PHYSICAL AND MATHEMATICAL MODELLING OF DYNAMICS OF IONOSPHERIC WAVE PRECURSORS OF EARTHQUAKES

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Abstract

In the present article a new standpoint is proposed on the theoretical modelling of transference of interaction from underground and overground perturbation sources to the above lying layers of atmosphere.

The new nonlinear mechanism of transformation of long-wave-length acoustic waves (of natural or artificial origin) into electromagnetic ones in ionospheric F-layer is suggested. It is shown that the acoustic wave involves in collective motion the charged particles of media by means of collisions. The relative motion of charged particles excites the alternating current with an arbitrary phase and, consequently, there happens the parametric generation of electromagnetic fields in the ionospheric layers.

Nonlinear propagation of the lowfrequency seismic origin acoustic-gravity (AG) perturbations in a nonuniform ionospheric E-layer has been investigated. Analytically calculated amplification of the night-sky green line intensity stipulated by nonlinear AG vortexes spreading has been compared with observed green radiation intensity increase of the night sky before the earthquake. Well matching of these data suggests that ionospheric AG vortexes can be considered as the wave forerunner of the strong earthquakes.

The results of investigation indicate the leading role of the acoustic channel of connection between lithosphere-ionosphere-magnetosphere.

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1. *Introduction*

More than hundred destroying earthquakes take place in the world every year. Untold human victims, physical and moral traumas, enormous material damage are connected with them. The danger from this point of view is gradually increasing because many densely populated regions of the globe today are situated in immediate proximity to the centers of tectonic tension. In addition a "suddenness" of earthquake is remaining dangerous by its social, economic and ecological consequences. Diminution of damage may be provided with division into districts, prognosis of place, time and force of earthquake. In prospect the prevention of earthquake's destructive consequences may be considered by engineering methods as the total direction of the prediction. Therefore the problem of prediction and prognosis of an earthquake remains one of the actual problems of the contemporary mankind.

An earthquake, that is the intrinsic destruction of the Earth's crust, in fact is not such a sudden event! The results of numerous observed data show that perturbations of various geophysical fields of deformation and electromagnetic nature and changes in seismic regime as well precede an earthquake. Today it is well known that before an earthquake different kinds of physical, geological, geochemical and biological phenomena are observed. From physical phenomena preceding earthquake should be noted: origin of intensive variations of geomagnetic field, electromagnetic radiation, night airglow and acoustic-gravitational waves (AGW). The significance of precursors is difficult to overrate. They play the decisive role in understanding processes of preparation and passing of earthquake. It is assumed today that the problem of long-term prediction is solved in principle. As to the problem of operating (short-term) prediction of an earthquake it remains complicated and complex and is still far from completion.

Search of precursors of an earthquake is conducted not only in solid envelope of the Earth (lithosphere), but also in air (atmosphere) and plasma (ionosphere) ones. Up to the present ionospheric perturbations were considered on the whole in connection with the solar action on it. However the experimental investigations of the last decade have shown that atmosphere - ionosphere layers are sensitive indicator of processes in troposphere, hydrosphere and lithosphere.

From this point of view the atmosphere - ionosphere envelope of the Earth is the unique medium. Preparing in it, nonordinary phenomena are appearing in the form of wave precursors over and under the Earth, while anthropogenic actions themselves are appearing in the form of wave perturbations.

So that the Nature itself informs in advance about coming nonordinary phenomena, it is necessary only to learn operatively and purposefully to make use of these information with maximum advantage for problems of predictions and managements.

From aforementioned it follows that for improvement of operating forecast of strong earthquakes it is necessary in the first place to manifest the class of evidently expressed physical processes preceding an earthquake. This is possible on the basis of analysis of experimental data of already performed earthquakes and anthropogenic actions.

The interest in the study of wave phenomena connected with the seismic and anthropogeneous activity has recently increased. Generation of acoustic waves in the atmosphere and ionospheric plasma at the very moment of an earthquake and immediately before it has been observed more than once [8,13,20,29]. An increase of intensity of electromagnetic disturbances in the ionosphere, at the altitude of the order of 150 km, on the eve of an earthquake has been recorded [7,12,15,18,26,30]. The analysis of the night airglow (of optical range) emission data also shows an increase of the emission line intensity before an earthquake [11,23,25,32], which also indicates the presence of acoustic and electromagnetic disturbances.

The acoustic waves in question may come from any natural and artificial sources, including earthquakes; when propagating upwards, they reach the ionospheric layers [16,28]. On satellite "Oreol-3", for instance, the effective energy transformation of anthropogeneous explosions - the acoustic perturbation of infrasonic band - into the energy of electromagnetic perturbations is detected [5]. Such an approach, however, excludes from consideration the ion motion with respect to the neutrons, which may turn crucial for studying electromagnetic disturbance generation. The approximation considering all the three components of the media, like that in [6,19,20], appears to be more general.

We hold that the most important drawback of all the previous works is the fact that they neglected the nonlinear term $(\vec{v}\nabla)\vec{v}$ in the velocity equation. In doing so they neglected the only available possibility of trasformation of the waves in a quasihomogeneous medium: nonlinear interaction of waves. Therefore the authors had to assume the existence of a sharp border line between a neutral gas and a fully ionized plasma, which is not the case, and so to reduce the problem, in fact, to a slightly modified mechanism of the linear transformation of waves on inhomogeneities of the medium. As a result, the coefficient of transformation for waves in the ionosphere has been found to be unimportant [22]. It is clear that for a large amplitude of a sound wave, the contribution of the nonlinear term when treating the problem of transformation of an acoustic wave into electromagnetic one, becomes crucial. In an ordinary, completely ionized plasma, when the amplitudes of all interacting waves are small, the nonlinear transformation is described by a decay instability. But if an initial wave ("pump-wave") is of a large amplitude, then the very formulation of the problem essentially changes for it is treated as that for a parametric instability [31] and so does the evaluation technique.

The nature of our case essentially differs from the ordinary parametric instability: the pump-wave is an acoustic one, i.e. oscillations of neutral particles. So a direct interaction of the pump-wave with electromagnetic modes is impossible. An intensive acoustic wave reaching the ionospheric altitudes, draws charged particles into the collective motion by the friction with the neutrons. It is important that the coefficients of frictions of the electrons and the ions with the neutral particles differ from each other, which causes a relative motion of the charged particles. The alternating currents induced by this motion can generate electromagnetic fields due to the parametric pumping. It is also important that such alternating currents have a chaotic phase that must be taken into account [4]. The specific nature of the case in question requires a new technique to be developed, somewhat different from that for the ordinary parametric instability [31], which we are going to propose below.

Supposing the ionospheric medium to be 3-component and quasihomogeneous, and neglecting the inertia of the electrons (which is true for Alfven and magneto-acoustic disturbances), the dynamics of wave disturbances is described by the equations [9]:

$$0=-\nabla P_e-en_e\left(\vec{E}+\frac{1}{c}[\vec{v}_e\vec{B}]\right)+m_en_e\nu_{ei}(\vec{v}_i-\vec{v}_e)+m_en_e\nu_{en}(\vec{v}_n-\vec{v}_e), \quad (1.1)$$

$$Mn_i\left\{\frac{\partial\vec{v}_i}{\partial t}+(\vec{v}_i\dot{\nabla})\vec{v}_i\right\}=-\nabla P_i+en_e\left(\vec{E}+\frac{1}{c}[\vec{v}_i\vec{B}]\right)+m_en_i\nu_{ei}(\vec{v}_e-\vec{v}_i)+$$
$$+Mn_i\nu_{in}(\vec{v}_n-\vec{v}_i), \quad (1.2)$$

$$Mn_n\left\{\frac{\partial\vec{v}_n}{\partial t}+(\vec{v}_n\dot{\nabla})\vec{v}_n\right\}=-\nabla P_n-m_en_e\nu_{en}(\vec{v}_n-\vec{v}_e)-$$
$$-Mn_i\nu_{in}(\vec{v}_n-\vec{v}_i), \quad (1.3)$$

$$\frac{\partial n_\alpha}{\partial t}+div(n_\alpha\vec{v}_\alpha)=0, \quad (1.4)$$

$$\nabla P_\alpha=\gamma_\alpha T_\alpha\nabla n_\alpha, \quad \alpha=e,i,n. \quad (1.5)$$

In eqs (1.1)-(1.5) $P_\alpha=P_{0\alpha}+\tilde{P}_\alpha, n_\alpha=n_{0\alpha}+\tilde{n}_\alpha$, the pressure and the density of ionospheric plasma, $\vec{B}=\vec{B}_0+\tilde{\vec{B}}$ and $\vec{E}$ are magnetic and electric fields (the index "0" means an equilibrium value, and a tilde above average quantities means their disturbed values), $\nu_{\alpha\beta}$ is an average frequency of collisions for a particle of the kind "α" with those of the kind "β",c is the speed of light, n_iv_{in} is the mass of an ion or a molecule (neutral), e and m_e are the charge and the mass of the electron, $\gamma_\alpha\sim 1$ is the ratio of specific heats, T_α is the temperature. In the sequel we shall omit, as a rule, the tilde for disturbed quantities.

Neglecting the displacement of the currents, we describe the disturbances in magnetic and electric fields by Maxwell equations:

$$
\begin{aligned}
&rot\vec{B} = \frac{4\pi}{c}\vec{j}, \quad rot\vec{E} = -\frac{1}{c}\frac{\partial \vec{B}}{\partial t},\\
&div\vec{E} = 4\pi e(n_e - n_i), \;\; div\vec{B} = 0, \;\; \vec{j} = e(n_i\vec{v}_i - n_e\vec{v}_e).
\end{aligned} \tag{1.6}
$$

Eqs. (1.1)-(1.6) describe the interaction of a gas disturbance with magnetized ionospheric plasma.

Considering that the ionospheric plasma may be treated, accurately enough, as quasi-neutral, $n_{0e} \approx n_{0i} = n_0$ and summing up eqs. (1.1),(1.2), we get the equation of motion for charged particles. One can see that the electric field has been removed from the equation

$$
Mn_i\left\{\frac{\partial \vec{v}_i}{\partial t} + (\vec{v}_i\nabla)\vec{v}_i\right\} = -\nabla(P_e + P_i) + \frac{1}{4\pi}[rot\vec{B}, \vec{B}] + Mn_i\nu_{in}(\vec{v}_n - \vec{v}_i). \tag{1.7}
$$

Here the following inequalities were taken into account:

$$
\frac{\nu_{en}}{\nu_{in}} \sim \frac{M\nu_{in}}{m_e\nu_{en}} \sim \left(\frac{M}{m_e}\right)^{\frac{1}{2}} \gg 1. \tag{1.8}
$$

Let us use the generalized Ohm's law for the transformation of the Faraday's equation of induction (1.6). The generalized Ohm's law can be obtained from the equations (1.1) and (1.2) (taking into account the inertia of electrons) (compare [14]):

$$
\frac{\partial \vec{j}}{\partial t} + \nu_e\vec{j} + \frac{\omega_{ci}}{B_0}[\vec{j}, \vec{B}] = -\frac{e^2 n_e}{m}(\vec{E} + \frac{1}{c}[\vec{v}_i, \vec{B}] - \frac{e}{m}\nabla P_e + en_e\nu_{en}(\vec{v}_n - \vec{v}_i), \tag{1.9}
$$

where $\nu_e = \nu_{ei} + \nu_{en}, \omega_{c\alpha} = eB_0/m_\alpha c, \alpha = i, e$. During the selfconsistent hydrodinamical description it is necessary to neglect the term $\partial\vec{j}/\partial t$ in the left hand side, because $\partial/\partial t \sim \omega \ll \nu_e$ (where ω is the frequency of the considered waves). Besides, if we take into account that $\vec{v}_n \approx \vec{v}_i$ for the ionospheric perturbations and in particular for the processes which we discuss, from (1.9) one can obtain

$$
\vec{E} = -\frac{1}{c}[\vec{v}_i, \vec{B}] + \frac{m}{e^2 n_e}\nu_e\vec{j} - \frac{\nabla P_e}{en_e} - \frac{m}{eMcn_e}[\vec{j}, \vec{B}]. \tag{1.10}
$$

Estimating each term in the right hand side of (1.10) using (1.14), one can find that in case when the collision is scarce enough ($\omega_{ce} \gg (M/m)^{1/2}\nu_e$, which is fulfilled with great stock in the ionospheric layers starting from the height $h \geq 150$ km) and for slow processes ($\omega \ll \omega_{ci}, \omega \approx$

kv_i, here k is the characteristic wave number of the generated disturbances), in the right hand side of (1.10) the most important is the first term. Therefore Faraday's induction law will take the following form:

$$\frac{\partial \vec{B}}{\partial t} = rot[\vec{v}_i, \vec{B}]. \tag{1.11}$$

(1.11) means that the magnetic field is frozen in the ion liquid, which is peculiar for the magnetosonic perturbations discussed by us [14,24].

In the equation (1.7) the term proportional to ∇P_e is decisive, because it describes the motion of electrons relative to ions. To define $\nabla P_e (= \gamma_e T_e \nabla n_e)$ correctly, one must have the equation connecting velocities of electrons and ions. In our case (ionospheric medium) this connection is achieved by collisions of electrons and ions with neutral particles. To find this connection we will use the schema of calculation described in [3].

Following [3], let us consider the motion of particles of α sort colliding with neutral particles in the average $\nu_{\alpha n}$ times per second. After collisions the particle velocity is distributed chaotically, or is equal to zero with respect to the neutral particles. It is assumed that neutrals are immovable. Between the collisions the motion of the particle with the charge q_α undergoes only the influence of electric and magnetic fields. The electric field is the alternating field $\vec{E}$ generated by the acoustic wave. Let us assume that the frequency of the change of the generated field ω is much less than the cyclotron frequency, $\omega \ll \omega_{c\alpha}$. We can therefore believe that the electric field remains constant during the Larmour rotation period, and the magnetic field of perturbation is small compared with the external magnetic field. Let the homogeneous geomagnetic field $\vec{B}_0$ be directed along z axis, and the electric field have components along x and z axes, $\vec{E}(E_x, 0, E_z)$. In such a model one can easily find the solution of the equation of motion of charged particles. For the particle, motionless at the moment $t = 0$, we have well known solutions:

$$\tilde{v}_{\alpha x} = \frac{q_\alpha}{\mid q_\alpha \mid} \frac{cE_x}{B_0} \sin \omega_{c\alpha} t, \quad \tilde{v}_{\alpha y} = -\frac{cE_x}{B_0}(1 - \cos \omega_{c\alpha} t), \quad \tilde{v}_{\alpha z} = q_\alpha \frac{E_z}{m_\alpha} t. \tag{1.12}$$

The interval between collisions obeys the laws of the statistical distribution, so the probability that the collision happens between t and $t + dt$ is $f_{\alpha n} = \nu_{\alpha n} exp(-\nu_{\alpha n} t) dt$. Then multiplying each of the velocity components (1.12) by $f_{\alpha n}$ and integrating over time we find the components of mean or drift velocity of charged particles:

$$v_{\alpha x} = \frac{q_\alpha}{\mid q_\alpha \mid} \frac{\nu_{\alpha n} \omega_{c\alpha}}{\nu_{\alpha n}^2 + \omega_{c\alpha}^2} \frac{cE_x}{B_0}, \quad v_{\alpha y} = \frac{\omega_{c\alpha}^2}{\nu_{\alpha n}^2 + \omega_{c\alpha}^2} \frac{cE_x}{B_0}, \quad v_{\alpha z} = \frac{q_\alpha}{m_\alpha \nu_{\alpha n}} E_z. \tag{1.13}$$

Rewriting these expressions separately for electrons and ions ($\alpha = e, i$) and excluding $\vec{E}$ and $\vec{B}_0$ one can easily find the connection between velocity components $(v_e)_j$ and $(v_i)_j$ (j=x,y,z or 1,2,3):

$$(v_e)_j = \beta_{jj}(v_i)_j. \tag{1.14}$$

Here

$$\beta_{11} = -\frac{\nu_{en}\omega_{ce}}{\nu_{in}\omega_{ci}}\frac{\nu_{in}^2 + \omega_{ci}^2}{\nu_{en}^2 + \omega_{ce}^2}, \quad \beta_{22} = \frac{\omega_{ce}^2}{\omega_{ci}^2}\frac{\nu_{in}^2 + \omega_{ci}^2}{\nu_{en}^2 + \omega_{ce}^2}, \quad \beta_{33} = -\frac{m_i}{m_e}\frac{\nu_{in}}{\nu_{en}}. \tag{1.15}$$

The self-consistent set of equations (1.3)-(1.5),(1.7),(1.11),(1.14) describes generation of electromagnetic field when an acoustic wave propagates in the ionosphere.

1.1. *Dispersion equation for coupled acoustic and electromagnetic modes*

Hereafter we will consider the problem in the frame of the oscillating neutrals, in order to take into account the influence of the acoustic wave on the charged particles.

At the hydrodynamic description the identity of the medium particles behaviour is usually assumed. For instance, in two fluid hydrodynamics, at the parametric generation of waves in plasma by an external electromagnetic field, the pump wave affects all charged particles simultaneously. Because of the difference in inertia of the plasma components, electrons move with respect to ions causing thereby the parametric instability.

In our case the influence of the acoustic pump wave on the charged particles is achieved only by collisions. It is obvious that collisions with different particles happen at the different moments of time. After each case of collision of neutrals with charged particles their momenta change - as if the particles begin to move with a new phase. Hence for the adequate description of the statistical state of many particles one has to introduce the arbitrary phase. In the case of two fluid hydrodynamics one has to introduce the arbitrary phases φ and ψ for the plasma electrons and ions respectively, with the consequent averaging over these phases.

Let us consider the case when the frequency of generated perturbation ω is much more than the frequency of collisions of neutral particles with ions $\nu_{ni}, \omega \gg \nu_{ni}(\nu_{ni} \sim (n_{0i}/n_{0n})\nu_{in} \ll \nu_{in})$. In this frequency range the motion of the charged particles doesn't influence the motion of neutrals. So the solution of the equation of the neutral particle dynamics (1.3),(1.4) can be accomplished independently, and the velocity $\vec{v}_n$ can be presented in the form of oscillations with an arbitrary phase ψ,

$$\vec{v}_n = \vec{V}_0 \cos(\omega_0 t + \psi). \tag{1.16}$$

Here $\vec{V}_0$ is a time-invariant amplitude of neutral particle velocity, ω_0 is the acoustic (pump) wave frequency, ψ is the initial phase of neutrals or ions. We describe an enormous set of particles. So it is clear that ψ is a random value.

Therefore the quasi-homogeneous atmosphere is supplied by the external constant source of perturbation (1.16) (for instance, by the oscillation of the terrestrial surface which, because of superficial Rayleigh wave excitation, generate the sound propagating up to ionospheric heights).

The form of the external source (1.16) presumes that $\omega_0 t \gg \vec{k}_0\vec{r}$, where $\vec{k}_0$ ($\parallel \vec{V}_0$) is the wave vector of pump wave, i.e. the area of the generation of ionospheric perturbation is less than the pump wave length $\lambda_0 = 2\pi/k_0$ (see also section 2.3). This consideration distinguishes substantially our way of setting up the problem from the previously considered analogous ionospheric problems [5,6,19,22]. In the laboratory of plasma the formulation of the problem in this way is popular and well grounded [31].

Let us investigate temporal evolution of the ionospheric perturbation created by the external source (1.16) in the frequency of $\nu_{ni} \ll \omega \ll \omega_{ci}$. As it was mentioned above, one should pass to the frame of neutrals' rest, and for convenience, let us introduce new variables

$$\vec{V} = \vec{v}_i - \vec{v}_n = \vec{U}(t)e^{-a\sin(\omega_0 t + \psi)}, \quad a = \frac{\vec{k}\vec{V}_0}{\omega_0}, \tag{1.17}$$

where $\vec{k}$ is a wave vector of the generated perturbation. For these new variables we get from eqs.(1.7),(1.11), with an accuracy of the terms of the second order of smallness:

$$\frac{\partial \vec{U}(t)}{\partial t} - \nu_{in}\vec{U}(t) = \Bigg\{\omega_0\vec{V}_0\sin(\omega_0 t + \psi) - \frac{\nabla(P_e + P_i)}{Mn_0} -$$

$$-\frac{1}{4\pi M n_0}\int^t \left[\vec{B}_0[\nabla[\nabla[\vec{U}(t'), \vec{B}_0]]]\right] e^{-ia\sin(\omega_0 t' + \psi)}\,dt'\Bigg\} e^{ia\sin(\omega_0 t + \psi)}. \tag{1.18}$$

It should be noted that the ions and the electrons interact with the neutrals not in a synchronous way. So the initial phases of oscillations for the electrons φ and the ions ψ are not the same. Taking this (as well as (1.14)) into account, the solution of the continuity equation for electrons and ions (1.4) could be presented in the form:

$$n_e(t) = -n_0 e^{-ib\sin(\omega_0 t + \varphi)} \sum_{n,s=-\infty}^{\infty} J_n(b) J_s(a) e^{i(n\varphi - s\psi)}$$

$$\times \int_{-\infty}^{\infty} \frac{k_j \beta_{jj} U_j(\omega') e^{i[\omega' + (n-s)\omega_0]t}}{\omega' + (n-s)\omega_0} d\omega', \tag{1.19}$$

$$n_i(t) = -in_0 e^{-ia\sin(\omega_0 t + \psi)} \int^{t} \vec{k}\vec{U}(t')\, dt', \tag{1.20}$$

where

$$b = \frac{\beta_{jj} k_j V_{0j}}{\omega_0}, \quad j = x, y, z(1,2,3), \tag{1.21}$$

Summing up is over the index j. The expansion

$$e^{ia\sin(\omega_0 t + \psi)} = \sum_{n=-\infty}^{\infty} J_n(a) e^{in(\omega_0 t + \psi)},$$

was used in eq. (1.19). Here $\vec{U}(\omega)$ is the Fourier component of $\vec{U}(t)$.

We substitute the expressions (1.19) and (1.20) into (1.18). Then we integrate (1.18) over t' to Fourier transformation and pass back from the variable $\vec{U}(\omega)$ to $\vec{V}(\omega)$ using the expression

$$\vec{U}(\omega) = \sum_{l=-\infty}^{\infty} J_l(a) \vec{V}(\omega - l\omega_0) e^{il\psi},$$

that is the Fourier transform of eq. (1.17). If here the phase ψ is fixed as it takes place in the case of usual parameters [31], the infinite sum of various harmonics is obtained. To describe such a system one has to cut off the infinite system of the overlapping equation, using the additional resonance conditions. As for our case, at the averaging over arbitrary phases φ and ψ, the system of equations for $\vec{V}$ is automatically simplified. Indeed, as the result of such an averaging over arbitrary phases equation (1.18) eventually takes the form:

$$\omega(1 + i\frac{\nu_{in}}{\omega})\vec{V}(\omega) - \vec{k}_\perp(\vec{k}_\perp \vec{V}_\perp(\omega))\frac{V_A^2}{\omega} - \vec{k}(\vec{k}\vec{V}(\omega))\frac{V_{Ti}^2}{\omega} -$$

$$-\vec{k}\beta_{jj}(k_j V_j(\omega)) V_{se}^2 \sum_{n=-\infty}^{\infty} \frac{J_n^2(b)}{\omega - n\omega_0} \equiv A_{ij} V_j(\omega) = 0. \tag{1.22}$$

Here $V_A^2 = B_0^2/(4\pi n_0 M), V_{Ti}^2 = \gamma_i T_i/M, V_{se}^2 = \gamma_e T_e/M; k = (k_\perp^2 + k_\parallel^2)^{\frac{1}{2}}; \vec{k}_\perp, \vec{k}_\parallel$ are the orthogonal and parallel (to the outer magnetic field $\vec{B}_0$) components of the wave vector, respectively.

Vector equation (1.22) or the corresponding system of three scalar equations defines the relation between the acoustic and electromagnetic waves in the ionosphere.

The non-trivial solution of the equation (1.22) exists if the determinant formed by the coefficients of the system is zero. This expression is the desired dispersion equation:

$$\begin{aligned} \parallel A_{ij} \parallel = A_{11}A_{22}A_{33} + A_{12}A_{23}A_{31} + A_{13}A_{21}A_{32} - \\ -A_{13}A_{22}A_{31} - A_{11}A_{23}A_{32} - A_{12}A_{21}A_{33} = 0, \end{aligned} \tag{1.23}$$

where

$$\begin{aligned}
&A_{11} = 1 + i\frac{\nu_{in}}{\omega} - \frac{k_x^2}{\omega^2}(V_A^2 + V_{Ti}^2 + \beta_{11}A_0V_{se}^2),\\
&A_{12} = -\frac{k_xk_y}{\omega^2}(V_A^2 + V_{Ti}^2 + \beta_{22}A_0V_{se}^2),\\
&A_{13} = -\frac{k_xk_z}{\omega^2}(V_{Ti}^2 + \beta_{33}A_0V_{se}^2), \quad A_{21} = -\frac{k_yk_x}{\omega^2}(V_A^2 + V_{Ti}^2 + \beta_{11}A_0V_{se}^2),\\
&A_{22} = 1 + i\frac{\nu_{in}}{\omega} - \frac{k_y^2}{\omega^2}(V_A^2 + V_{Ti}^2 + \beta_{22}A_0V_{se}^2),\\
&A_{23} = -\frac{k_yk_z}{\omega^2}(V_{Ti}^2 + \beta_{33}A_0V_{se}^2),\\
&A_{31} = -\frac{k_zk_x}{\omega^2}(V_{Ti}^2 + \beta_{11}A_0V_{se}^2), \quad A_{32} = -\frac{k_zk_y}{\omega^2}(V_{Ti}^2 + \beta_{22}A_0V_{se}^2),\\
&A_{33} = 1 + i\frac{\nu_{in}}{\omega} - \frac{k_z^2}{\omega^2}(V_{Ti}^2 + \beta_{33}A_0V_{se}^2), \quad A_0 = \sum_{n=-\infty}^{\infty}\frac{\omega J_n^2(b)}{\omega - n\omega_0},\\
&b = \beta_{jj}\frac{k_jV_{0j}}{\omega_0},
\end{aligned} \tag{1.24}$$

which determines the spectra of generated perturbations $\omega = \omega(\vec{k})$.

We note that in (1.24) the arguments of Bessel's functions comprise the velocity of neutral particles $V_0 (b = \beta_{jj}k_jV_{0j}/\omega_0)$. Thanks to this (1.23) defines the nonlinear dependence of the spectra of generated perturbations of the amplitude of pumping wave. In this sense the obtained dispersion equation describes the nonlinear effect of impact of the acoustic signal (pumping) on ionospheric plasma.

1.2. *Generation of electromagnetic noises on the harmonics of acoustic frequencies*

It is well-known from the theory of parametric instabilities that in a conducting medium, in an outer electric field, when the frequency of the outer field ω_0 is close to one of the selfoscillation frequencies of the medium $\omega(\vec{k})$, increasing disturbances may possibly appear.

In our case the motion of neutral particles (acoustic pumping wave) in the ionosphere causes separation of charges through collisions and, subsequently, induction of an electric field. The induced field can change the parameters of ground state of ionospheric medium and lead to parametric instabilities, i.e. to the generation of electromagnetic field. To demonstrate this, let us consider the solution of (1.23) in extreme cases - transversal and longitudinal propagation of generated perturbations.

1. Let the perturbation propagate transversely to the geomagnetic field, $\vec{k} \perp \vec{B}_0$.

Assuming $k_z = 0$ in (1.24), from (1.23) one can obtain:

$$\omega^2 + i\nu_{in}\omega - k_\perp^2(V_A^2 + V_{Ti}^2) - (\beta_{22}k_y^2 - |\beta_{11}|k_x^2)V_{se}^2 \sum_{n=-\infty}^{\infty} \frac{\omega J_n^2(b)}{\omega - n\omega_0} = 0. \quad (1.25)$$

Now let us consider the frequencies of medium oscillations, close to the pumping wave harmonics, i.e. sonic wave,

$$\omega(\vec{k}) = \omega_k + i\gamma, \quad (1.26)$$

were $\omega_k \simeq l\omega_0 (l = 1, 2, ...)$ and $\gamma \ll \omega_k$. Then, substituting (1.26) by (1.25) and expanding over the degrees of γ/ω_k, we will obtain for the oscillation spectrum:

$$\omega_k^2 = k_\perp^2(V_A^2 + V_{Ti}^2) + (\beta_{22}k_y^2 - |\beta_{11}|k_x^2)V_{se}^2 J_l^2(b). \quad (1.27)$$

Here, the effect of the outer acoustic (pump) wave manifests itself in the correction terms and specifies the dependence of the frequency on the wave vector (because $b = \beta_{jj}k_j V_{0j}/\omega_0$). For the wave lengths greatly exceeding the characteristic pump-wave length $V_0/\omega_0 (b \ll 1)$, the oscillation frequency (1.27) equals the frequency of a fast magneto-acoustic wave.

The growth rate of the oscillations (1.27) is equal to

$$\gamma = -\frac{\nu_{in}}{2} + \left[\frac{\nu_{in}^2}{4} - \frac{1}{2}(\beta_{22}k_y^2 - |\beta_{11}|k_x^2)V_{se}^2 J_l^2(b)\right]^{\frac{1}{2}}. \quad (1.28)$$

It follows that the generation of electromagnetic perturbations ($\gamma > 0$) is possible in ionospheric plasma caused by external acoustic waves, in the area of the following wave numbers:

$$k_x > \left(\frac{\beta_{22}}{|\beta_{11}|}\right)^{\frac{1}{2}} k_y \sim \left(\frac{M}{m_e}\right)^{\frac{1}{4}} k_y. \quad (1.29)$$

Accordingly from (1.27) one can obtain the relation for the threshold values of the pumping wave amplitude where the instability becomes possible:

$$J_l^2\left(|\beta_{11}|\frac{k_x V_{0,thr}}{\omega_0}\right) \leq \frac{V_A^2 + V_{Ti}^2}{V_{se}^2}. \quad (1.30)$$

When $V_0 \to 0$, the expression (1.28) gives the known result $\gamma = -\nu_{in}/2$ and we have the strong damping of electromagnetic disturbances in the collisional medium.

Far from the resonance of the external frequency harmonics $|\omega_k - n\omega_0| \gg \omega_k$, i.e. for great $n = N \gg 1$ one can obtain for the frequency and the growth rate the following expressions:

$$\omega_{k\pm} = -(\beta_{22}k_y^2 - |\beta_{11}|k_x^2)V_{se}^2\frac{\varepsilon_0}{2\omega_0} \pm \left[(\beta_{22}k_y^2 - |\beta_{11}|k_x^2)^2 V_{se}^4\frac{\varepsilon_0^2}{4\omega_0^2} + k_\perp^2(V_A^2 + V_{Ti}^2)\right]^{\frac{1}{2}}, \tag{1.31}$$

$$\gamma = -\nu_{in}\left[2 + (\beta_{22}k_y^2 - |\beta_{11}|k_x^2)V_{se}^2\frac{\varepsilon_0}{\omega_0\omega_{k\pm}}\right]^{-1}, \tag{1.32}$$

where $\varepsilon_0 = \sum\limits_{N=n\gg 1}^{\infty}\frac{1}{n^2} \le \frac{\pi^2}{6}$. In the area of wave numbers $k_x > (M/m_e)^{\frac{1}{4}}k_y$ the fast wave ω_{k+} is excited, and in the area $k_x < (M/m_e)^{\frac{1}{4}}k_y$ -the wave ω_{k-} is. Besides, for the generated frequencies the following conditions must be satisfied:

$$|\omega_{k\pm}| < \varepsilon_0|\beta_{11,22}|\frac{k_{x,y}^2V_{se}^2}{2\omega_0},$$

and for the threshold value of the pump field one has:

$$\beta_{jj}k_jV_{0,thr}/\omega_0 \approx N \gg 1, \quad (j = 1,2 \ \ or \ \ x,y).$$

2. Let us consider the case of the longitudinal propagation $\vec{k} \parallel \vec{B}_0 \quad (\vec{k}(0,0,k_z))$.

The solution of the dispersion equation (1.23) in this case, in the range of frequencies $\omega_k \approx l\omega_0 \, (l = 1,2,...)$ is

$$\omega_k^2 = k_z^2V_{Ti}^2 - |\beta_{33}|k_z^2V_{se}^2J_l^2(b), \tag{1.33}$$

$$\gamma = -\frac{\nu_{in}}{2} + \left[\frac{\nu_{in}^2}{4} + |\beta_{33}|k_z^2V_{se}^2J_l^2(b)\right]^{\frac{1}{2}}. \tag{1.34}$$

Therefore the longitudinal electrostatic waves of ion-sonic type are excited, having the threshold of excitation of $J_l^2(\beta_{33}k_zV_{0,thr}/\omega_0) \le V_{Ti}^2/V_{se}^2$.

In the range of high harmonics $\omega_k \ll N\omega_0$, $N \gg 1$, it follows from (1.23) that longitudinal ($k_\perp = 0$) fluctuations are excited with frequencies and growth rates of:

$$\omega_{k\pm} = \frac{\varepsilon_0}{2}|\beta_{33}|\frac{k_z^2V_{se}^2}{\omega_0} \pm \left(\frac{\varepsilon^2}{4}\beta_{33}^2\frac{k_z^4V_{se}^4}{\omega_0^2} + k_z^2V_{Ti}^2\right)^{\frac{1}{2}}, \tag{1.35}$$

$$\gamma = -\nu_{in}\left(2 - \frac{\varepsilon_0|\beta_{33}|k_z^2V_{se}^2}{\omega_0\omega_{k\pm}}\right)^{-1}. \tag{1.36}$$

For the development of instability the following condition must be satisfied:

$$|\omega_{k\pm}| < \frac{\varepsilon_0}{2}|\beta_{33}|\frac{k_z^2 V_{se}^2}{\omega_0}.$$

Accordingly, from (1.35) and (1.36) follows the generation of electrostatic fast wave of ion-acoustic type, with frequency of $\omega_k = \omega_{k+}$ and threshold value of pump amplitude $V_{0,thr} \simeq N\omega_0/(k_z|\beta_{33}|)$.

The set of eqs.(1.22) also allows the polarization of the generated waves under consideration to be determined. It we take the quantity $P = V_x/V_y$ as a wave polarization characteristic, then, from (1.22), on account of (1.25), we have

$$P = \frac{V_x}{V_y} = \frac{k_x}{k_y}. \tag{1.37}$$

It is clear that the factor of polarization P is the real number and generated perturbations (1.27) and (1.31) have linear polarization in the plane x, y. The wave (1.33) and (1.35) also have linear polarization ($V_x = V_y = 0, V_z \neq 0$).

It should be noted that we ignore inhomogeneity of the pump-wave and the ionospheric medium. This imposes restrictions on the wavelengths of the disturbances under study. And the wavelengths λ are to be considered small as compared with the characteristic dimension of the inhomogeneity of the outer field L_0 and the medium L_n,

$$\lambda < \min(\lambda_0, L_0, L_n), \tag{1.38}$$

where λ_0 is the pump-wave length.

In the above discussion we neglected also the atmosphere viscosity. During the upward propagation of acoustic waves the higher frequencies are absorbed due to the viscosity of the atmosphere. The viscosity effect is small for the frequencies that satisfy the condition $\omega_0 \ll V_s/l$ (here V_s is the sound speed, l is the mean free path for the molecules) [17]. Therefore, the upper atmospheric layers are only reached by low-frequency, long acoustic waves (for instance, the waves of frequencies under $10^{-2} sec^{-1}$ propagate up to the altitude of $400 km$ without any absorption).

2. *Theory of Intensification of the Airglow in E-Region of the Ionosphere Before the Earthquakes*

2.1. *Statement of the problem and dynamic equations*

Recently, particular attention has been given to direct analysis and measurements of both steady-state and nonsteady-state wave structures in the atmosphere and ionosphere of the Earth under the action of external disturbing factor (natural or artificial). Much of the data from observations confirm that, before and after earthquakes and during man-made influence, the activity of wave processes in the atmosphere and ionosphere increases [7,8,12,13,15,18,20,26,29,30]. In particular, the excitation of acoustic-gravity (AG) waves has been observed. These disturbances are detected most easily in the infrasonic range, because, in this frequency range, the natural level of atmospheric turbulences is lower than in the range of internal gravity waves. The analysis of the data on the pre-twilight night-sky luminescence shows that, before earthquakes, an increase of the intensity of green radiation of atomic oxygen at a wave-length of 557.7 nm (the radiating layer at heights of 85-110 km) occurs. Such an increase of radiation begins several hours before the earthquake [11,23,25,32] and confirms the presence of AG (infrasonic) disturbances.

In connection with this fact, a problem of current interest is a theoretical research of the nonlinear dynamics of the AG waves (seismic and man-made origin) in the middle atmosphere and ionosphere of the Earth. On the basis of the available data from observations and experiments [7,8,11,12,13,15,18, 20,23,25,26,29,30,32], we can assume that, before earthquakes, the processes in the atmosphere develop according to the following physical scenario: Tectonic processes occurring in the Earth's crust generate the Earth's seismic-gravity oscillations or Rayleigh waves propagating over the Earth's surface from the epicenter at supersonic velocities. These waves give rise to atmospheric-pressure disturbances (AG waves) by vertical pulsed action on air. AG waves propagate vertically almost without damping to height of E-layer (their amplitudes increase exponentially as the height increases) and give rise to collective motions in the form of nonlinear AG vortices [1]. Because of the convective vortical motion, the density of atomic oxygen at these heights and, consequently, the intensity of green night-sky radiation increases. The present part of this work is devoted to the theoretical justification of this model.

We describe the propagation of wave disturbances in a nonuniform, neutral, nondissipative (the dissipation will be included below) atmosphere by the equation of motion

$$\frac{\partial \vec{v}}{\partial t} + (\vec{v}\nabla)\vec{v} = -\frac{\nabla P}{\rho} + \vec{g}, \tag{2.1}$$

the continuity equation

$$\frac{\partial \rho}{\partial t} + \vec{v}\nabla\rho + \rho div\vec{v} = 0, \tag{2.2}$$

and the equation of state (entropy)

$$\frac{\partial P}{\partial t} + (\vec{v}\nabla)P + \gamma P div\vec{v} = 0. \tag{2.3}$$

Here, as usual, $\vec{v}$ is the hydrodynamic velocity; $\rho = \rho_0(z) + \rho'$ is the density; $P = P_0(z) + P'$ is the pressure; $\vec{\mathbf{g}} = -g\vec{\mathbf{e}}_z$, g is the gravity acceleration; $\vec{e}_z$ is the unit vector along the OZ-axis; $\gamma = c_p/c_v$ is the ratio of the specific heats; the subscript zero denotes parameters of the atmosphere in undisturbed state; and the prime denotes deviations from this average state.

Let us introduce the local frame of reference x, y, and z with the x-axis directed toward the east, the y-axis directed toward the north, and the z-axis directed vertically. We will study the dynamics of short-wavelength ($|k| \gg H^{-1}$, where k is the wave number, and $H = dln\rho_0/dz$ is the characteristic scale on which the density of the atmosphere varies, i.e.,the atmosphere depth) low-frequency ($\omega \ll kc_s$, where ω is the characteristic disturbance frequency, and $c_s = (\gamma P_0/\rho_0)^{1/2}$ is the velocity of sound) atmospheric waves. The choice of this spectral interval of atmospheric disturbances is conditioned by the specific features of the problem. As the wave disturbance propagates upward, the high frequency components are absorbed rapidly due to atmospheric viscosity. The influence of viscosity is weak for the frequencies $\omega \ll kc_s/l$ (where l is the mean free path of molecules) [17]. For heights of $\sim$110 km, we have c_s $\sim$300 m/s, l $\sim$1m, and, consequently, $\omega \ll 3 \times 10^2 s^{-1}$. This means that AG waves with a frequency of (and less than) $100s^{-1}$ propagate to a height of 110 km without damping, and we can neglect viscosity in equations describing these waves.

Thus, we consider short-wavelength low-frequency waves and assume that both the pressure disturbance and the compressibility of the medium are relatively small, $\tilde{P}/P_0 \ll \tilde{\rho}/\rho_0 \ll (kH)^{-1}$ and $div\vec{v} \ll \partial v_z/\partial z$; hence, we can rewrite the set of equations (2.1)-(2.3) in the form

$$-\frac{\partial \Delta v_z}{\partial t} + (rotrot(\vec{v}\nabla)\vec{v})_z \simeq \Delta_\perp \Phi \frac{dP_0}{dz}, \tag{2.4}$$

$$\frac{\partial \Phi}{\partial t} + (\vec{v}\nabla)\Phi \simeq \frac{N^2}{dP_0/dz} v_z. \tag{2.5}$$

Here, the velocity is represented by the expression

$$\vec{v} = v_z\vec{e}_z - \nabla_\perp \Delta_\perp^{-1} \frac{\partial v_z}{\partial z} \tag{2.6}$$

the new variable is introduced as

$$\Phi = \rho^{-1} - \rho_0^{-1}, \tag{2.7}$$

N^2 is the square of the Brunt-Vaisala frequency

$$N^2 = -\left(\frac{d\rho_0^{-1}}{dz} + \frac{1}{\gamma\rho_0 P_0}\frac{dP_0}{dz}\right)\frac{dP_0}{dz}, \tag{2.8}$$

$\Delta = \Delta_\perp + \partial^2/\partial z^2$, $\Delta_\perp = \partial^2/\partial x^2 + \partial^2/\partial y^2$ is the horizontal Laplacian, and $\Delta_\perp^{-1}$ is the operator inverse to $\Delta_\perp$.

The set of two coupled equations (2.4) and (2.5) defines the nonlinear dynamics of three-dimensional waves in the nonuniform atmosphere. This set of equations conserves the energy integral:

$$E = \int\left(\frac{dP_0}{dz}(\nabla_\perp\Phi)^2 + \frac{N^2}{dP_0/dz}(\nabla v_z)^2\right)d\vec{r}. \tag{2.9}$$

In the analysis of the solution of the set of equations (2.4) and (2.5), the presence of a non-one-dimensional nonlinearity of the $rot_z rot[(\vec{v}\nabla)\vec{v}]$ type involves difficulties. If the problem is characterized by some anisotropy (e.g., in our case, a disturbance propagates along the vertical plane, and $v_y = 0$), then the set of equations (2.4) and (2.5) can be simplified by introducing the current function $\Psi(x, z, t)$ and reduced to the classical Charney-Obukhov (or Hasegava-Mima) equation [1]. In fact, in terms of the current function, the velocity components can be written in the form of the derivatives

$$v_x = -\frac{\partial\Psi}{\partial z}, \quad v_z = -\frac{\partial\Psi}{\partial x} \tag{2.10}$$

In this case, the condition for weak compressibility is satisfied automatically.

The current function must satisfy the following equations:

$$\frac{\partial\tilde{\Delta}\Psi}{\partial t} + J(\Psi, \tilde{\Delta}\Psi) = -\frac{dP_0}{dz}\frac{\partial\Phi}{\partial x}, \tag{2.11}$$

$$\frac{\partial\Phi}{\partial t} + J(\Psi, \Phi) = \frac{N^2}{dP_0/dz}\frac{\partial\Psi}{\partial x}, \tag{2.12}$$

which are obtained by substituting (2.10) into equations (2.4) and (2.5). Here, we introduced the notation

$$J(a, b) = \frac{\partial a}{\partial x}\frac{\partial b}{\partial z} - \frac{\partial a}{\partial z}\frac{\partial b}{\partial x}, \quad \tilde{\Delta} = \frac{\partial^2}{\partial x^2} + \frac{\partial^2}{\partial z^2}. \tag{2.13}$$

These equations, which describe the dynamics of two-dimensional AG atmospheric disturbances, not only resemble the Charney-Obukhov equation but also can be solved in the same way, as will be shown below.

2.2. *Self-organization of Nonlinear AG Disturbances in the Middle Atmosphere*

In accordance with the general theory of nonlinear regular structures [1,2], vortices are formed in bunches of oscillations with a phase velocity smaller than the sound velocity. The low-frequency waves under consideration satisfy this condition rather well $\omega/k \ll c_s$. In order to find other features of these waves, first, we will seek a solution to the linearized system (2.11) and (2.12) in the form of planar waves $exp\{i(k_x x + k_z z - \omega t) - z/2H\}$. In this case, we obtain the dispersion relation

$$\omega^2 = \frac{k_x^2}{K^2} N^2, \tag{2.14}$$

where $K^2 = k_x^2 + k_z^2 + 1/4H^2$.

As follows from (2.3), in atmospheric regions, in which $N^2 < 0$ or $d \ln T_0/dz > (\gamma - 1) d \ln \rho_0/dz$ (i.e., at heights of $h \sim 110$ km, where the temperature gradient is large), AG waves arise as a result of convective instability. The growth rate of these waves is $\gamma_a \sim |k_x N/K|$. The disturbance amplitudes grow with time, and the wave process becomes nonlinear. Competition with dispersive spreading may result in equilibrium. The process ends with the formation of steady-state, localized, nonlinear vortex structures. We show that, in reality, the set of nonlinear equations (2.11) and (2.12) has a steady-state solution in the form of solitary two-dimensional vortex structures.

We seek the solution to equations (2.11) and (2.12) in the form of steady-state waves propagating along the x-axis with the velocity $u = const$. We assume the Brunt-Vaisala frequency to be constant $N = N_0$; $\Psi = \Psi(\eta, z)$, $\Phi = \Phi(\eta, z)$, where $\eta = x - ut$. It is easy to show that, in the variables (η, z), the set of equations (2.11)and (2.12) is equivalent to the following equation:

$$J\left(\Psi + uz,\ \tilde{\Delta}\Psi - \frac{N^2}{u} z\right) = 0 \tag{2.15}$$

Following [2], we introduce the polar coordinates $r = (\eta^2 + z^2)^{1/2}$, $\tan\theta = z/\eta$ and a circle of radius a. Further, we require $\Psi(r, \theta)$ to tend exponentially to zero as $r \to \infty$ and to be double continuously differentiable with respect to its arguments. Then, equation (2.15) has the solution

$$\Psi = -\frac{u}{N^2}\frac{dP_0}{dz}\Phi = auF(r)\cos\theta, \tag{2.16}$$

where

$$F(r) = \begin{cases} (\beta/k)^2 J_1(kr)/J_2(ka) - (k^2 + \beta^2) r/k^2, \ r < a \\ -K_1(\beta r)/K_1(\beta a), \ \ r \geq a, \end{cases} \tag{2.17}$$

J_n is the Bessel function of the first kind, K_n is the modified Bessel function, and the following relationship between the parameters γ and k holds:

$$-\frac{J_2(ka)}{kJ_1(ka)} = \frac{K_2(\beta a)}{\beta K_1(\beta a)}, \quad \beta^2 = -\frac{N_2}{u^2}. \tag{2.18}$$

Since we construct solutions that exponentially decreases at infinity, the values of the parameter $\beta^2 = -N^2/u^2$ must be only positive, i.e., $N^2 > 0$ [or $d\ln T_0/dz > (\gamma-1)d\ln\rho_0/dz$].

Taking into account the dispersion equation (2.18), we have the two independent parameters u and a in the solution. As follows from (2.16), the solution has the asymptotic $\Psi \sim r^{1/2}\exp(-\beta r)$, as $r \to \infty$ and the wave is localized in the (η, z) plane. Current-function lines of the disturbance correspond to a dipole. Hence, these structures are pairs of oppositely rotating vortices (cyclone-anticyclone) of equal intensity, which move rectilinearly in the rest medium. Vortices can travel either toward the east $(u > 0)$ or toward the west $(u < 0)$.

The motion of medium particles in vortex structures (2.16)-(2.18) is characterized by nonzero vorticity, $rot\vec{v} \neq 0$, i.e., the particles rotate in vortices. As follows from (2.10)-(2.16), the characteristic velocity of this rotation is on the order of (or more than) the vortex velocity $v \geq u$. In this case, the vortex entrains the group of particles (the number of these particles is approximately the same as the number of transit particles); rotating, these particles move simultaneously with the vortex structure. The characteristic diameter of the vortex is $d \sim a \sim \beta^{-1} \simeq |u/N|$, and its velocity is restricted by the condition $|u| < H|N|$. Therefore, for the vortex with the height $h \sim 110$ km, we obtain the maximum vortex diameter $d \sim H \simeq 6,5$km.

Above, we considered the localization of AG waves in a nonuniform (stratified) nondissipative atmosphere. Nonlinear vortex structures are formed in those atmospheric layers in which $d\ln T_0/dz > (\gamma-1)d\ln\rho_0/dz$, i.e., at heights $h \sim 110$km. Although, at these heights, for the modes under consideration $(\omega \ll 300s^{-1})$ the influence of dissipative effects (viscosity, thermal conductivity, collisionless damping, etc.) is not substantial, [16,17,28] the dissipation cannot be ignored. The dissipation is important for the energy redistribution in the system and maintains steady nonlinear structures in the medium.

In the presence of dissipative factors, a vortex is a nonsteady wave. Hence, it is necessary to use an appropriate transport equation in order to study the dynamics of the nonlinear structures. We restrict ourselves to an analysis of the viscous atmosphere. With allowance for viscosity, we must add the term $\nu\Delta v$ (where ν is the kinetic viscosity) [17] to the right-hand side of the motion equation (2.1). In this case, the integral wave characteristic is not conserved; in particular, the energy (2.9) varies slowly with time (because of the weak dissipation). Note that the latter term of (2.9) describes the energy exchange between the vortex and the atmosphere ($N \neq 0$). In the region of the atmosphere where AG vortices are generated ($d\ln T_0/dz > (\gamma-1)d\ln\rho_0/dz$), (i.e.,at a height $h \sim 110$ km), the energy of the structure can be negative; as it will be shown below, this leads to amplification of the wave structures in the dissipative atmosphere.

Following [2], we present each wave quantity in the set of equations (2.4) and (2.5) (with the dissipation term taken into account) in the form $F = F_0 + F_1$, where F_0 is the steady-state solution (for $\nu = 0$), F_1 describes variations of the vortex space structure due to viscosity ($F_0 \gg F_1$). Hence, we find that the vortex energy varies with time according to the law

$$\frac{\partial E}{\partial t} = -\nu \int \frac{N^2}{dP_0/dz}(\tilde{\Delta} v_z)^2 d\vec{\mathbf{r}}. \tag{2.19}$$

For low dissipation, we can substitute the solution given in the form (2.16), (2.10) with parameters u and a, which vary slowly with time, into (2.19). Then, as it follows from (2.19), because the energy E in (2.9) is negative, the vortex structure should be amplified due to its dissipative energy exchange with the atmosphere.

from the energy standpoint, the region in which the atmosphere is unstable against convective perturbations (i.e., at a height $h \sim$110 km, where $N^2 < 0$) is the most appropriate for the formation of these structures; there, they are self-sustaining (like autosolitons) due to dissipative processes.

The question arises of whether AG vortices are topologically stable in a dissipative atmosphere. The numerical experiment [27] shows that flute dipole vortices in a nonuniform plasma are topologically stable and grow under the action of dissipation. Probably, a similar picture will be observed (because of the familiarity of both equations and physical causes for the generation of these structures) for AG vortices in the stratified atmosphere as well.

2.3. *Increase of the Intensity of Green E-layer Night-sky Radiation by vortices during Seismic Activity*

The main source of excitation of green night-sky luminescence ($\lambda = 557.7$ nm) at heights from 80 to 120 km is the two-stage Bart-Hildebrant mechanism consisting of the following reactions [10]:

$$\begin{aligned} &O(^3P) + O(^3P) + M = O_2^* + M, \\ &O_2^* + O(^3P) \rightarrow O_2 + O(^1S), \end{aligned} \tag{2.20}$$

which are followed by the emission of the oxygen green line via the reaction

$$O(^1S) \rightarrow O(^1D) + h\nu_0, \tag{2.21}$$

where $O(^3P)$ is an atomic oxygen in the ground state; O_2^* is a molecular oxygen in the excited state; M is any neutral particle; $O(^1S)$ and $O(^1D)$ are oxygen atoms in excited states; and ν_0 is the green-line frequency. The radiation intensity I depends strongly on the concentration of the atomic oxygen $n(O)$ in the radiating layer [32]:

$$I \sim \left(\frac{T_0}{300}\right)^{-2.9} n^3(O). \tag{2.22}$$

The variations of $n(O)$ at a height of ~100 km may be associated with the occurrence of AG vortices (2.16), which cause the convective mixing of oxygen in the vertical direction and, consequently, increase the concentration $n(O)$ in the radiating atmospheric layer. Let us estimate how the concentration of the atomic oxygen $n'(O)$ in the radiating layer ($h \sim 100$ km) increases due to AG vortices. In accordance with (2.7) and (2.16)

$$\Phi = -\frac{N^2}{u}\left(\frac{dP_0}{dz}\right)^{-1}\Psi = -\frac{\rho'}{\rho_0^2}, \tag{2.23}$$

where ρ' is the variation of the density of atomic oxygen due to vortical convection.

In accordance with (2.23) and (2.16), we have

$$\frac{\rho'}{\rho_0} = \frac{n'(O)}{n(O)} \sim \rho_0 u \left|\frac{N}{dP_0/dz}\right|. \tag{2.24}$$

Taking into account that $\rho_0 = m(O)n(O)$ is the atomic mass of oxygen and $\rho = \rho_0 + \rho'$, from (2.22) and (2.24) we obtain

$$\frac{I_{obs}}{I_{seas}} \sim \left(\frac{\rho}{\rho_0}\right)^3 \sim \left|\rho_0 u \left|\frac{N}{dP_0/dz}\right| + 1\right|^3. \tag{2.25}$$

Here, I_{obs} and I_{seas} are the observed intensity and that averaged over the season, respectively.

In [11,23,32], on the basis of long-term observations, the variations of the intensity of the main kinds of radiation of the upper atmosphere before earthquakes were shown. This effect is most pronounced in variations of green night radiations of atomic oxygen ($\lambda = 557.7$ nm); its radiating layer is localized at a height about 100 km. The difference $\Delta I = I_{obs} - I_{seas}$ was analyzed. It was found that 24 h before earthquakes the intensity of the emission $\lambda = 557.7$ nm begins to increase. The amplitude $\Delta I/I_{seas}$, expressed in percents, is $6.9 \pm 1.4\%$.

For a comparison with the observed intensity, by using (2.25), we determine the relative variation of the intensity $\Delta I/I_{seas}$ of green radiation due to AG vortices:

$$\frac{\Delta I}{I_{seas}} \sim \left(\rho_0 \left|\frac{uN}{dP_0/dz}\right| + 1\right)^3 - 1. \tag{2.26}$$

We now perform some estimations. For typical values of the atmospheric parameters at a height of 110 km ($dP_0/dz \sim \rho_0 c_s^2/H,\ c_s \sim 300\ m/s$, $H = 6.5 \times 10^3$ m, and $N = 2 \times 10^{-2} s^{-1}$), in accordance with (2.26), we obtain $\Delta I/I_{seas} = 7\%$ for the vortex velocity u as small as 16 m/s (the velocity u is restricted only by the condition $u < H|N| = 130$ m/s), i.e., the observed variations in the intensity (before earthquakes) $\Delta I/I_{seas}$ can be easily ensured.

Hence, the observed variations of the emission at $\lambda = 557.7$ nm, which are related to enhanced seismic activity, are quantitatively consistent with the proposed theory of the propagation of AG vortices (seismic in nature) in the middle atmosphere.

Note that it is difficult to observe AG vortices in the middle atmosphere, because it is difficult to identify properly the perturbations of the parameters of the neutral atmosphere. However, in the lower ionosphere (E-layer, $h \sim 100$ km), by means of vertical probing, we can find many effects related to disturbances in the atmospheric parameters and indicating the propagation of AG vortices. Although the density of charged particles in the ionospheric E-layer is low and weakly affects the dynamics of AG vortices, these vortices may cause a strong plasma disturbance in the E-layer. Actually, in this layer, electrons move predominantly along the geomagnetic field, because their gyrofrequency is much higher than their collisional frequency, whereas most ions are entrained by the neutral wind across the geometric field. Thereby, ions generate a current in the transverse direction with respect to the field and strongly disturb the ionospheric plasma. Let us estimate the disturbance of the geomagnetic field $\vec{B}_1$, which is caused by the propagation of AG vortices.

In accordance with the Maxwell equation,

$$rot\vec{B}_1 = \frac{4\pi e}{e} n_{0i} \vec{v}_\perp, \tag{2.27}$$

where n_{0i} is the equilibrium ion density, and $\vec{v}_\perp$ is the neutral-wind velocity across the geomagnetic field. Taking into account (2.10), from (2.27), we obtain

$$B_{1y} \simeq \frac{4\pi e n_{0i}}{c} \Psi. \tag{2.28}$$

It is seen that the disturbed magnetic field is localized in the same way as AG vortices. By means of earthbound and satellite observations, we can determine the character of disturbances of the geomagnetic field (2.28); the localized closed field lines can be associated with AG vortices.

3. *Discussion of Results and Conclusion*

1) New mechanism of transformation of Low-frequency acoustic waves (of natural or artificial origin) into electromagnetic ones in a conductivity medium is suggested. The possibility of nonlinear transformation of acoustic waves of seismic origin into electromagnetic ones in the F-layers of a quasi-homogeneous atmosphere is examined. The range of frequencies covering geomagnetic pulsations $\nu_{ni} \ll \omega \ll \omega_{ci}$ has been studied. The phase of the outer acoustic wave (pumping) is arbitrary. This is the main difference of our consideration from the usual parameters [31]. As a result of averaging over the arbitrary phases we will have one vector equation instead of the infinite system of connected equations for particle velocities (which is intrinsic to the parameters). It is important that the equation contains the contribution of electrons motion relative to ions (the last term in the equation (1.22)), which causes the existence of non-stable modes. This instability can be called the parametric instability in the general meaning of this word [24].

2) It is shown that when low-frequency acoustic wave propagates in the ionosphere, it involves charged particles in a collective motion due to collisions, the friction of electrons with neutral particles not being the same as that of ions. Relative Motion of charged particles excites the alternating current with arbitrary phase and consequently, it causes generation of electromagnetic and electrostatic fields due to parametric interaction of low-frequency acoustic waves with the ionospheric medium. The frequencies of the generated oscillations and the growth rate of the parametric instability are determined. The excitement thresholds and the polarization of the generated disturbances are estimated. This mechanism can explain an excitation of electromagnetic perturbations within ranges of frequency $10^{-4}s^{-1} < \omega < 10^{2}s^{-1}$, periods (T) 2.8 $h > T > 10^{-2}s$, and wavelengths $0,5m < \lambda < 10^{3}km$ on ionospheric F-layers ($h \geq 150km$).

3) The theory stated here makes possible a qualitative interpretation of the initiation of the acoustic and electromagnetic disturbances in the atmosphere that have been observed in several experiments: low-frequency acoustic waves (infrasound) generated at the Earth-Air boundary (by the Earth's seismic-gravity oscillations or due to Rayleigh waves), after reaching the ionospheric layer, bring about disturbances of the current-carrying

medium parameters and transform their energy into that of electromagnetic emissions in the Ultra-Low-Frequence (ULF) and Extreme-Low-Frequence (ELF) range with high efficiency.

4) Thus the emergence of intense low-frequency acoustic, electromagnetic, and electrostatic disturbances in the atmosphere and ionospheric layers can be identified as a forerunner of an earthquake or with medium response on the influence of strong perturbing factors (perturbation source). Knowing the polarization of the generated field, one can evaluate the direction toward the perturbation source center.

5) The suggested mechanism has been also successfully applied for explanation of a very elegant and largely unexplained phenomenon, Sonoluminescence, which has been puzzling physicists for more than six decades.

6) We would like to emphasize that the above mentioned results, which certainly need further elaboration, clearly demonstrate that the mechanism is capable to explain quite a diverse range of phenomena, starting from generation of electromagnetic radiation in the certain layers of the ionosphere to a tiny air bubble in the flask with water (Sonoluminescence). Thus, the mechanism can possibly be corroborated by satellite as well as by laboratory equipment.

7) The problem of the propagation of low-frequency nonlinear acoustic-gravity disturbances (seismic or man-made in nature) in the stratified middle atmosphere and ionosphere is investigated. The simplified set of three and two-dimensional nonlinear equations for the dynamics of AG waves in the stratified atmosphere and ionosphere was obtained.

8) It was shown that AG disturbances propagating from lower atmospheric layers practically without damping reach heights of $h \sim 110km$, where they become unstable against convection ($N^2 < 0$), and, as a result of this process, strongly localized AG vortices are formed. The solitary wave consists of a pair of dipole vortices rotating in the opposite directions in the x, z plane (the cyclone and the anticyclone). Vortex structures entrain neutral particles and can travel with them either toward the east or toward the west. Under the action of disturbances, they can also move upward and downward. In this case, they will be damped, because they leave the appropriate active region ($N^2 > 0$).

9) It was found that the energy of AG vortices becomes negative at a height of $h \sim 110km$; the formation of these structures is advantageous from the energy standpoint, and their generation is self-maintained due to the dissipative redistribution of the internal potential energy of the medium. At this height, the source of the potential energy is the heat released in the recombination of the atomic oxygen.

10) The additional mixing of neutrals, which is related to AG vortices and caused by the transport of atmospheric components in the vertical

direction, increases the neutral density in the atmosphere at this height. The increase of the density of oxygen atoms increases the efficiency of their recombination and, consequently, increases the intensity of green ($\lambda = 557.7nm$) night-sky radiation.

11) The fluctuations of the density of the atomic oxygen, which are caused by AG vortices, and the corresponding increase in the intensity of green night-sky luminescence are estimated. The comparison of the theoretical estimate for the increase of the intensity of the green line $\lambda = 557.7nm$ with the data from observations conducted before strong earthquakes shows that they are in good agreement.

12) In the ionosphere, AG vortices can be detected from the pulsations of the geomagnetic field, which are caused by the currents generated by ions entrained by the neutral wind. In the E-layer, the disturbed magnetic field is localized in the same way as AG vortices; hence, the closed field lines can be identified as AG vortices.

13) We can conclude that the variations of the radiation at $\lambda = 557.7nm$ at periods of the seismic activity may be caused by the generation of AG vortices in the middle atmosphere; hence, they can be considered to be a short-term $(24 - 48h)$ sign of earthquakes.

14) Results of the above investigation permit to deduce, that dynamical processes in neutral atmosphere play an important role in transference of lithosphere perturbations into ionosphere and magnetosphere.. So, the noises of lowfrequent electromagnetic and optical range which are observed in ionosphere can be stipulated by propagation acoustic type waves.

15)On the whole, it is possible to confirm, that several days before earthquake variation spectrums of ionosphere parameters were modified on different heights over the earth. Information about the modifications of spectrum are used for creating physical and mathematical models for phenomena, which proceed in the bowel and casings of the earth and for decision of problem about operative forecast earthquake.

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