PULSATIONS OF TYPE IV SOLAR RADIO EMISSION: THE BOUNCE-RESONANCE EFFECTS

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Abstract. A mechanism of excitation of radial oscillations of a 'magnetic tube' is proposed for the interpretation of a periodic modulation of type IV radio burst intensity in the meter and decimeter range. After the flare a configuration with denser plasma extended along the magnetic field can be formed in the corona. Eigenoscillations of such a system are damped by MHD-wave emission into the external coronal plasma. However, if high energy protons with $\beta \gtrsim 0.2$ are trapped by this configuration, the damping of oscillations can be made up for by an amplification due to bounce-resonant plasma instability. The regularity of the pulse period is explained by presence of a maximum in the wave growth rate dependence on the frequency.

1. Introduction

Solar type IV radio bursts in the meter and decimeter wavelength region are sometimes characterized by a periodic or quasiperiodic intensity modulation (Hughes and Harkness, 1963; de Groot, 1970; McLean *et al.*, 1971; Abrami, 1972). Generally accepted models of the periodic pulsations are based upon the hypothesis, that the radio emission is modulated by MHD-oscillations (Rosenberg, 1970; McLean *et al.*, 1971). The sources of type IV emissions, i.e. the magnetic traps in the high corona can have the form of 'magnetic tubes'* and be the resonator for MHD-waves. However, the quality of such resonators is finite and energy feeding to the wave is required for the observations to be adequately interpreted. The plasma instabilities arising due to the presence of a group of fast particles (electrons, protons), mirroring in a trap, can act as such a source of energy. We wish to discuss here the possibility for bounce-resonant interaction between the MHD-waves and fast particles to compensate the emissive damping of oscillations in the source of the type IV radio emission.

2. Observations and Interpretations

The period of pulsations of the type IV radio bursts depends on the wavelength

* The term 'magnetic tube' has no meaning unless one implies inhomogeneity of the density, perpendicular to the magnetic field (see below).

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range. At the decimetric wavelengths the period is of the order of several tenths of a second, at metric ones it comes up to several seconds (Achong, 1974). The pulsation phase usually lasts for several minutes. The modulation of the continuum is about 5 to 10% of the mean level, but sometimes it reaches 50%.

At the present time there are two mechanisms of IV type radio emission complementing each other. There are the synchrotron and the plasma mechanisms. The former attributes the radio bursts to the synchroton radiation of relativistic electrons injected into the coronal magnetic trap (Boischot and Denisse, 1957; Zheleznyakov, 1970). The high frequency instabilities exciting plasma waves can also arise in such a system due to the existence of a 'loss-cone' for the trapped particles. The conversion of plasma waves into electromagnetic waves may result in the radio emission. This is the plasma mechanism of type IV emissions (Stepanov, 1973; Kuijpers, 1974).

To explain the pulsations, Rosenberg (1970) considered the modulation of the radio emission caused by an oscillating magnetic field in the source. In Rosenberg's model the source was treated as an infinitely long straight magnetic tube. The change in the magnetic field strength occurred with a period of eigenoscillations of the tube (magnetoacoustic, or fast mode), i.e. with the period of $T \simeq R/(C_A^2 + C_S^2)^{1/2}$. Here R is the tube radius, $C_S = (\gamma T/M)^{1/2}$ is the velocity of sound, $C_A = B_0/(4\pi\rho_i)^{1/2}$ is the Alfven velocity; T, plasma temperature; M, proton mass; γ, specific heat ratio; ρ_i, plasma density inside the tube. The pulse period coincides with observational values provided $R \approx 10^8$–10^9 cm.

In Rosenberg's model the plasma outside the tube was not taken into account. The ambient plasma effect, as was shown by Zaitsev and Stepanov (1975a), leads to dissipationless (emissive) damping of a fast mode because of its emission into the ambient plasma. The ratio of the oscillation decrement to the frequency for sufficiently long tubes is estimated as

$$|\nu|/\omega \sim \rho_e/\rho_i, \tag{2.1}$$

where ρ_e is the plasma density outside the tube. It follows from (2.1) that without energy feeding the high quality of observed pulsating sources cannot be explained, because it would require a very large difference between the external and internal plasma densities. No emissive damping exists for the Alfvén and slow magnetoacoustic waves with a group velocity directed along the magnetic field. The typical periods of these waves, however, are defined by the tube length $L \gg R$ and considerably exceed the values observed (Zaitsev and Stepanov, 1975a). Further, according to McLean *et al.* (1971), MHD-oscillations of the tube can be excited by the action of a shock-wave propagating through the corona. It is questionable, however, whether the basic mode of the high-quality system can be effectively excited by an external agent (reciprocity theorem).

Because of these difficulties involved in the pulsation interpretation with the help of MHD-waves, Zaitsev and Stepanov (1975b) proposed a pulsation model based on a plasma mechanism of type IV emission. In this model, the intensity of plasma turbulence, generated by energetic protons in the coronal magnetic trap, pulsates, the pulsations causing the radio emission modulation. In a similar way, Zaitsev (1970) explained the 'rain-type' bursts. The high frequency loss-cone instability in such a system can also develop in a pulsating regime (Trakhtengerts, 1968). In both cases the cause of modulation is connected with induced scattering of plasma waves by the thermal coronal particles and the subsequent wave damping in the nonresonant region. In comparison with the 'MHD-model' of pulsations, the 'plasma model' has the advantage of being not critical to the magnetic field structure. In this model, the pulse period is governed by the collision frequency of thermal electrons with ions ν_{ei}, and by the parameters of an energetic particle stream.

On the other hand, since the pulse period $T \sim \nu_{ei}^{-1}$ the plasma model can hardly explain the events with invariable period in a wide (>100 MHz) frequency band. Therefore Rosenberg's model of MHD-modulation explaining monochromativity of pulsations, continues to be attractive. However, as mentioned above, this model should be supplemented with an energy source for magnetoacoustic oscillations (MA).

The amplification of oscillations is possible if there is a resonant interaction between oscillations and magnetically trapped energetic particles, accelerated in the process of a flare. There is a wide class of resonant plasma instabilities arising due to the presence of a small fraction of fast particles in the plasma (see, e.g., Mikhailovsky, 1974). When studying plasma instabilities in adiabatic traps, a consistent account of particle bouncing in longitudinally nonuniform magnetic field is of importance. If the wave frequency, ω, is comparable with the typical bounce frequency of fast particles, Ω, the bounce-resonance effects should be considered. A number of works was devoted to various aspects of this interaction. Investigations were carried out into the role of bounce-resonance in the problem of particle transport in the Earth's magnetosphere (Dungey, 1965; Roberts and Schulz, 1968; Eviatar and Schulz, 1969; Tamao and Ishihara, 1976; Meerson, 1976; Meerson and Pokhotelov, 1978) and of the bounce-resonance effects on the plasma ULF-waves evolution (Southwood *et al.*, 1969; Coppi, 1970; Southwood, 1973; Karpman *et al.*, 1977). Recently the nonlinear and quasilinear theories of bounce-resonance were developed (Meerson and Sasorov, 1978; Meerson and Pokhotelov, 1978).

In the case under consideration the bounce-resonance condition takes the form

$$\omega = n\Omega , \quad n = 1, 2, \ldots . \tag{2.2}$$

For MHD-waves in the source of type IV emission this condition implies that

$$(C_A^2 + C_S^2)^{1/2}/R \simeq v_T/L \quad \text{for} \quad n \sim 1 , \tag{2.3}$$

v_T being a typical fast particle velocity. For $C_A/R \simeq 1\ \mathrm{s}^{-1}$ and $L \simeq 10^{10}$ cm the particle velocity v_T must be of the order of $10^{10}\ \mathrm{cm\ s}^{-1}$. We see that for high energy particles the bounce-resonance condition can be fulfilled. Thus, in the radio emission from the solar corona the bounce-effects may occur. Specifically these effects may play an important role in causing pulsations of type IV radio emission.

3. The Pulsating Source Model

In this section we describe the main special features of the proposed mechanism of the intensity modulation of type IV radio bursts.

3.1. The formation of the source

Let us assume that in a high region of the corona (at the distance of $\sim 1R_\odot$ from the photosphere) 'magnetic tube' with denser plasma is formed. Such a configuration may result from a flare ejection out of the chromosphere along a magnetic arch of an active region. If a transverse size of the inhomogeneity in density is $R \sim 3\times10^8$ cm, the mass of ejection $\sim 3\times10^{12}$ g is enough for a plasma density within the tube to reach $\rho_i \sim 10^8$ cm^{-3}. Such an ejection from the flare region is quite real (see, e.g., Zirker, 1971). The ejection motion is attended with a shock-wave responsible for a type II radio emission. After the tube is full of plasma, the density inside the tube ρ_i becomes 10–100 times higher than the corresponding value outside the tube, ρ_e. In this case the configuration of a magnetic field with a strength of 3–10 G is disturbed but slightly ($\beta_c \ll 1$), while the hydromagnetic stability of such a system can be provided by the stabilizing effect of tube ends frozen into the photosphere (Kadomtsev, 1966).

To motivate our assumption of an appreciable ratio of ρ_i/ρ_e it would be desirable to obtain at least indirect estimation of this ratio from observational data. If we attribute the type II and type IV radio events described by McLean *et al.* (1971) to the plasma emission mechanism, then a comparison of frequencies of type II ($f_{\rm II} \sim 30$ MHz) and type IV ($f_{\rm IV} \sim 150$ MHz) events gives the ratio $\rho_i/\rho_e \simeq 2(f_{\rm IV}/f_{\rm II})^2 \simeq 50$. Here we assume the electron density at the shock-wave front being twice as high as the corresponding value for the undisturbed corona. This estimate indicates that ratio ρ_i/ρ_e may really be of an appreciable value. Let us consider the oscillatory processes which may arise in such a system.

3.2. The source quality

A real structure of a magnetic tube with a dense plasma is rather complicated. To determine the eigenmode spectrum of such a system we consider a simplified model. Let the source be represented as a straight cylindrical magnetic tube immersed in coronal plasma. In Appendix A the eigenmode equation for fast and slow MA-oscillations is obtained (see Equation (A.8)). At first let us consider the case of $\rho_e = 0$ and $L \to \infty$, first analyzed by Rosenberg (1970) (new designations are given in the Appendix A). In this case Equation (A.8) describes the properties of the fast mode. For axial-symmetrical oscillations ($m = 0$) we get

$$\omega = \omega_0 = K_\perp (C_{\rm A}^2 + C_{\rm S}^2)^{1/2}, \tag{3.1}$$

where the set of eigennumbers $K_\perp = \lambda_j/R$ is defined by the equation (compare with Rosenberg's result)

$$J_0(\lambda_j) = 0\,. \tag{3.2}$$

Now let us elucidate the effects of an ambient plasma on the properties of MHD-oscillations. We consider the case $\rho_{0e}/\rho_{0i} \ll 1$. The fact that the plasma pressure is considerably lower than the magnetic pressure, $\beta_0 = 8\pi p_0/B_0^2 \ll 1$, allows us to treat the fast mode independently from other modes, and to reduce the equilibrium Equation (A.2) to the approximate condition

$$B_{0e} \simeq B_{0i}\,. \tag{3.3}$$

Under these assumptions Equation (A.8) for the main mode with $m = 0$ takes the form

$$\varkappa_i R \frac{J_0(\varkappa_i R)}{J_1(\varkappa_i R)} = \eta R \frac{H_0^{(1)}(\eta R)}{H_1^{(1)}(\eta R)}, \tag{3.4}$$

where

$$\varkappa_i^2 = \omega^2/C_{Ai}^2 - K_\parallel^2, \qquad \eta^2 = (\rho_{0e}/\rho_{0i})K_\perp^2 - K_\parallel^2\,.$$

Designating $\varkappa_i R = \lambda_j + \Delta$, where Δ takes account of $\rho_{0e}/\rho_{0i} \ll 1$ and $(K_\parallel/K_\perp)^2 \ll 1$, and expanding (3.4) in series we get

$$\omega = \omega_0 \times \begin{cases} 1 + \dfrac{K_\parallel^2}{2K_\perp^2} - \left(\dfrac{\rho_{0e}}{\rho_{0i}} - \dfrac{K_\parallel^2}{K_\perp^2}\right) \ln \dfrac{2}{\eta R} - i\dfrac{\pi}{2}\left(\dfrac{\rho_{0e}}{\rho_{0i}} - \dfrac{K_\parallel^2}{K_\perp^2}\right), & \eta^2 > 0\,, \\ 1 + \dfrac{K_\parallel^2}{2K_\perp^2}, \quad \eta^2 = 0\,, \\ 1 + \dfrac{K_\parallel^2}{2K_\perp^2} + \left(\dfrac{K_\parallel^2}{K_\perp^2} - \dfrac{\rho_{0e}}{\rho_{0i}}\right) \ln \dfrac{1}{|\eta| R}, & \eta^2 < 0\,. \end{cases} \tag{3.5}$$

Here ω_0 is defined by Equations (3.1) and (3.2); for $\beta_0 \ll 1$,

$$\omega_0 = \lambda_j C_{Ai}/R\,. \tag{3.6}$$

As follows from (3.5) when the external plasma density is finite, an imaginary part of the frequency arises corresponding to emission of the fast mode by the plasma cylinder. This result could be expected because the group velocity of the fast mode is directed almost across the cylinder axis when $K_\parallel^2 \ll K_\perp^2$. The wave energy emitted from the tube transforms into energy of motion of the external plasma. The external plasma also affects the elastic properties of the system. As a consequence, the real part of the frequency is also altered in comparison with the case $\rho_{0e} = 0$. Similar results were obtained by Zaitsev and Stepanov (1975a) having considered another relation between C_A and C_S.

The emissive damping rate of the fast mode

$$\nu = -\frac{\pi}{2}\omega_0\left(\frac{\rho_{0e}}{\rho_{0i}} - \frac{K_\parallel^2}{K_\perp^2}\right) \tag{3.7}$$

vanishes when $K_\parallel/K_\perp \geqslant (\rho_{0e}/\rho_{0i})^{1/2} \simeq C_{Ai}/C_{Ae}$. Then total internal reflection of a wave appears, i.e. the magnetic tube becomes an ideal resonator.

The analysis of Equation (A.8) shows (Zaitsev and Stepanov, 1975a) that the fast modes do not exist when outer and inner densities are comparable and $R \ll L$. This result is evident because such a system is not radially limited and free oscillations are known not to exist here. Thus the requirement of a large ratio ρ_i/ρ_e is essential.

Let us now estimate the quality $Q \simeq \omega_0/2|\nu|$ of a pulsating source, for example of the event described by Maxwell and Rinehart (1974). If the intensity of the type IV emission is a power function of B_0 (incoherent synchroton radiation of relativistic electrons with power energy spectrum) we have

$$Q \simeq \pi t/T, \tag{3.8}$$

where t is the time of e-fold decrease in the amplitude of radio emission pulsations, T is the pulse period. For $t = 200$ s and $T = 2$ s we have $Q \simeq 300$, i.e. without the fast mode amplification it is necessary that $\rho_{0e}/\rho_{0i} \simeq 3 \times 10^{-3}$ (at least without applying a finite $K_\parallel$). In the event described by Akin'yan *et al.* (1973), a still lower value of ρ_{0e}/ρ_{0i} is required.

Such a difference between the inner and outer plasma densities is unlikely to exist. As follows from these examples, an amplification of the fast mode should operate to provide the high pulsation quality observed.

3.3. The Bounce-Resonant Excitation of the Fast Mode by Energetic Particles

The amplification of the oscillations can be provided by the bounce-resonant interaction between the waves and high-energy particles (electrons, protons) trapped by a magnetic arch in the high corona. As the energy density of the fast particles is considerably lower than the energy density of a magnetic field, their contribution to the real part of the frequency is small (see, e.g. Southwood, 1976; Meerson *et al.*, 1977). These particles, however, may significantly contribute to the imaginary part of frequency and lead to instability (growth of the wave amplitude).

The importance of the bounce-motion effects can be seen from evaluation of (2.3): pulse periods observed are comparable with bounce-periods of fast particles. Consistent linear theory of bounce-resonant excitation of hydromagnetic waves was developed by Southwood (1973) and Karpman *et al.* (1977). Before applying a general theory to the situation considered we assume, for simplicity, that the stationary magnetic field B_0 depends on the longitudinal coordinate s as follows:

$$B_0(s) \simeq B_a(1 + s^2/L^2), \tag{3.9}$$

B_a being the magnetic field at the top of the arch. In this case the undisturbed motion of a trapped particle is described in drift approximation by the harmonic oscillator equation

$$\ddot{s}+\Omega^2 s=0, \tag{3.10}$$

where $\Omega=(2\mu B_a/m_j)^{1/2}/L$ is the bounce-frequency; $\mu=m_j v_\perp^2/2B(s)$ the magnetic moment of the particle ($\mu=$ const.); $v_\perp$ the velocity component transverse to the magnetic field; m_j the particle mass.

If the fast mode exists in a plasma, the force $F_\parallel=-\mu\nabla\tilde{B}_\parallel$ affects the particle, where $\tilde{B}_\parallel$ is the component of the magnetic field of the wave parallel to B_0. For a wave field of the form

$$\tilde{B}_\parallel(s,t)=B_{\parallel 0}\sin(\omega t-K_\parallel s) \tag{3.11}$$

the general expression by Karpman *et al.* (1977) for the wave growth rate takes the form

$$\gamma/\omega=8\pi^4 n_0\sin^2\theta\sum_{K=1}^{\infty}K\int \mathrm{d}I\,\mathrm{d}\mu\mu^2\frac{\partial f}{\partial I}J_K^2(K_\parallel a)\delta(\omega-K\Omega), \tag{3.12}$$

where n_0 is density of the fast particles at $s=0$; $I=m_j\Omega a^2/2$ is the longitudinal adiabatic invariant; a amplitude of particle bounce; θ angle between the wave vector and the stationary magnetic field, $J_K(w)$ – Bessel function, $\delta(w)$ – Dirac's δ-function. The distribution function $f(\mu, I)$, introduced in (3.12) is related to the velocity distribution function (normalized to density) in the following way:

$$f(\mu,I)=\frac{2B_a}{n_0 m_j^2 L}F(\mathbf{v},s). \tag{3.13}$$

Equation (3.12) shows, that contribution to the wave growth rate is due to particles satisfying the bounce-resonance condition

$$\Omega(\mu)=\omega/K, \qquad K=1,2,\ldots. \tag{3.14}$$

As follows from (3.12), the energetic particles with distribution function monotonically decreasing with I, $\partial f/\partial I<0$, cause the damping of fast mode. For the instability to arise the condition $\partial f/\partial I>0$ must be fulfilled in a some region (μ, I).

Suppose that the distribution function F is isotropic, i.e. it depends only on the total energy*

$$\varepsilon=(m_j/2)(v_\parallel^2+v_\perp^2)=\mu B_a+\Omega I. \tag{3.14}$$

We approximate this dependence as

$$F(\mathbf{v},s)\sim\varepsilon^p\exp\left[-\varepsilon(p+\tfrac{3}{2})/T\right], \qquad p>0, \tag{3.15}$$

where T is the mean energy of fast particles. Distribution function (3.15) describes a lack of low energy particles in the trap. Possibility of such a choice of F will be discussed later.

* Actually because of the loss-cone presence, some anisotropy takes place, but for a large 'mirror ratio' it can be neglected.

Under these assumptions the expression (3.12) for the growth rate takes the form

$$\gamma(\omega, K_\parallel) = \omega \frac{\pi^2(p+\frac{3}{2})}{\Gamma(p+\frac{3}{2})} \beta_h \Phi[(\omega/\Omega_T)^2, (K_\parallel L)^{-2}], \tag{3.16}$$

where

$$\beta_h = 8\pi n_0 T/B_a^2, \qquad \Omega_T^2 = 2T/m_j L^2(p+\tfrac{3}{2}),$$

$$\Phi(\lambda, \eta) = \sum_{K=1}^{\infty} \Phi_K(\lambda/K^2, \eta), \tag{3.17}$$

$$\Phi_K(x, y) = yx^{p+5/2} e^x \times$$
$$\times \left\{ \sum_0^{p-1} (pC_{p-1}^n - x(C_p^n)(-y)^n \psi_K^{(n)}(xy) - x(-y)^p \psi_K^{(p)} \right\}, \tag{3.18}$$

$\psi_K(x) = (\frac{1}{2}x)\exp(-\frac{1}{2}x) I_K(\frac{1}{2}x)$, $I_K(w)$ – modified Bessel function, $\Gamma(w)$ – gamma-function, $\psi_K^{(n)} = \mathrm{d}^n\psi_K/\mathrm{d}\omega^n$. It follows from (3.16) that the value $\gamma/\omega\beta_n$ depends only on dimensionless parameters, p, ω/Ω_T and $K_\parallel L$. The analysis of expression (3.16) (see Appendix B) allows to find the region of instability in the plane $(\omega/\Omega_T, K_\parallel L)$. The fact of importance is that the dependence of the growth rate upon the frequency has a maximum in the region of $\omega/\Omega_T \sim 1$. Thus, if the typical bounce-frequency of fast particles is comparable with the frequency of eigen-oscillations of the 'resonator', effective energy transformation from particles to oscillations can occur.

For not a large p the maximum growth rate is in the order of magnitude

$$\gamma_{\max}/\omega \simeq 0.1\beta_h, \tag{3.19}$$

when $K_\parallel L \simeq 2-3$.

Since $R/L \ll 1$, the 'optimal' wave vector is nearly orthogonal to the stationary magnetic field. In this case we can neglect the Čerenkov' damping of the fast mode caused by thermal coronal electrons (this neglect was done in derivation of Equation (3.16)).

It should be noted that the estimate (3.19) depends slightly on the specific form of the distribution function $F(\mathbf{v}, s)$; it is essential only for this function to have a sufficiently abrupt fall in the low energy region of a spectrum.

The lack of 'soft' particles in the coronal trap can be formed in principle due to the following reasons:

(a) the peculiarity of the accelerating mechanism – the preferential acceleration to given energy;

(b) lower life-time for 'soft' particles as a result of precipitation into the chromosphere due to the wave-turbulence and Coulomb scattering*;

* As it has been pointed out by Dr J. Kuijpers, Coulomb scattering alone may be not enough for maintaining of abrupt fall in the low energy region of a spectrum.

(c) dependence of the radial diffusion coefficient on the energy.

It follows from (3.19) that the growth rate determined by the component of the fast particles (electrons or protons) that provides higher β_h. What component plays the main role? The observational data are ambiguous. However, the spacecraft measuremenrts show, that for the case of proton flares, protons with energies $\gtrsim 10$ MeV leaving the active region, contain appeciable part of the flare energy (see, e.g., Somov and Syrovatskii, 1976) in contrast to energetic electrons. So it seems to us, that for the coronal traps at an altitude $\sim 1R_\odot$ the higher β_h can be most probably provided by the proton component. Our choice of the proton component as an energetic 'reservoir' for the MHD-waves is in agreement with the correlation between pulsations of type IV emission and 'protonability' of the flares. (McLean *et al.*, 1971). However, it's quite possible that electrons also should be taken into account.*

For successful interpretation of pulsations in the frame of the model involved, the important requirement is one of absence of more effective high frequency instabilities, that can 'destroy' the distribution function of type (3.15) and the bounce-resonant interaction. The isotropic distribution function with $\partial F/\partial v_\perp > 0$ can cause the ion-cyclotron instabilities (see, e.g., Mikhailovsky, 1974). The corresponding criteria of instability differ from the bounce-resonant instability one. These criteria are, as a rule, more rigid. Thus the bounce-resonant instability can be considered independently from others.

3.4. Condition of Amplification and Parameters of Pulsating Source

It follows from the Sections 3.2 and 3.3 that two contrary mechanisms regulate the amplitude of MHD-oscillations of a magnetic tube. The former is bounce-resonant excitation of the fast mode by energetic protons, the latter is the emission of the wave from the 'resonator'. If $\gamma_p \gtrsim \nu$, the MHD-oscillations are supported and the tube vibrates with the frequency $\omega \simeq C_A/R$. The self-excitation condition may be written as

$$0.1\beta_p \gtrsim \frac{\pi}{2}\left(\frac{\rho_e}{\rho_i} - \frac{K_\parallel^2}{K_\perp^2}\right) \quad \text{when} \quad K_\parallel^2 \leq K_\perp^2 \frac{\rho_e}{\rho_i}. \tag{3.20}$$

The wave numbers $K_\parallel$ and $K_\perp$ are determined by tube dimensions (see Appendix A): $K_\parallel \sim 3L^{-1}$, $K_\perp \sim 3R^{-1}$. So $K_\parallel^2/K_\perp^2 \sim R^2/L^2 \sim 10^{-4}$. It is inferred from (3.20) that when $\rho_e/\rho_i \sim 0.02$ and $\beta_p \gtrsim 0.2$, the self-excitation condition is fulfilled.

It is essential that in the model discussed an external source of excitation (for instance, a shock-wave) is not required. The possible role of the shock-wave in such a mechanism lies in additional injection and acceleration of energetic particles. After the passage of the shock-wave β_p may achieve the threshold value (3.20) (start of pulsation phase).

It should be noted that 'the bounce-resonant model' of pulsations is in agreement with the observed increase of a period of modulation with increase of the altitude of

* McLean and Rosenberg (1974) supposed, for example, that electron component can have $\beta_h \sim 1$.

the source of radio emission (Achong, 1974). The pulse period in our model, $\sim L/v_T$, grows with L.

At the end of this section we formulate the main properties of the described model of a pulsating source.*

(1) The amplification of MHD-waves can be provided by resonant interaction with energetic particles (first of all, the protons). Required proton pressure, $\beta_p \gtrsim 0.2$, the fast proton distribution function must have a fall in the low energy part of spectrum.

(2) It follows from the condition $\omega \sim \Omega$ (i.e. $C_A/R \sim v_T/L$) that at $L \simeq 10^{10}$ cm and $C_A \simeq 10^8 - 10^9$ cm, protons must be sufficiently energetic, $v_T \sim 10^{10}$ cm s^{-1}.

(3) The type IV radio emission is determined by the electron component of trapped energetic particles, because characteristic incremenrts of the amplification of high frequency waves by electrons is considerably higher than corresponding values for protons.

If electrons are relativistic (synchroton mechanism of radio emission), pulsations arise due to modulation of the synchroton radiation by the variable magnetic field of the fast MHD mode.

If electrons are nonrelativistic (plasma mechanism of radio emission), pulsations are possible as a result of adiabatic effects ($B \sim v_\perp^2$) of the MHD-wave on the electrons, emitting plasma waves. The similar mechanism was considered by Benz and Kuijpers (1976).

(4) The plasma density inside the magnetic tube must be at one-two orders higher, than the plasma density outside the tube. Pulsating sources must be rather compact: the characteristic length $\lesssim 10^{10}$ cm**, the radius $\simeq$(1–3) 10^8 cm. The fact, that observed sizes of pulsation sources of IV type emission considerably exceed the above values may be explained by low resolution of the radio telescopes in the metre range and by screening of the compact pulsating source by more extended continuum source (Caroubalos *et al.*, 1973; Kai and Takayanagi, 1973).

4. Conclusion

The analysis of pulsations of type IV radio emission indicates the variety of its properties. Apart from 'classical' pulsations (the period does not depend on the frequency of emission; smooth grow and decrease of amplitude (McLean *et al.*, 1971), the events of a sudden start of modulation (McLean and Sheridan, 1973) also takes place. The time profile of pulsations has sometimes an asymmetric shape with a slow rise and sudden decay (McLean and Sheridan, 1973). Sometimes the pulsations take the form of spikes, sometimes – the form of absorption bursts.

* The numbers given below have a character of mean estimates. Really for various events they can change significantly.

** The length of source can not exceed 10^{10} cm, because in the opposite case the 'erosion' of pulsations should be observed because the speed of light is finite.

Evidently such a variety of pulsation properties can not be explained in the frame of a single mechanism. Nevertheless, pulsations of 'classical' type are well explained in terms of modulation of radio emission by MHD-wave (fast mode), where damping due to radiation is compensated by bounce-resonant instability (interaction with energetic trapped protons). Despite the fact, that some of our reasonings are, in general, of an illustrative character, we managed to formulate the main requirements to pulsating sources of type IV radio emission (see (1)–(4) at the end of Section 3). The given example shows that the correct interpretation of fine structure of type IV solar radio emission requires an investigation of the role of low frequency plasma instabilities in the coronal magnetic traps.

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Appendix A: Eigenoscillations of a Magnetic Tube

Let us consider a magnetic tube (plasma cylinder with radius R) immersed in external plasma. Ends of the cylinder are frozen in superconducting planes spaced the distance $L \gg R$ perpendicularly to the axis $Z (Z \parallel \mathbf{B}_0)$. Suppose that the surface current flows along the boundary of cylinder. Consequently, when passing from the internal plasma to the external one, the plasma pressure and magnetic field have finite jumps. External and internal plasmas are considered as homogeneous and superconducting. Linearized MHD-equations take the form (see, e.g., Kadomtsev, 1966)

$$\begin{aligned}
&\rho_0 \frac{\partial^2 \boldsymbol{\xi}}{\partial t^2} = -\nabla\phi + \frac{1}{4\pi}[(\mathbf{B}_0 \nabla)\mathbf{b} + (\mathbf{b}\nabla)\mathbf{B}_0]\,,\\
&\rho = -\operatorname{div}(\rho_0 \boldsymbol{\xi})\,, \qquad\qquad (\text{A.1})\\
&p = -\boldsymbol{\xi}\nabla p_0 - \gamma p_0 \operatorname{div}\boldsymbol{\xi}\,,\\
&\mathbf{b} = \operatorname{rot}[\boldsymbol{\xi}, B_0]\,.
\end{aligned}$$

Here ρ_0, p_0, $\mathbf{B}_0$ are equilibrium values of pressure, density and magnetic fields, p, ρ, $\mathbf{b}$ and $\mathbf{v} = \partial\boldsymbol{\xi}/\partial t$ are small deviations from the equilibrium values, $\phi = p + (1/4\pi)\mathbf{B}_0 \cdot \mathbf{b}$. All disturbed quantities in (A.1) are sought as

$$\boldsymbol{\xi}(r)(a_1 \cos K_\parallel z + a_2 \sin K_\parallel z)(a_3 \cos m\varphi + a_4 \sin m\varphi) \operatorname{Re} e^{-i\omega t}\,,$$

where r, φ, z are the cylindrical coordinates, a_j are constants.

From (A.1) with allowance made for the equilibrium equation

$$\nabla(p_0 + B_0^2/8\pi) = 0\,, \qquad (\text{A.2})$$

we obtained the equation for the perturbating potential,

$$\frac{1}{r}\frac{\partial}{\partial r}\left[r\varepsilon_1(r)\frac{\partial\phi}{\partial r}\right]+\varepsilon_2(r)\phi=0\,, \tag{A.3}$$

where

$$\varepsilon_1(r)=[\rho_0(\omega^2-K_\parallel^2C_A^2)]^{-1},$$

$$\varepsilon_2(r)=[\rho_0(C_A^2+C_S^2)]^{-1}\left[1-\frac{m^2}{r^2}\frac{C_A^2+C_S^2}{\omega^2-K_\parallel^2C_A^2}-\right.$$

$$\left.-\frac{K_\parallel^2C_S^4}{\omega^2(C_A^2+C_S^2)-K_\parallel^2C_A^2C_S^2}\right].$$

The solution of Equation (A.3) must be bounded at $r=0$ and satisfy Sommerfeld's condition at $r\to\infty$. Then taking into account the condition of the disturbances going at zero at the ends of the plasma cylinder we obtain the perturbation potential inside ($r<R$) and outside the cylinder:

$$\phi_i(\mathbf{r},t)=\sum_{s=1}^{\infty}A_{sm}J_m(\varkappa_i r)\sin K_\parallel z\cos m\varphi\,\mathrm{Re}\,e^{-i\omega t},$$

$$\phi_e(\mathbf{r},t)=\sum_{\substack{s=1\\m=0}}^{\infty}C_{sm}H_m^{(1)}(\varkappa_e r)\sin K_\parallel z\cos m\varphi\,\mathrm{Re}\,e^{-i\omega t}. \tag{A.4}$$

Here $H_m^{(1)}(\varkappa r)$ is Hankel' function of the first kind, A_{sm} and C_{sm} are constants, depending on boundary conditions,

$$\varkappa^2=\omega^4[\omega^2(C_A^2+C_S^2)-K_\parallel^2C_A^2C_S^2]-K_\parallel^2\,. \tag{A.5}$$

At the boundary of the cylinder the potential must be continuous,

$$\phi_i(R)=\phi_e(R)\,, \tag{A.6}$$

because a jump of the potential leads to infinite acceleration of the plasma.

Second boundary condition can be obtained by integration of Equation (A.3) along the transition layer $(R-\sigma, R+\sigma)$ with subsequently letting $\sigma\to 0$:

$$\xi_{ri}(R)=\xi_{re}(R)\,, \tag{A.7}$$

where the radial displacement component ξ_r is connected with the perturbation potential by relation

$$\xi_r=\rho_0^{-1}(\omega^2-K_\parallel^2C_A^2)^{-1}(\partial\phi/\partial r)\,.$$

From Equations (A.4), (A.6) and (A.7) we obtain the uniform algebraic system of equations in coefficients A_{sm} and C_{sm}. The condition of existence of nontrivial solutions yields the eigenmode equation for a magnetic tube (Zaitsev and Stepanov, 1975a):

$$\frac{J_m'(\varkappa_i R)}{J_m(\varkappa_i R)}=\alpha\frac{H_m^{(1)\prime}(\varkappa_e R)}{H_m^{(1)}(\varkappa_e R)}\,, \tag{A.8}$$

where

$$\alpha = \frac{\varkappa_e}{\varkappa_i} \frac{\rho_{0i}}{\rho_{0e}} \frac{\omega^2 - K_\|^2 C_{Ai}^2}{\omega^2 - K_\|^2 C_{Ae}^2}.$$

The Equation (A.8) describes the properties of slow and fast magnetoacoustic modes (Alfvén waves $\omega = K_\| C_A$ do not contribute to potential ϕ).

Appendix B

Here we list the results of analysis of increment (3.16). For $p = 2$ stability and instability regions obtained numerically are illustrated in Figure 1. At the same place the levels of constant increment for $p = 2$ are shown. When $\omega/\Omega_T \sim 1$ one can take in Equation (3.17) only the main term $K = 1$. Then the cases $K_\| L \ll 1$ and $K_\| L \gg 1$ can be analyzed analytically. Tne maximum growth rate is achieved, as may be seen from Figure 1, when $K_\| L \simeq 3$ and $\omega/\Omega_T \simeq 0.5$. The limit $\omega/\Omega_T \gg 1$ corresponds to Landau damping of the wave (Karpman *et al.*, 1977).

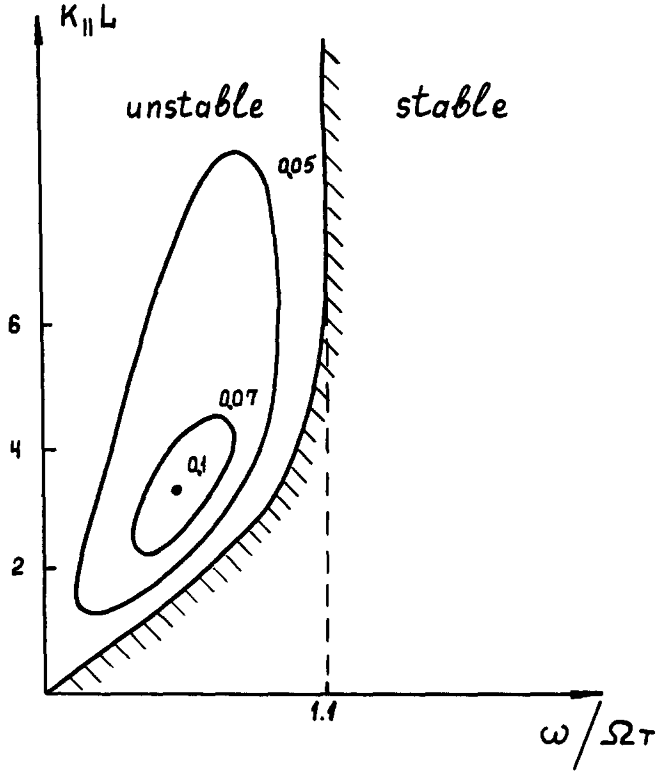

Fig. 1. Unstable and stable regions and the lines of constant increment in the plane $(\omega/\Omega_T, K_\| L)$ for $p = 2$.

Approximate dependence of maximum increment on parameter p is shown by Figure 2. The case $p \to \infty$ can be easily analyzed analytically, because then the distribution function (3.15) behaves as $\delta(\varepsilon - T)$ that permits simple integration over energies.

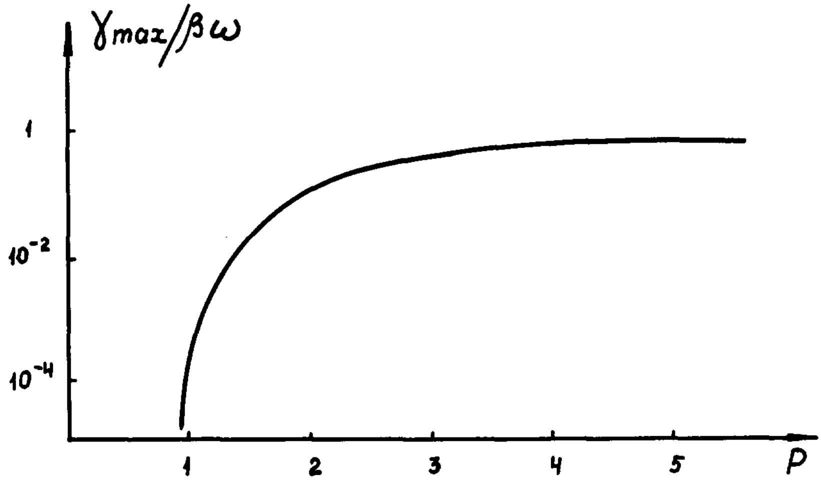

Fig. 2. The plot of dependence of maximum increment on parameter p (see Equation (3.15)).

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