\beta_e \to \beta \to 1$) we get $\gamma \approx 3$, corresponding to a proton energy $\gamma m_p c^2 \approx 3$ GeV.

We note that at large values of N_1 the induction rest mass of the electron is

$$m = m_0(1 + N_{10} r_0 L_{10}) >> m_0,$$

where $N_{10} = N_1/\gamma_e$, r_0 is the classical radius of the electron, L_{10} is the running inductance in the rest system of the jet ($N_{10\ KD} \approx 1/r_0 L_{10} \approx 1/r_0 \ln(\ell_0/d) \approx 1/r_0 \ln(\ell\gamma_e/d) \approx 10^{12}$ el/cm) and ℓ is the length of the electron jet in the laboratory system, which greatly decreases γ_e at a specified value of $m\gamma_e$ during the initial stage of the acceleration, when this does not matter, but γ_e continues to increase with decelerating current, owing to self-induction, thereby improving the acceleration conditions during the relativistic stage. Self-induction effects occurring in the deceleration of powerful electron beams can be so large that the remaining part of the electrons can be accelerated to high energies (see [5] concerning self-induction acceleration). The reflection of the electron stream by the frozen-in field may turn out to be more convenient than the reverse Cerenkov acceleration of a plasmoid in an electron stream proposed in [6] (in our case we use the stronger mechanism of interaction between the electron stream and the plasma).

There are possible variants of more intensified interaction in which collective effects are excited in the plasmoid, but they heat the plasmoid and accelerate its disintegration. In the case of a frozen-in magnetic field it is possible to avoid collective effects, for example, by preventing the electron beam from entering the plasma, and scattering the electrons by the magnetic-field force lines emerging from the plasmoid.

Axial stability of the plasmoid can be ensured either by acceleration in a metal tube that repels from its walls the plasmoid with the field lines emerging from it (the repulsion of the plasmoid magnetic field from the conducting walls may be comparable with the accelerating force, since the magnetic pressure causing the plasmoid to be pushed away from the walls is close to the accelerating magnetic pressure of the electron current), or by producing a special magnetic field to press the plasmoid towards the axis, or else by choosing the profile of the electron stream such as to ensure both acceleration and axial stability of the plasmoid.

We note that the considered acceleration mechanism is the inverse of the Fermi acceleration mechanism (in our case we use the reaction of the recoil due to particle reflection from magnetic clouds).

It is possible to realize an explosive-reaction acceleration of a plasmoid with a frozen-in magnetic field by unilateral heating of the plasmoid by means of an electron or microwave beam and pressure of the heated layer on the remaining part of the plasma (the frozen-in magnetic field makes the plasma impenetrable to the expanding heated plasma layer). The pressure of the heated plasma can in this case greatly exceed the pressure of the electron beam itself [7].

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FLEXURE-DRIFT RESONANCE IN SEMICONDUCTORS

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In semiconductors with carriers of one or both signs, in the presence of an external electric field (E) and in general also a magnetic field, there occur the so-called drift waves of the carrier density and field. In the absence of temporal and spatial dispersion, the frequency of these waves equals $\omega = k_i\mu_{ik}E_k = (\vec{k}\cdot\vec{v})$ ($\vec{k}$ - wave vector, μ_{ik} - carrier mobility in the case of charged waves or ambipolar mobility in the case of quasineutral waves). If the frequency greatly exceeds the damping decrement $4\pi\sigma/\varepsilon + Dk^2$ (charged waves) or Dk^2 (quasineutral waves), then these waves attenuate weakly (D, ε, σ - diffusion coefficient, dielectric constant, and electric conductivity). At definite values of the wave vector, the drift branch of the crystal plasma oscillations can cross the branches corresponding to other osicllations in the crystal, provided the frequencies of the latter depend on the wave vector in nonlinear fashion; for example, it can cross the spin, plasma, optical, or flexural waves. If these waves cause crystal plasma oscillations, and the crossing of the frequencies occurs in a region where the wave damping is weak, then it is possible to have a resonant increase of the absorption of the indicated waves, on the one hand, and in the case of non-ohmic contacts also a resonant change of the impedance of the crystal.

We investigate in this paper resonance of a charged drift wave with a flexural wave in the presence of a piezoelectric interaction, in crystals having carriers of one sign. We consider a plate of thickness h (z axis) and length L (x axis) in an external electric field $E = E_x$, with $L \gg h$. If the piezoelectric tensor $\beta_{ik,\ell}$ is homogeneous in the z direction, then it can be shown that the flexural oscillations of the plate do not interact with the drift wave in first approximation in $\beta^2/\rho s^2$. We consider a plate with different values of $\beta_{i,k\ell}$ when $z < 0$ and $z > 0$. Using the continuity and Poisson equations and the equation for