the exchange constant and a is the lattice constant). Taking as a very rough estimate $\theta \sim H_E$, we find that in the experiments at 9.4 GHz the largest values of q = ak for the excited spin waves are $q_{max} \sim 3 \times 10^{-3}$. For 36 GHz we get $q_{max} \sim 2 \times 10^{-2}$. It should be noted that at $q \sim 10^{-2}$ an important role can already be assumed by the anisotropy of the exchange interaction [2].

However, the smallness of the dipole interaction of the spin waves in an antiferromagnetic raises difficulties when it comes to explaining theoretically the relatively low threshold of the parametric resonance observed by us. In particular, it follows from Ozhogin's calculations [4] that, in the first approximation, parametric excitation of AFMR and spin waves is possible only in the presence of weak ferromagnetism.

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Addendum. After submitting this article, we became acquainted with a paper by Seavey [5], who observed parametric excitation of the electron spin waves in $CsMnF_3$ at 17.5 GHz.

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MAGNETORELATIVISTIC MODEL OF A PULSAR PULSE

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The pulsed character of pulsar radiation is usually connected (this being the most probable explanation) with the existence of a definite directivity of their radiation pattern, which can be either "pencil-like" or "knife-like." It is the rotation of this diagram in space which determines the observed radiation in the form of individual pulses. In the heretofore proposed models, the directivity of the radiation pattern is attributed to geometric factors, viz., detachment of plasma clouds in the equatorial plane of the part of the magnetosphere of a neutron star rotating with relativistic velocity [1], a special distribution of the radiating particles in "convenient" parts of the magnetosphere [2], etc. It has been impossible to date, however, to obtain from physical considerations quantitative estimates of the aperture $\Delta\nu$ of the directivity pattern. On the other hand, it is well known that the directivity is an invariant property of relativistic objects.

In this connection, we can propose the following model of a pulsar.

1. Assume that the pulsar is a neutron star with a dipole (or more complicated) magnetic field, rotating or oscillating in such a way that the pulsar surface regions containing the magnetic fields move with relativistic velocities. We are unable to justify such a model rigorously at present, but we can advance some qualitative considerations in

it favor.

It is known that at least part of the pulsar has two periods. The larger period is usually connected with the pulsar rotation, and the smaller with its pulsation. The pulsar rotation velocity determined from the observed period $T = 3 \times 10^{-3} - 4$ sec and from the specified radius of the neutron star, $a \approx 10^{6}$ cm, is much smaller than the speed of light. But the pulsar surface velocity during the pulsations is unknown, and the assumption that it is close to that of light, during certain phases, does not contradict the observations.

Another model is also possible. For example, the pulsar itself may rotate and have a relativistic linear velocity on its equator, and in addition it may process with a much smaller velocity. Favoring this variant is the following consideration. First, if a star such as the sun (and all the more a rapidly-rotating star of earlier spectral class) collapses, while conserving its angular momentum, to a neutron star with dimension close to the gravitational radius, then relativistic rotation is inevitable. The neutron star then turns out to be non-spherical and the appearance of precession as a result of the special features of the gravitational field in general relativity theory is perfectly possible here. Of course, this question should be considered in greater detail, but we note that precession in the field of a rotating body is well known in general relativity theory (cf. [3, 4]). In analogy with these problems, we can expect also precession of a body with relativistic angular velocity $\Omega \sim GM/ca^2$, where G is the gravitational constant, M and a the mass and radius of the body, and c the velocity of light (particularly if account is taken of the possible rotation inhomogeneity, analogous, say, to the equatorial acceleration observed in the sun). Incidentally, within the framework of the assumed precessional motion, it is easier to explain the jumps in the period of the Vela pulsar.

2. The poles of a dipole having a magnetic moment m and moving with relativistic velocity radiate long-wave electromagnetic fields. Such a radiation can be calculated by means of the usual formulas of synchrotron radiation, in which the electric charge q is replaced by the magnetic charge $\mu = m/2a$ (where 2a is the distance between the charges $\pm \mu$ of the dipole), and replacing $\vec{E}$ by $\vec{H}$ and $\vec{H}$ by $-\vec{E}$ (where $\vec{E}$ and H are the intensities of the electric and magnetic fields of the long-wave radiation. It follows therefore that this radiation is directional; the aperture angle of the directivity pattern is determined by the degree of relativistic motion of the magnetic charges: $\Delta\nu \approx 1/\gamma = (1 - u^2/c^2)^{1/2}$ (u is the linear velocity of the poles). In the rotation + precession model, this diagram rotates rapidly in the equatorial plane ("filling" it, as it were), and more slowly around the precession axis. In the model with relativistic pulsations, the directivity pattern rotates only in the equatorial plane of rotation. The rotation of the directivity pattern of the long-wave radiation forms its "video pulse."

3. The electrons of the plasma surrounding the pulsar, falling into the region of the directivity pattern of this long-wave radiation, are accelerated to ultra-relativistic energies, moving predominantly along the instantaneous axis of the directivity pattern. Emerging from the long-wave radiation cone, these electrons are rapidly decelerated in the residual external magnetic field, producing synchrotron and Compton radiation. The long wave

radiation of the relativistic magnetic poles is thereby converted into high-frequency electromagnetic radiation (including x-rays and gamma rays), i.e., the "carrier" of the video pulse.

We now consider quantitative relations. The relativistic character of the magnetic-pole motion, regardless of the assumptions concerning their character, is determined by the ratio of the observed pulse duration Δt and their repetition period T. According to the general properties of the emission of relativistic charges (including magnetic poles) we have $\Delta t/T \approx 1/2\gamma^3$. The observed pulsed character of the pulsar radiation can therefore be attributed to the very weak relativism of the magnetic poles, $\gamma \approx 2 - 3$.

Let us assume that the long-wave radiation of the magnetic pole μ, once converted into high-frequency radiation, accounts for the entire pulsar luminosity L in all the bands. The electric and magnetic fields of this radiation, in the wave zone at a distance R from the source, are then given by

$$E \approx H \approx \frac{2\pi\mu}{\lambda R} \approx \frac{\gamma}{R}\left(6\frac{L}{c}\right)^{1/2}, \tag{1}$$

where λ is the wavelength. For a pulsar rotating with relativistic velocity $\lambda = \pi a/\gamma^3$. For real pulsars, $\lambda \approx 10^4 - 10^5$ cm. Assuming $L = 10^{38}$ erg/sec we obtain from (1) $\mu \approx 10^{19}$ cgs esu, i.e., if the magnetic field on the pulsar surface is $H_0 \approx 10^{12}$G, then the pole dimension is $\sqrt{\mu/H_0} \approx 3 \times 10^5$ cm. The electric and magnetic fields in the wave zone at $R \approx 10^8$ cm are of the order of 10^6 cgs esu (Gauss). If we neglect the polarization of the medium around the pulsar, then the electric and magnetic fields of the long-wave radiation are crossed and are equal in magnitude. In this case they accelerate the charges to an energy on the order of $E \approx (mc^2 \mathcal{E}_0^2)^{1/3}$, where $\mathcal{E}_0 = e\lambda E \approx 10^1$ erg (see [3, 5], also dealing with acceleration of charges by long-wave radiation of the magnetic dipole moments of pulsars, but with relativism). Thus, the electrons are accelerated at least to an energy of the order of 10^5 MeV. The electron concentration needed to fill the video pulse is determined from the obvious formula $n \approx L\lambda\,\gamma^2/c\mathcal{E}R^3\ 10^9\ \text{cm}^{-3}$.

The polarization of the medium (the dipole moment of the volume), which is produced by the long-wave radiation, cannot exceed $P = ne\lambda \approx 10^4$ cgs esu, i.e., it is smaller by two orders of magnitude than the field itself. Consequently the long-wave radiation can propagate over large distances. On the other hand, allowance for the polarization makes E and H no longer equal (E increases), contributing to the acceleration of the particles to high energies. In addition, an additional narrowing of the directivity pattern of the long-wave radiation is possible here, owing to the self-focusing in the relativistic plasma. The authors plan to examine the proposed scheme in detail in a future article.

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HYBRID SURFACE PLASMONS ON A METAL-SEMICONDUCTOR BOUNDARY

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Interband transitions in metals have a strong influence on the character of the plasma oscillations [1]. In some cases, allowance for the dispersion of the electronic part of $\varepsilon(\omega)$, connected with the interband transitions, leads to the occurrence of "hybrid" volume plasmons, the frequency of which is much lower than that of the plasma frequency of the free electrons. Such a "hybrid" plasmon was observed experimentally in Ag [2]. A similar situation arises on the boundary between a metal and a semiconductor, where allowance for the dispersion of Re $\varepsilon(\omega)$ near the singularities of the interband density of states of the semiconductor leads to the appearance of surface waves of a new type.

Let us consider the contact between a metal and a semiconductor. Assume for simplicity that their work functions are equal or are such that the region where the bending of the bands plays a role is small compared with the wavelengths that are of importance.

Proceeding in the usual manner [3], without allowance for the retardation effects, we obtain the solutions of Maxwell's equations of the surface type in the form:

$$E_0^{(1)} \exp[i(k_1x + k_2y + i\kappa_1 z) \quad (z > 0 - \text{semiconductor})$$

$$E_0^{(2)} \exp[i(k_1x + k_2y - i\kappa_2 z) \quad (z < 0 - \text{metal})$$

Using the ordinary boundary conditions, neglecting spatial dispersion, and assuming the semiconductor and the metal to be isotropic, we obtain

$$\kappa_1 = \kappa_2 = \sqrt{k_1^2 + k_2^2}, \tag{1}$$

$$\varepsilon^{(1)}(\omega) + \varepsilon^{(2)}(\omega) = 0, \tag{2}$$

where $\varepsilon^{(1)}(\omega)$ and $\varepsilon^{(2)}(\omega)$ are respectively the dielectric constants of the semiconductor and of the metal. Equation (2) determines the possible frequencies of the surface oscillations. Taking for $\varepsilon^{(2)}(\omega)$ the simple plasma expression

$$\varepsilon^{(2)}(\omega) = 1 - \frac{\omega_{p2}^2}{\omega^2}, \tag{3}$$

where ω_{p2} is the plasma frequency of the electrons in the metal, and disregarding the frequency dispersion of $\varepsilon^{(1)}(\omega)$, we obtain from (2) the well known solution for the frequency of the surface plasmon, $\omega = \omega_{p2}[1 + \varepsilon^{(1)}(\infty)]^{-1/2}$. Actually, $\varepsilon^{(1)}(\omega)$ can have a strong frequency dependence near certain singular points of the interband density of states [2, 4]. In these cases, Eq. (2) should contain additional new solutions at frequencies $\tilde{\omega} \ll \omega_{p2}[1 + \varepsilon^{(1)}(\infty)]^{-1/2}$.