of the nuclear absorption ($\varepsilon$ = 5% and not zero at $\theta = \theta_B$) is due mainly to the insufficiently small angular divergence of the $\gamma$ rays incident on the crystal. 2) The $T_{res}(\theta)$ dependence has a somewhtt asymmetrical dispersion form, and the $T_{res}$ levels to the right and to the left of $\theta_B$ are different. This is due to the fact that the crystal investigated by us is not thick enough. According to the theory [1], diffraction is accompanied, besides absorption by the wave state, also by formation of the strongly absorbable state, which is not completely absorbed in our crystal and produces the asymmetry of $T_{res}(\theta)$.

For our Fe + 3% Si crystal, the ratio of the amplitudes of the scattering of the $\gamma$ rays by nuclei and electrons is $f_{nuc}/f_{el} \approx 0.08$. Therefore the wave field in the crystal is formed mainly as a result of the diffraction by electrons. We are presently carrying out measurements on an $\alpha$-$Fe_2O_3$ crystal made of enriched iron (~85% $Fe^{57}$). In this case $f_{nuc}/f_{el} \approx 10$ and the effect of normal transmission of the resonant $\gamma$ rays is due mainly to diffraction by the nuclei.

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# ACCELERATION OF COSMIC ELECTRONS AND POSSIBLE MECHANISM OF FORMATION OF THE SPECTRUM OF METAGALACTIC X-RADIATION

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The problem of the mechanism whereby the cosmic-electron spectrum is produced has by now become very important. On the one hand, it is closely connected with the problem of the origin of cosmic rays, and on the other it is of independent significance, since the electrons are the sources of the background radiation of the universe. Apparently, the diffuse x-radiation of the metagalaxy is produced by Compton scattering of relativistic electrons by radiation of longer wavelength (relict, infrared, etc.). The spectrum of the diffuse x-radiation of the metagalaxy has been investigated in sufficient detail. For photons in the energy interval 1.5 < E < 40 keV, it is described by the expression $I(E) < E^{-\alpha}$ at $\alpha \approx 0.7 - 0.8$. At E ≈ 40 keV, the spectrum experiences a kink, after which $\alpha$ becomes equal to 1.2 - 1.3. At low energies, there is only one experimental point at E = 280 eV. The high value of the x-ray flux at this point has led a number of authors [1, 2] to the conclusion that there exists one more kink in the energy region below 1 keV. However, the available experimental data are not sufficient to make the conclusion unambiguous [3]. Most authors [2 - 4] have reached the conclusion that the kink at E ≈ 40 keV is due to the kink in the spectrum of the electrons responsible for the x-radiation. The spectrum of these electrons corresponds to the spectrum of the electrons of our galaxy, which was determined both on the basis of radio-emission data and by direct measurements [5, 6]. However, the mechanism of formation of this spectrum was not considered.

We shall show that it is precisely such a spectrum that can arise in the case of acceleration and scattering of electrons by inhomogeneities of the magnetic field of a moving plasma. A detailed theory of the scattering of charged particles by moving magnetic inhomogeneities is contained in a number of papers by Dolginov and Toptygin [7]. An equation describing the scattering of relativistic electrons by magnetic clouds [7] with allowance for the energy loss is given by

$$\frac{\partial n}{\partial t} = \frac{\partial}{\partial \epsilon}\left[\frac{\overline{\Delta u^2}}{3c\Lambda(\epsilon)}\,\epsilon^4\,\frac{\partial}{\partial \epsilon}\left(\frac{n}{\epsilon^2}\right)\right] - u\,\frac{\partial n}{\partial r} - \frac{\partial}{\partial \epsilon}\,[b(\epsilon)n] + \frac{2u}{3r}\,\epsilon^3\,\frac{\partial}{\partial \epsilon}\left(\frac{n}{\epsilon^2}\right) + \frac{1}{3}\,c\Lambda(\epsilon)\nabla^2 n\,. \tag{1}$$

Here $n \equiv n(r, \varepsilon, t)$ is the density of electrons with energy $\varepsilon$ at the point $r$. We assume a radial spreading of the magnetic inhomogeneities with a velocity $u$, which is subject to fluctuations $\Delta u$. Since the principal role is played by losses to synchrotron radiation and to the inverse Compton effect, it follows that $b(\varepsilon) = b_0\varepsilon^2$. The second and last term in the right side of (1) describe the loss of the electrons from the acceleration region by convective transport and by diffusion, respectively. $\Lambda(\varepsilon)$ is the mean free path of the electron relative to scattering. The electron energy is in units of $mc^2$, where $m$ is the electron mass.

At $u = 0$ and $\Lambda$ = const, Eq. (1) goes over into a well known equation describing the Fermi acceleration of the electrons. The electron free path depends on their Larmor radius $r_L$ and on the spectrum of the magnetic inhomogeneities. This spectrum is characterized, first, by the correlation length, i.e., by the spatial scale within which a noticeable change of the magnetic field takes place, and second, by the relative contribution of the irregularities with different correlation lengths. If the principal role is played in the spectrum by irregularities having a definite dimension $L_c$, then $\Lambda$ decreases with increasing energy so long as $r_L < L_c$, depends very little on the energy when $r_L \sim L_c$, and increases with energy when $r_L > L_c$. A similar dependence holds also for cosmic protons in the scattering of solar-wind plasma by magnetic fields.

Let us solve Eq. (1) under the assumption that $\Lambda(\varepsilon) = \Lambda_0\varepsilon^{-\beta}$, where $\beta \neq 0$. The energy loss is most appreciable at large electron energy, whereas at low energy the escape of electrons from the acceleration region may become more important. At a certain energy $\varepsilon_1$, the characteristic time for energy loss $t_\varepsilon$ becomes comparable with the time $t_e = r/u$ of the escape of electrons carried away by the outflowing plasma. For an energy $\varepsilon < \varepsilon_1$ we can retain the first two terms in the right side of (1), and when $\varepsilon > \varepsilon_1$ only the third term.

A solution shows that when $\beta \leq 0$ it is impossible to obtain the observed electron spectrum. For $\beta > 0$, the analytic solution of Eq. (1) has under the stationarity condition the form

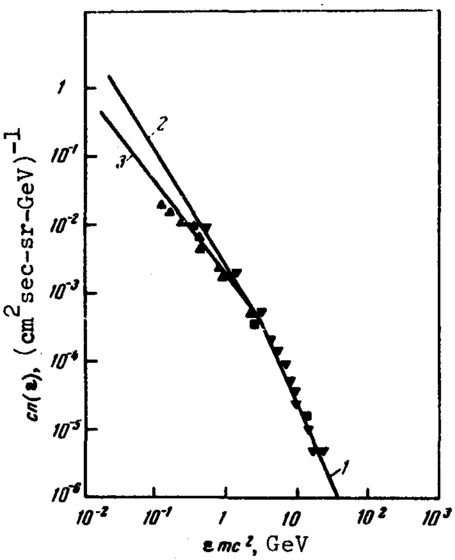
$$n(R_0, \epsilon) = \frac{A}{\beta\epsilon_0 \tau}\exp -\left(\frac{1}{\beta^2\tau}\right)\left(\frac{\epsilon}{\epsilon_0}\right)^{(1-\beta)/2} I_{1+(3/\beta)}\left[\frac{2}{\beta^2\tau}\left(\frac{\epsilon_0}{\epsilon}\right)^{\beta/2}\right]. \tag{2}$$

Here $I_\nu(x)$ is the Bessel function of imaginary argument. $\tau = R_0\overline{\Delta u^2}/3c\Lambda(\varepsilon_0)u$, with $\beta^2\tau < 1$. $R_0$ is the characteristic dimension of the acceleration region.

$\varepsilon_0$ is the energy at which the assumption that $\Lambda = \Lambda_0 \varepsilon^{-\beta}$ becomes valid, $\beta$ being assumed constant in the entire energy region under consideration up to $\varepsilon_2 \sim 100$ GeV. The initial distribution of the accelerated electrons is assumed to decrease steeply with energy (e.g., the tail of the Maxwellian distribution)

$$A = \int_0^\infty n(0, \varepsilon)\, d\varepsilon .$$

Using the asymptotic form of the Bessel function, it can be easily shown that when $\varepsilon \gg \varepsilon_0$ it follows from (2) that the electrons have a power-lay energy spectrum: $n(R_0 \varepsilon) \sim (\varepsilon_0/\varepsilon)^{1+\beta}$. In the region $R_0 < r \leq R_1$, the energy losses cause the electron spectrum to become steeper, and its exponent increases by unity [2, 3], namely, $n(R_1 \varepsilon) \sim (\varepsilon_0/\varepsilon)^{2+\beta}$ when $\varepsilon \geq \varepsilon_1$. The figure shows a comparison of the theoretical spectrum with the observed spectrum of the cosmic electrons of our galaxy [5, 6]. A comparison makes it easy to determine the physical parameters of the model:




1 – $n(\varepsilon) \sim \varepsilon^{-2.5}$, 2 – $n(\varepsilon) \sim \varepsilon^{-1.6}$, 3 – $n(\varepsilon) \sim \varepsilon^{-1.4}$. 2 and 3 correspond to the electron spectra calculated on the basis of radio observations [5, 6].

$$\beta \frac{R_{2-i}}{u} (H^2 + 2 \cdot 10^{-11} \rho)\, \varepsilon_i = 5 \cdot 10^8 ; \tag{3}$$

$$\frac{3c\,\Lambda(\varepsilon_0)\, u}{\beta^2 R\, \overline{\Delta u^2}} \left(\frac{\varepsilon_0}{\varepsilon}\right)^{\beta/2} < 1, \quad i = 1,2 .$$

Here $\rho$ is the radiation-energy density in eV/cm$^3$ and $R_1$ is the dimension of the region around the electron source where the effective time of the energy loss is smaller than the escape time. The best agreement is obtained at the following values of the parameters: $\beta = 0.4 - 0.6$, $H \sim 10^{-3}$ G, $\Delta u \sim u \sim 3 \times 10^7$ cm, $R_0 \sim 10^{17}$ cm, $R_1 \sim 3 \times 10^{18}$ cm, $\Lambda(\varepsilon_0) \sim 10^{15}$ cm, $\varepsilon_0 > 10^{-2}$. These values are close to the values of the corresponding parameters for the envelopes of supernova stars.

As another possibility, let us consider the discrete sources of powerful infrared radiation of the type recently observed near the galactic center [8, 9]. M.S. Longair and R.A. Syunyaev have already discussed the possible role of these objects in the formation of the background x-radiation [2]. The specified values of the parameters of these objects were determined from observations [8] (luminosity $L \sim 10^7 L$, $R_L \sim 3 \times 10^{18}$ cm, $u \sim 10^8$ cm) make it easy to reconcile the theoretical spectrum with the experimental one (see the figure). The spectrum of the diffuse x-radiation, under the assumption that the spectrum of the electrons corresponds to that shown in the figure, was calculated in [3] and agrees well with the observation results.

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## METASTABLE "OVERTURNED" STATE OF FERROMAGNET AS A SOURCE OF AMPLIFICATION OF ELECTROMAGNETIC OR ACOUSTIC WAVES


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It is known [1] that a uniaxial ferromagnet in a magnetic field parallel to the easy magnetization axis has two equilibrium states if $H < \beta M_0$, where H is the intensity of the magnetic field in the magnet, $M_0$ is the value of the spontaneous magnetic moment, and β is the anisotropy constant (β > 0). In one of these states (thermodynamically stable), the vector of the magnetic moment $\vec{M}$ is directed along the field $\vec{H}$, and in the other (thermodynamically metastable) it is directed against the field $\vec{H}$. If $H > \beta M_0$, then only one stable direction of $\vec{M}$ along the field $\vec{H}$ remains.

We wish to call attention to the fact that a metastable "overturned" state of a ferromagnet (a ferrodielectric) can be readily used for amplification (to a considerable degree, coherent) of electromagnetic waves or acoustic oscillations. Two methods of such a utilization of the "overturned" state can be indicated.

1. Let the z axis be the easy magnetization axis, and let the magnetic moment of the crystal $\vec{M}$ be parallel to this axis: $\vec{M} = -M_0\vec{n}$ ($M_0 > 0$), where $\vec{n}$ is the unit vector along the z axis. We assume that the external constant magnetic field is directed opposite to $\vec{M}$ ($\vec{H} = H_0\vec{n}$, $H_0 > 0$) and is such that $H_0 - \beta M_0 = H_1 < 0$ and $|H_1| \ll H_0 \simeq \beta M_0$. In the region of low temperatures, the stability of the "overturned" state relative to thermal fluctuations will be conserved also at very small $H_1$.

We now turn on a weak magnetic field $\vec{h}$, parallel to the z axis and increasing slowly with time. Assume that at the instant of time when the condition $H_1 + h(t) = 0$ is satisfied, the "overturned" state becomes unstable and the ferromagnet goes over, after a suitable relaxation time τ, into a stable equilibrium state with a vector $\vec{M}$ directed along $\vec{H}$. If $g\beta M_0\tau \gg 1$, where g is the gyromagnetic ratio, then the magnetic moment $\vec{M}$ will precess during the course of such a transition about the direction $\vec{n}$ with increasing frequency, approaching the resonant frequency $\omega_0 = g(H + \beta M_0) = 2g\beta M_0$. It is clear that this will be accompanied by electromagnetic radiation and amplification of the high-frequency electromagnetic wave with the required circular polarization and the frequency $\omega \leq \omega_0$, satisfying the condition $\omega\tau \gg 1$, is possible. The spectral characteristics of the radiation or of the amplification can be analyzed in analogy with the description given in [2]. The maximum amplification should occur when $\omega = \omega_0$. We note that the resonant frequency $\omega_0$ can be varied by changing the inclination of the magnetic field $\vec{H}$ to the easy-magnetization axis.