Gill SM

National Aeronautics and Space Administration
Goddard Space Flight Center
Contract No. NAS-5-9299

ST-PF-SP-CR-IGA-10 452

DYNAMIC DISSIPATION OF THE MAGNETIC FIELD AND ACCELERATION OF PARTICLES

by

S. I. Syrovatskiy

(USSR)

FACILITY FORM 602

N66 87193
(ACCESSION NUMBER)

27
(PAGES)

CR 78099
(NASA CR OR TMX OR AD NUMBER)

(THRU)

None
(CODE)

(CATEGORY)

3 MARCH 1966

ST-PF-SP-CR-IGA-10452

DYNAMIC DISSIPATION OF THE MAGNETIC FIELD AND ACCELERATION OF PARTICLES *

Preprint **
Fizicheskiy Institut im.
P. N. Lebedeva ANSSSR
MOSKVA, 1965

by S. I. Syrovatskiy

SUMMARY

A mechanism of "dynamic" dissipation of the magnetic field is suggested, which leads to direct magnetic energy transfer into the energy of fast particles. This mechanism consists in that, under specific conditions, the motion of the conducting medium leads to the emergence of regions with a great magnetic field gradient and a simultaneous high rarefaction. In such regions, the current maximum, which may be assured by the conducting medium, is insufficient for the preservation of the magnetic field gradients. As a result, there appears a strong electric field accelerating the charged particles. The reserve of the remnant magnetic energy passes into kinetic energy of charged particles.

The application of the obtained criteria shows that dynamic dissipation conditions may be realized at chromospheric flares on the Sun, in turbulent shells of supernovae, in exploding galactic nuclei and in radiogalaxies; they may also be materialized in the Earth's magnetosphere and in certain laboratory experiments. For such objects, the dynamic dissipation apparently serves as the basic acceleration mechanism of particles to greater energies and, in particular, of cosmic ray generation.

* * *

A series of events are known in the cosmic plasma, for which there is either observed, or has to be inescapably admitted a process of rapid magnetic field dissipation accompanied by the emergence of high-energy particles. To such events we refer: the supernova bursts, the explosions of galactic nuclei and of radiogalaxies, at which high-energy cosmic rays are emerging [1]; solar chronospheric flares generating sub-cosmic particles

---

* DINAMICHESKAYA DISSIPATSIYA MAGNITNOGO POLYA I USKORENIYE CHASTITS.

** Destined for publication in the next issue of "Astronomicheskiy Zhurnal".

and cosmic rays of moderate energies [2, 3]; the appearance of fast particles in the magnetosphere and its train [4, 5] and also certain processes in laboratory plasma, at which neutron and hard radiation appear [6].

All these events are linked with essentially nonstationary dynamic processes in the plasma with the presence of intense magnetic fields. At present, the most practical object for the study of such processes are the chromospheric flares on the Sun; this is explained by the large volume of accumulated information and by comparatively favorable possibilities for complex observations. As a concrete example we shall emphasize below on the chromospheric flares, though the results actually do have a more general scope.

The dissipation and the changeover of the magnetic field during flares on the Sun, as well as the acceleration of particles last for 10^2 — 10^3 sec, whereas the spatial scale of the event constitutes from 10^9 to 10^{10} cm. This means that the process unfolds with a velocity up to 10^8 cm/sec, that is, with a velocity corresponding to fast magnetohydrodynamic waves in the magnetic fields of sunspots. If, at the same time, we have in mind particles with energy of several hundred Mev and higher, the energy of accelerated particles for powerful flares, that is, $\sim 10^{30}$ erg, constitutes the essential part of the total energy yield, which is $\sim 10^{32}$ erg.

As is shown by the closest consideration, the dissipation of the magnetic field in this process cannot be viewed as a Joule dissipation [7] (the meaning of this term is clarified in #3). On the other hand, the brief appearance of energetic particles coul most naturally of all be connected with strong electric field [8, 9]. However, because of high plasma conductivity, such fields cannot be assumed under standard conditions.

All the above-said leads to the conclusion of the existence of a substantially more effective dissipation mechanism, which may, under specific conditions, ensure a rapid transformation of magnetic energy into the energy of accelerated particles. It will be shown below that for a nonstationary plasma such a mechanism does exist indeed, and the conditions for its realization will be ascertained. Inasmuch as the plasma dynamics as a whole is closely linked with that of separate particles, we shall designate it in the following, if only for brevity, as a dynamic dissipation mechanism.

1. - TWO-DIMENSIONAL MAGNETOHYDRODYNAMIC FLOWS

In connection with the problem of solar flares the dynamics of the magnetized plasma was the object of consideration in a series of works. Thus, in the work [10], consideration was given to the one-dimensional plasma compression in the vicinity of the zero or neutral magnetic field plane. However, as was subsequently shown in [11], it was impossible to agree with the conclusion obtained in [10] on the instability of the magnetohydrodynamic equilibrium near the neutral plane. The uni-dimensional compression leads either to a new state of magnetohydrodynamic equilibrium, or, when taking into account the nonadiabadicity of the compression, (see [13]) — to a compression with a velocity determined by the rate of heat removal from the system. In this connection, any rapid change-over and dissipation processes of the magnetic field should at any rate be sought for in the category of two-dimensional flows.

The first step in this direction was made in the work [14], in which the uni-dimensional stationary plasma compression with the field, taking into account the plasma outflow in the transverse direction along the layer, is treated approximately. In this case, the dissipation at the expense of field gradients increase, may be increased; however, it will remain Joule-type, without providing a sufficiently rapid transformation of the magnetic energy into the thermal energy [7], without even speaking of absence in such a scheme of conditions for the generation of fast particles.

An essential step forward in the development of the model [14] was made in [15]. In the latter work the two-dimensional stationary plasma compression near the neutral point was considered, taking into account the inclined magnetohydrodynamic waves. This scheme allows to convert the energy of the magnetic field into the energy of heated gas at an apparently sufficient rate, as a result of the Joule dissipation; in this way it can explain the rapid heating and the ejections of plasma in the region of solar flares. At the same time, the question of fast particle generation, alongside with a series of other peculiarities characteristic of flares, remains open, for it cannot be explained by this model.

Considered in the model [15] was an earlier assigned stationary flow pattern near the neutral point, appearing to be a generalization of the pattern considered in [14]. At the same time, the question of the emergence of precisely such a flow has been set aside.

In this connection it appears to be indispensable to consider a more generalized two-dimensional problem of magnetized plasma flow, emerging from the given initial state.

Let us consider the two-dimensional flows for which the velocity vectors $\vec{v} = \{v_x, v_y, 0\}$ and the intensities of the magnetic field $\vec{H} = \{H_x, H_y, 0\}$ are parallel to the plane (x, y) and, besides the time t, are dependent only on the coordinates x and y. It is practical to introduce for such a case the vector-potential of the magnetic field $\vec{A}$, such that

$$\begin{aligned} \vec{A} &= \{0, 0, A(x, y)\} \\ \vec{H} = \operatorname{rot} \vec{A} &= [\nabla A \cdot \vec{k}] = \left\{\frac{\partial A}{\partial y}, -\frac{\partial A}{\partial x}, 0\right\} \\ \vec{j} &= \frac{c}{4\pi} \operatorname{rot} \vec{H} = -\frac{c}{4\pi} \Delta A \vec{k}, \end{aligned} \tag{1}$$

where $\vec{j}$ is the current density in the plasma, $\vec{k}$ is a unitary vector along the axis z. With such a determination of the vector-potential the equation of the magnetic field's line of force is simply

$$A(x, y) = \text{const} \tag{2}$$

and the magnetic flux in the tube bounded by the lines of force $A(x, y) = A_1$ and $A(x, y) = A_2$ is equal to

$$\Phi = \int_{A_1}^{A_2} \vec{H} [\vec{k} d\vec{l}] = -\int_{A_1}^{A_2} \nabla A \, d\vec{l} = A_1 - A_2. \tag{3}$$

In these denotations the magnetohydrodynamics equation will take the form:

$$\begin{aligned} \frac{d\rho}{dt} &\equiv \frac{\partial \rho}{\partial t} + (\vec{v}\nabla)\rho = -\rho \operatorname{div} \vec{v} \\ \frac{dA}{dt} &\equiv \frac{\partial A}{\partial t} + (\vec{v}\nabla) A = 0 \\ \frac{d\vec{v}}{dt} &= -\frac{1}{\rho} \nabla p - \frac{1}{4\pi\rho} \Delta A \nabla A \end{aligned} \tag{4}$$

where p and ρ are respectively the gas pressure and the plasma density, the dissipative terms being dropped.

The first two equations have integrals which are easy to find, passing to Lagrange coordinates. Precisely, the continuity equation is reduced to

$$\rho(\vec{r}, t) = \rho_0 (\vec{r} - (\vec{r}, t) \frac{D(\vec{r} - \vec{\xi}(\vec{r}, t)}{D(\vec{r})}, \quad (5)$$

where $\rho_0(\vec{r})$ is the plasma density at the time $t = 0$, $\vec{\xi}(r, t)$ is the vector of the initial position shift for a particle situated at the moment of time t at the point $\vec{r}$, and $D(\vec{r}_0)/D(\vec{r})$ are the Jacobi transformations from the coordinates $\vec{r}_0 = \vec{r} - \vec{\xi}(\vec{r}, t)$ to coordinates $\vec{r}$.

The second of the equations (4) expressed the well known condition of "freeze-in", according to which, at medium's motion, the magnetic line of force remains linked with the same particles of the medium. If $A_0(\vec{r})$ is the value of the vector-potential at the initial moment of time, then at the moment of time t we shall have

$$A(\vec{r}, t) = A_0 [\vec{r} - \vec{\xi}(\vec{r}, t)] \quad (6)$$

Therefore, for a well known plasma shift $\vec{\xi}(\vec{r}, t)$ the equations (5) and (6) determine the density of the plasma and the line of force pattern of the distorted magnetic field directly for any moment of time.

In order to determine the shift $\vec{\xi}(\vec{r}, t)$ (or the velocities of the medium $\vec{v}(\vec{r}, t) = d\vec{\xi}/dt$) one must, however, resolve the entire system of equations (4) with the corresponding initial conditions. Because of the nonlinearity of the equations, this problem is found to be very complex in the general case. Alongside with this, the system (4) may, under specific conditions, be considerably simplified.

Let us examine, first of all, the motion of plasma in regions, where the condition

$$p \ll \frac{H^2}{8\pi} \quad (7)$$

In such regions the dynamics of the plasma is practically entirely determined by magnetic intensities, so that the term in the last equation (4) containing the pressure gradient, may be dropped. As a result, we have

$$\frac{d\vec{v}}{dt} = - \frac{1}{4\pi p} \Delta A \nabla A \quad (8)$$

The condition (7) is equivalent to the condition

$$s^2 \ll V_A^2 \tag{9}$$

where $s^2 = (\partial p/\partial \rho)_S = \gamma p/\rho$ is the square of sound's adiabatic velocity in the plasma, and $V_A = H/\sqrt{4\pi\rho}$ - is the characteristic magnetohydrodynamic (Alfvén) velocity.

As may be seen from the equation (8) the plasma acceleration, and consequently, at small perturbation the shift $\vec{\xi}(\vec{r}, t)$, are directed across the magnetic lines of force, that is, along ∇A. At the same time it is easy to ascertain that the perturbation propagates in the plasma with (local) Alfvén velocity. Let us consider, in reality, the small perturbation of the initial state of equilibrium

$$\vec{v} = 0, \quad p = p_0 = \text{const}, \quad \rho = \rho_0 = \text{const}, \quad \Delta A_0\,(x, y) = 0.$$

Then, we shall obtain for the quantities $\vec{v}$, A' and ρ', characterizing the small perturbation, from the equations (4) in linear approximation :

$$\begin{aligned} \frac{\partial \rho'}{\partial t} &= -\rho_0 \operatorname{div} \vec{v} \\ \frac{\partial A'}{\partial t} &= -\vec{v}\nabla A_0 \\ \frac{\partial \vec{v}}{\partial t} &= -\frac{s^2}{\rho_0}\nabla\rho' - \frac{1}{4\pi\rho_0}\cdot\Delta A'\nabla A_0 . \end{aligned} \tag{10}$$

Hence

$$\begin{aligned} \frac{\partial^2 A'}{\partial t^2} &= \frac{(\nabla A_0)^2}{4\pi\rho_0}\Delta A' + \frac{s^2}{\rho_0}\nabla A_0 \nabla\rho' \\ \frac{\partial^2 \rho'}{\partial t^2} &= s^2\Delta\rho' + \frac{1}{4\pi}\nabla A_0\,\nabla(\Delta A') \end{aligned} \tag{11}$$

As is clearly seen from (11), when neglecting the density gradient, that is, at fulfillment of the condition (9), the perturbations of the magnetic field spread with local magnetohydrodynamic velocity V_A, where

$V_A^2 = (\nabla A_0)^2 / 4\pi\rho_0 = \frac{H_0^2}{4\pi\rho_0}$ / H_0 is the intensity of the unperturbed magnetic field.

As to plasma shifts along the magnetic field, that is, in directions for which $\vec{\xi}\nabla A = 0$, such perturbations, as may be seen from (10) - (11), correspond to density disturbances propagating with the sonic speed along the lines of force, that is, by the strength of (9) — substantially slower than the transverse perturbations.

Thus, at fulfillment of the conditions (7), (9), the perturbation of the initial state of equilibrium leads in the first place to establishment of a certain rapid (with velocity V_A) quasi-equilibrium state by way of shift across the lines of force, and only then — substantially slower (that is, with the speed of sound s) to the establishment of equilibrium along the lines of force.

As may be seen from the equation (8), the condition for the "transverse quasiequilibrium" has the form

$$\Delta A\ (x, y, t) = 0 \qquad (12)$$

This condition simply implies that when neglecting plasma pressure, only a free (currentless) field may be equilibrium *.

We shall estimate that the boundary conditions of the problem ry sufficiently slowly by comparison with the Alfvén velocity, but sufficiently rapidly by comparison with the speed of sound (sonic velocity). By the strength of (9) such a region does exist. Then from the equation (8) we may drop the acceleration, that is, we may estimate the deformation as being adiabatically slow. At the same time, the system will pass through a continuous series of quasiequilibrium states satisfying the equation (12). This equation, alongside with the "freeze-in" condition (6) and the condition for shift's "transversality"

$$[\ \vec{v}(\vec{r}, t)\ \nabla A\ (\vec{r}, t)\ ] = 0 \qquad (13)$$

allow to determine the shift $\vec{\xi}(\vec{r}, t)$ for any given dosplacement of field's sources.

* In the considered two-dimensional case the so-called force-free fields, with current density different from zero, do not exist.

2. - DEFORMATION OF THE MAGNETIC FIELD IN THE VICINITY OF ZERO (NEUTRAL) POINTS

If we have in mind a rapid dissipation (the term "annihilation" of the magnetic field is sometimes used), it is most natural of all to link it with the zero (neutral) points, inasmuch as at such points magnetic fields of opposite directions join. Attention to the role of neutral points in the solar flare events was first drawn by Giovanelli [16]. Subsequently, numerous observations in that direction were conducted by Severnyy [10].

In the first approximation the magnetic field of bipolar groups of solar flares in the chromosphere may be considered as a field of the magnetic dipole. The neutral points emerge in the presence of two or more such groups (multipolar groups) in places where the magnetic fields of separate pairs of spots have the opposite direction.

It is essential for the following that only events near the neutral point are to be considered and that the concrete form of field sources plays no part; that is why we shall consider, for the sake of simplicity, a neutral point having emerged between two parallel currents.

Let two parallel currents of equal strength I_o be situated at the points $\pm I$ along the axis x (Fig. 1). Let us choose the vector-potential of these currents' magnetic field in the form:

$$A_o(x, y) = \frac{h_o}{4}\left\{ \ln\left[(x-1)^2 + y^2\right] + \ln\left[(x+1)^2 + y^2\right]\right\}. \quad (14)$$

Expanding (14) in series x, y near the point $x = 0$, $y = 0$, we obtain at $r^2 = x^2 + y^2 \ll 1$

$$\begin{aligned} A_o(x, y) &= \frac{h_o}{2}(y^2 - x^2) \\ \vec{H}_o(x, y) &= h_o\left\{y, x, 0\right\} \end{aligned} \quad (15)$$

The point $x = 0$, $y = 0$ is here the zero point of the simplest type (type-X) h_o is the gradient of the absolute value of magnetic field intensity near the neutral point ($h_o = 4 I_o / c$).

We shall consider the gas pressure and the plasma density at the initial state as constant, that is,

$$p(x, y) = p_o = \text{const.} \quad \rho(x, y) = \rho_o = \text{const} \tag{16}$$

The characteristic parameter of the problem is the distance r_s, over which the Alfvén velocity becomes equal to the speed of sound, namely,

$$r_s^2 = \frac{4\pi\rho_o s^2}{h_o^2} = \frac{4\pi\gamma\, n_o \varkappa T}{h_o^2} \tag{17}$$

where γ is adiabate indicator, $\varkappa$ is the Boltzmann constant.
At the same time the square of the Alfvén to sonic velocity ratio, as a function of the distance from the neutral point, is

$$V_A^2 / s^2 = (r / r_s)^2 \tag{18}$$

If the intensity of the magnetic field is sufficiently great while the density of the plasma is small, $r_s \ll 1$ and the condition (9) is fulfilled in a broad region of values

$$r_s \ll r \ll 1 . \tag{19}$$

We shall bring forth as an example the characteristic values for the active regions on the Sun (regions near sunspots). The gradient of the magnetic field near the neutral points reaches the values

$$h_o = 0.1 \frac{\text{gauss}}{\text{km}} = 10^{-6} \frac{\text{gauss}}{\text{cm}} .$$

At concentration of plasma, consisting mainly of hydrogen, $n_o = 10^{11}$ cm^{-3} and temperature $T = 10^4$ °K, we obtain

$$r_s^2 \simeq 3 \cdot 10^{12} \text{ cm}^2,$$

whereas the characteristic distance between spots, taken for the unity, constitutes $a \simeq 10^9 \div 10^{10}$ cm. Hence in dimensionless denotations

$$r_s^2 = 3 \cdot 10^{-6} \div 3 \cdot 10^{-8} \tag{20}$$

The values assumed at computation of r_s are apparently extreme. Thus, the gradient of the magnetic field may be by one order smaller.

But even in such a case the characteristic distance $r_s \simeq 10^7$ cm is found to be significantly smaller than the distance between spots, and consequently, there exists a large interval of values of r satisfying the condition (19).

Assume now that the currents shifted by a small distance δ along the axis x *. Since in the region $r \gg r_s$ the field of shifted currents is quasi-free (see the equation (12)), we may write at once the potential of the displaced field as follows :

$$A(x, y) = \frac{h_0}{4} \left\{ \ln[(x - 1 + \delta)^2 + y^2] + \ln[(x + 1 - \delta)^2 + y^2] \right\} \quad (21)$$

Hence, we have near the neutral point, with a precision to terms of the order δ

$$\begin{aligned} A(x, y) &= \frac{h_0}{2} (y^2 - x^2 - 2\delta) \\ \vec{H}(x, y) &= h_0 \{y, x, 0\} \end{aligned} \quad (22)$$

As may be seen from the comparison of (15) and (22), at such a displacement the field near the neutral point does not vary, while the potential gets a constant addition.

Let us determine now the plasma displacement vector $\vec{\xi}(\vec{r}, t) = \vec{r} - \vec{r}_0$. making use of the equations (6) and (13). Since at considered quasistationary compression the shift $\vec{\xi}$ is composed from infinitely small shifts $\delta\vec{\xi}$ across the lines of force, while by the strength of (15) and (22) the general pattern is invariable in time, the lines of plasma current constitute a family of curves orthogonal to the lines of force, that is, a family of hyperbolae

$$x y = x_0 y_0 \quad (23)$$

Subsequently, as a consequence of the adhesion condition, we have (see the equations (6), (15), (22)) :

$$y^2 - x^2 - 2\delta = y_0^2 - x_0^2 \quad (24)$$

The equations (23), (24) allow the expression of the original coordinates

* It is also easy to take into account the arbitrary variation of current magnitude. Inasmuch as the results do not change in essence, we shall limit ourselves to the indicted simplest case.

x_0, y_0 through the coordinates of the displaced particle and thus determine the vector $\vec{\xi}(\vec{r}) = \{x - x_0,\ y - y_0\}$, namely:

$$x_0^2 = \tfrac{1}{2}\left\{[(y^2 - x^2 - 2\delta)^2 + 4x^2y^2]^{1/2} - (y^2 - x^2 - 2\delta)\right\},$$
$$y_0^2 = \tfrac{1}{2}\left\{[(y^2 - x^2 - 2\delta)^2 + 4x^2y^2]^{1/2} + (y^2 - x^2 - 2\delta)\right\}. \qquad (25)$$

Here we took the positive value of the radical on account of the condition of materiality of the coordinates x_0, y_0.

The displacement determined by the expressions (25) is such that the lines of force, having crossed the axis x at points o, $\sqrt{\delta}$, $\sqrt{2\delta}$, will occupy the place of the lines of force having crossed the axis y respectively at the points $\sqrt{2\delta}$, $\sqrt{\delta}$, 0 (see Fig. 2 in the region $r \gg r_s$).

The obtained displacement pattern of plasma and of the frozen-in magnetic field refer, first of all, only to the region $r \gg r_s$, for which the conditions (7) and (9) are fulfilled. Secondly, regions of strong plasma compression may emerge in the process of displacement. In such regions the conditions (7), (9) may be disrupted and consequently, the expressions (22) and (25) will be inapplicable.

Thirdly, there may appear at field deformation regions with strong current not taken into account in the expression (21). Such currents will distort the shift pattern in the region $r \lesssim \sqrt{\delta}$.

As to regions of strong plasma compression, it is not difficult to establish their existence with the aid of the continuity equation (5). Computing the jacobian of the transformation from coordinates x_0, y_0 to the coordinates x, y, from (5) and (25) it is easy to obtain:

$$\frac{\rho(x, y)}{\rho_0} = \frac{D(x_0, y_0)}{D(x, y)} = \frac{x^2 + y^2}{\sqrt{(x^2 + y^2)^2 + 4\delta\,(x^2 + y^2) + 4\delta^2}} \qquad (26)$$

As follows from (26), in the region

$$y^2 < x^2 + \delta \qquad (27)$$

the shift leads to plasma rarefaction with it being greatest in the region of small r ($r^2 \ll \delta$), when

$$\frac{\rho(x, y)}{\rho_0} \simeq \frac{r^2}{2\delta}. \qquad (28)$$

In the region $y^2 > x^2 + \delta$ the plasma contracts and its density increases; at the same time, at the points

$$x = 0, \quad y = \pm\sqrt{2\delta} \tag{29}$$

the density passes to infinity. In the vicinity of these points the conditions (7), (9) are dusrupted and the true deformation will differ from the one found above. The character of displacements in these regions may be determined utilizing the "freeze-in" condition of the lines of force and taking into account that in the course of the phase of rapid transverse contraction the plasma has no time to shift substantially along the field. In these assumptions the deformation of the magnetic field has the shape indicated in Fig. 2; the regions of strong plasma compression are shown by double shading. Fig. 2 corresponds to current rapprochement ($\delta > 0$). When currents drift away from one another, [$\delta < 0$, see (21), (25)], the pattern differs by the rotation by the angle $\pi/2$. In the latter case the regions of strong plasma densification and the regions of dissipation (see the following point) will lie on the line, connecting the initial currents.

3. - DYNAMIC DISSIPATION OF THE MAGNETIC FIELD

Let us now consider the situation occurring in the nontri ial region $x \lesssim \sqrt{\delta}$. The qualitative field deformation pattern in this region, based on the freeze-in properties and the continuity of the lines of force is drawn in Fig. 2.

The only thing important for the forthcoming is the fact that the entire magnetic flux, having crossed the axis x in the region $0 < x < \sqrt{2\delta}$ (for definiteness we shall consider the right-hand half-plane $x > 0$), that is,

$$\Phi = A_0\,(0, 0) - A_0\,(\sqrt{2\delta}\,, 0) = h_0\delta \tag{30}$$

will be now included inside the band $x \lesssim r_s$ with lines of force creeping along the axis y. That is why, by order of magnitude, the gradient of the deformed field in this region will be

$$h \simeq \frac{H}{r_s} \simeq \frac{\Phi}{r_s^2} = \frac{h_0\delta}{r_s^2} \tag{31}$$

The extension of this region of contracted field (region of Fig. 2

shaded horizontally) along the axis y, taking into account the divergence of the lines of force with the drifting away from the point $y - r_s$, may also be taken equal to r_s.

As may be seen from (31), the field gradient in the region of magnetic contraction rises by δ / r_s^2 times. On the other hand, in the region $y^2 < \delta$ the density of the medium decreases by $y^2/2\delta$ times, as this could be seen above (see (28). Strictly speaking, this result can no longer by applied to the region of magnetic contraction, inasmuch as it was obtained in the assumption of ideal, currentless deformation. It is however easy to see that such an estimate of plasma rarefaction may only overrate the real density in the region of magnetic contraction, inasmuch as the entrapment of the lines of force near the neutral point (in the region $r \lesssim r_s$) only diminishes the gas inflow into the magnetic contraction region. At the same time, the plasma outflow along the axis y as a result of its carrying from the upper to the lower quadrants by the lines of force remains about the same as in the ideal pattern.

Therefore, the characteristic field gradient to plasma concentration ratio in the region of magnetic contraction ($|x| \lesssim r_s$, $r_s \lesssim |y| \ll \sqrt{\delta}$) is by order of magnitude equal to

$$\frac{h}{n} \simeq \frac{h_0}{n_0} \cdot \frac{\delta^2}{r_s^4} \qquad (32)$$

The principle role of this value consists in that it is directly related to the velocity of charge carriers and it allows to establish the absolute criterion of freeze-in disruption *. In reality, by the strength of the quasistationary equation

$$\mathrm{rot}\,\vec{H} = \frac{4\pi}{c}\vec{j} = \frac{4\pi}{c}\sum_i n_i e \vec{v}_i ,$$

determining in the magnetohydrodynamics the density of the electric current, the indicated ratio h/n cannot exceed the obvious limit

$$\frac{h}{n} < 4\pi e \qquad (33)$$

Here $\sum_i n_i = n$ is the charge concentration of both signs in the plasma; at sufficiently high ionization, $\underline{n}$ is approximately equal to the concentration of the plasma.

If the condition (33) is disrupted in any region, this implies that

* see the infrapaginal note next page..

even at relativistic velocities of all charged particles the medium will not ensure the current density required for quasi- stationary preservation of field gradients. In these conditions the assumption, utilized in magnetohydrodynamics, of the smallness of shifting currents is already inapplicable. By the strength of the affected equations

$$\mathrm{rot}\ \vec{H} = \frac{4\pi}{c}\vec{j} + \frac{1}{c}\cdot\frac{\partial\vec{E}}{\partial t}, \quad \mathrm{rot}\,\vec{E} = -\frac{1}{c}\cdot\frac{\partial\vec{H}}{\partial t},$$

there appears an induction electric field E. Its appearance leads to two important consequences.

First, this field is not directly connected with plasma displacement and it may already appear in a system of coordinates that is fixed relative to plasma. This implies the disruption of the freeze-in property, which is fundamental for magnetohydrodynamics and amounting to the condition of absence of electric field in the coordinates linked with the conducting medium. The disrutpion of the freeze-in property implies the possibility of line of force sliding relative to the medium, their interlocking and so forth. It is important that in the case considered these processes are not linked with the Joule dissipation and they may take place in a collisionless plasma.

Secondly, the electric field is directed along the current $\vec{j}$ and thus performs a positive work over the charged particles, increasing their energy. This is precisely the process that assures the transformation of magnetic energy into the kinetic energy of particles, that is, a sort of peculiar magnetic field dissipation which we shall designate as dynamic dissipation. The distinction from the Joule dissipation consists in the first place in that in case of dynamic dissipation there is no simple proprtionality between the current $\bar{j}$ density and the intensity of the electric field E** .

---

* [from the preceding page].- Strictly speaking, the representation of the freeze-in process is already disrupted at any finite conductivity of the medium. The lines of force "move" then relative to the medium with a velocity dependent on medium's conductivity [9]. When the criterion [33] is not fulfilled, this velocity attains its maximum possible value (equal to the speed of light c) and the represnetation of freeze-in is absolutely invalid.

** In the averaged equation of motion of a separate particle $dp/dt = -\beta v + Ee$, where p - is the particle impulse, $-\beta v$ is the dynamic friction force, the inertial term dp/dt, which is neglected in case of Joule dissipation, begins to play an essential role.

At disruption of the condition (33) the current density reaches the saturation while the field energy is expended on the increase of the total relativistic energy of particles $\mathcal{E} = \sqrt{c^2p^2 + m^2c^4}$. The acceleration of particles takes place along the zero lines (along the axis z), forming in the regions of magnetic contraction. Another important circumstance here is that the motion of particles along the neutral lines is steady [17], namely, at particle deflection the magnetic field returns it to the neutral surface, as this is clear also from the direct consideration of the Lorentz force $\frac{e}{c}[\vec{u}\vec{H}]$.

Taking into account the expression (32) we shall write the condition for dynamic dissipation in the form

$$\frac{h_o}{n_o}\cdot\frac{\delta^2}{r_s^4} \gtrsim 4\pi e \qquad (34)$$

The condition (34) is, by its own sense, extreme, requiring the accelerat n of all particles to relativistic energies. In reality, the acceleration of particles, and of electrons in the first place, will take place too at significantly softer conditions, when the dynamic dissipation of the field is linked in essence with the well known phenomenon of escape of particles.

In order to derive the condition, substituting in this case the condition (34), we shall take advantage of the expression for the critical electric field [18]

$$E_{cr} = \frac{4\pi e^3 L n}{\varkappa T},$$

where L is the Coulomb logarithm, T is the temperature of the plasma, and let us estimate the strength of the electric field in the plasma as

$$E = j/\sigma,$$

where $j = \frac{c}{4\pi}h$ and $\sigma = \frac{3\sqrt{3}}{4\pi L e^2\sqrt{m}}\cdot(\varkappa T)^{3/2}$ is the plasma conductivity.

The condition of total particle escape then is

$$\frac{h}{n} = \frac{h_o}{n_o}\cdot\frac{\delta^2}{r_s^4} \gtrsim 4\pi e\, 3\sqrt{3}\left(\frac{\varkappa T}{mc^2}\right)^{1/2} \simeq 4\pi e\left(\frac{T}{2\cdot 10^8}\right)^{1/2}. \qquad (35)$$

As may be seen from the comparison of (34) and (35), the condition for total particle escape at standard temperatures $T \sim 10^4$ o K is substantially weaker than the condition (34).

A more detailed consideration of particle acceleration, maximum acquired ener y and other questions connected with it, will be given elsewhere. Here we shall only pause at the estimate of the total magnetic energy dissipating in the contraction region and of the mean energy of accelerated particles.

The energy, concentrated in the magnetic contraction region, may be estimated, as is clearly seen from the above, by assuming the field to have a cross section of the order $\sim r_s^2$ and be equal to

$$H \simeq \frac{h_0 \delta}{r_s}.$$

Hence

$$W_M = \frac{\overline{H^2}}{8} V \simeq \frac{1}{8} h_0^2 \delta^2 \tag{36}$$

This quantity may be inquisitively compared with the total work performed at current displacement

$$\Delta W_0 = \frac{1}{c} I_0 H_1 \cdot 2\delta \simeq \frac{h_0^2 \delta}{8} \tag{37}$$

where $H_1 = h_0 / 4$ is the field of the first current near the second.

As may be seen from the comparison of (36) and (37), there is concentrated in the contraction region of dimension $\sim r_s$ a fraction δ of the total liberating energy. The remaining energy is expended on gas compression in the region $y \sim \sqrt{2\delta}$ and the deformation of the field in the surrounding space.

The total quantity of gas, carried by the field into the region of magnetic contraction is estimated as follows. All the gas from the sector near the axis x, bouded by the trajectories $y \simeq r_s^2/x$ and $y = 0$ (at $|x| < \sqrt{2\delta}$) is assembled in the region of magnetic contraction. That is why

$$N \simeq 2n_0 \int_{r_s}^{\sqrt{2\delta}} \frac{r_s^2}{x} dx = 2 n_0 r_s^2 \ln \frac{\sqrt{2\delta}}{r_s} \tag{38}$$

At fulfillment of the condition for dynamic dissipation (34) or (35), the energy (36) is expended on the acceleration of particles. By the strength

of (36) and (38) the mean energy per particle after dissipation of the entire energy W_M consititutes

$$\varepsilon \simeq \frac{W_M}{N} = \frac{h_o^2 \delta^2}{16 n_o r_s^2 \ln(\sqrt{2\delta}/r_s)} \tag{39}$$

Let us consider now the temporal characteristics of the process.

The above estimates are valid if during the time of the process, we which we shall denote by τ , the gas dynamical plasma flow along the lines of force lacks the time to equalize the pressure in the considered region. Inasmuch as the quantity of gas in the central region, of $\sim n_o r_s^2$, is small by comparison with the amount of plasma carried into the dissipation region by the magnetic field, we may neglect the inflow from the central region. It is important that the plasma inflow from the region $r \simeq \sqrt{\delta}$ have no time to materialize; hence the necessary condition

$$\tau \ll a \frac{\sqrt{\delta}}{s} \tag{40}$$

This condition implies that the displacement of field sources (basic currents) by the distance δ must take place with a velocity

$$V \simeq \frac{\delta a}{\tau} \gg s\sqrt{\delta} \tag{41}$$

Inasmuch as the motions in the region of a strong field take place with a magnetohdrodynamic velocity $V_A \gg s$ as a rule, the condition (41) may be considered as practically always fulfilled.

4. - DYNAMIC DISSIPATION IN THE REGION OF SOLAR FLARES

Let us apply the results obtained to the solar flare event. We shall make use for the estimates the following values of temperature, magnetic field gradient and of the characteristic distance $\underline{a}$:

$$\begin{aligned} &n_o = 10^{11} cm^{-3} ; \quad T = 10^4\, {}^oK; \quad u = 10^{10}\, cm ; \\ &s = 10^6\, cm/sec ; \quad h_o = 0.1 \frac{gauss}{km} = 10^{-6} \frac{gauss}{cm} \\ &r_s^2 \simeq 3 \cdot 10^{-8} \quad (r_s \simeq 2 \cdot 10^{-4} a = 2 \cdot 10^6\, cm). \end{aligned} \tag{42}$$

It follows from the criterion (34) that the relativistic dynamic

dissipation sets in for the displacements

$$\delta^2 \gtrsim \frac{4\pi e\, n_0 r_s^4}{h_0} \simeq 6 \cdot 10^{-9} \frac{n_0 r_s^4}{h_0} \qquad (43)$$

that is, at utilization of the values (42), for displacements

$$\delta^2 \gtrsim 0.5 \cdot 10^{-6}\ ; \quad \delta \gtrsim 10^{-3} a \simeq 10^7 \text{cm} \qquad (44)$$

By the strength of (40) and (41), the characteristic time of the process and the velocity of spot displacement must be

$$\tau \ll 3 \cdot 10^2 \text{ sec}\ ; \quad v \gg 3 \cdot 10^4 = 0.3 \text{ km/sec} \qquad (45)$$

The shift of the spots by one tenth of a percent of the distance between them and the corresponding times and velocities are quite acceptable from the viewpoint of the available observations. It should then be borne in mind that the utilized values of h and a are most rather extremal, and the characteristic time τ (see (45)) is related to the generation stage of fast particles, while the process of optical flare development may be slower.

For less favorable conditions, when $h \simeq 0.01 \frac{\text{gauss}}{\text{km}}$, the remaining parameters being the same as in (42), we have

$$\delta \gtrsim 0.2\ ; \qquad \tau \ll 5 \quad 10^3\ ; \quad v \gg 5 \cdot 10^5\ . \qquad (46)$$

Let us note here that the total escape of particles (and first of all of electrons) sets in at $T \sim 10^4$ (see (35)) at displacement by about one order lesser.

Therefore, fast electrons are generated in the active regions of the Sun comparatively easily; as is well known, these electrons are responsible for bursts of X- and radio emission. Because of harder conditions, there are, but substantially more seldom, flares where realtivistic particles are generated. For such flares there must exist strong field gradients and substantial spot shifts. The cause of such shifts must be sought for in the instability of spot configuration occurring at their emersion, and in the hydromagnetic waves appearing at such an emersion.

Let us now estimate the energy liberation and the number of accelerated particles for the initial values of (42). From the expressions (36) — (39) we obtain, taking into account that the extension of the region of acceleration along the axiz z is by order of magnitude equal to the distance a between the sources,

— energy passing into fast particles

$$W_M \sim 10^{31} \text{ erg} ;$$

— number of accelerated particles

$$N \sim 10^{34} ;$$

— mean energy of particles

$$\mathcal{E} = \frac{W_M}{N} \simeq 10^{-3} \text{erg} \simeq 10^{9} \text{ ev.}$$

Comparison of the expressions (36) and (37) shows that a comparatively small fraction of energy departs for the acceleration of fast particles. A significant part is expended on plasma compression. It should be noted here, first of all, two regions of strong plasma compression near the points $y = \pm\sqrt{2\delta}$. Such a compression, leading to variation of the ionization equilibrium, may be one of the causes of radiation observed in flares.

At subsequent stages, when gas dynamics processes begins to play a role, which lead to pressure equalization along the lines of force, there occur plasma flows, mostly along the lines of force. These flows may apparently be identified with the observed ejections and prominences.

Thus, the above consideration allows to clarify, at least qualitatively, the conjunction of events taking place at solar flares.

5. - ROLE OF DYNAMIC DISSIPATION IN OTHER EVENTS

In solar flares the field sources, that is, the sunspots, hold on by forces that are external relative to the region under consideration. The gas pressure in denser sub-photospheric layers and the gravitation are such forces. Precisely because of that the requirement (7) my be practically fulfilled for a certain length of time.

The position is different for such objects as supernovae, the exploding galactic nuclei and radiogalaxies. In these objects, the magnetized

plasma is in a state on intense chaotic motion. Here such quasistatic sources of the field, as the Sun is, no longer exist. However, the field irregularities, the densities and the magnetized gas clouds etc.. serve "de facto" as field "sources" at chaotic flow of plasma. The distinction from the conditions near sunspots consists only in that the favorable situation for the dynamic dissipation sets in only casually and briefly in various portions of turbulized plasma. At sufficiently intense motions in such sectors high-energy cosmic rays may be generated. In such a scheme one circumstance is clarified outright, namely that if the dimensions of the system are great and the emergence from it of fast particles takes place slowly by comparison with the velocity (rate) of their generation, there takes place an accumulation of cosmic rays. This accumulation of fast particles may proceed only up to the moment of time, when the pressure $p_{cr} = w_{cr}/3$ is still sufficiently small by comparison with the magnetic pressure $p_H \simeq w_H = \frac{H^2}{8\pi}$. In the opposite case the pressufe of cosmic rays, which must also be taken into account in (7), will stop the above-considered process of dynamic dissipation. Thus, the earlier-adopted assumption [19] about the preservation in the process, of particle escape from the quasi-equilibrium state

$$w_{cr} \sim w_H \sim w_{turb}, \tag{47}$$

is naturally finding in this scheme its justification.

Here, the first part of the correlation stems from the above considerations, and the second (w_{turb}, which is the density of plasma's kinetic energy) stems from magnetohydrodynamics [20].

Let us now pause at concrete application of the criterion (34) to certain objects.

For such events as bursts of supernovae, of galactic nuclei and radiogalaxies we dispose of incomparably poorer information than for solar flares. In the first place this refers to information about the scale of irregularities and the magnitude of field gradients. That is why a direct application of formulas (34) and (35) is beset with difficulties.

Under these conditions we may take advantage of the following ideas. Assuming by the order of magnitude of $h_o = H_o/a$, where H_o is the characteristic value of magnetic field intensity, utilizing the expression (17), and

passing to the usual dimensional quantities, we obtain from (34)

$$H_o^5 \gtrsim \frac{5 \cdot 10^{-38}}{\delta^2} n_o^3 \cdot T^2 a \qquad (48)$$

Note too, that inasmuch as we considered only the given currents and the flow of plasma across the field, there was no necessity above to require that the length of the free path

$$\lambda \simeq 2 \cdot 10^4 \frac{T^2}{n} \text{ cm} \qquad (49)$$

(we assume the Coulomb logarithm to be equal to $L \simeq 20$) be less than all the distances considered. If we now take into account that the displacement of initial currents and their very existence are linked with the dynamics of the gas, we must require that the limitation

$$\lambda \ll a \qquad (50)$$

be fulfilled. Taking into account the feeble dependence of the criterion (48) on a, we may postulate as an indispensable condition $\lambda < a$. At the same time the condition for dynamic dissipation takes the form

$$H_o^5 > \frac{10^{-33}}{\delta^2} n_o^2 T^4. \qquad (51)$$

For subsequent estimates we shall take as still another admissible value

$$\delta^2 \sim 0.1.$$

Then

$$H_o^5 > 10^{-32} n_o^2 T^4. \qquad (52)$$

We shall apply this criterion to various events and objects in which generation of fast particles is actually observed, or may be assumed.

1. - Supernovae. - Crab Nebula.

Let us assume for the nebula supernova shells of value $n_o \lesssim 10$ and $T \simeq 10^4$. For the fulfillment of the criterion (52) it is then required that

$$H_o \gtrsim 2 \cdot 10^{-3} \text{ gauss}, \qquad a \sim \lambda \sim 10^{11} \text{ cm}.$$

Such a value of field strength is in accord with radioastronomical observations. The Crab Nebula is known to be a powerful generator of relativistic particles, with this generation still presently taking place. This last circumstance is evidence that this nebula, or at least its central part, are in a state of intense turbulent motion.

2. - Interstellar Space. - Ionized Regions.

Here $n_0 \simeq 1$, $T \simeq 10^4$ and it would be necessary to have

$$H_0 \gtrsim 6 \cdot 10^{-4} \text{ gauss}, \qquad a \sim \lambda \sim 10^{12} \text{ cm}.$$

In the interstellar space the field is significantly less. The nonfulfilled criterion (52) agrees well with the conclusion of absence of interstellar acceleration [1]. It is easy to ascertain that the criterion (52) is not satisfied for the galactic halo either, when the density is $n_0 \simeq 5 \cdot 10^{-3}$ and the field is $H_0 \simeq 3 \cdot 10^{-6}$ gauss.

3. - Radiogalaxies. - Cygnus A

Assuming for the mean desnity $n_0 \sim 10^{-1} \div 10^{-2}$ and the temperature $T \simeq 10^4$, we obtain

$$H_0 \sim 1 \div 3 \cdot 10^{-4}, \qquad a \sim \lambda \sim 10^{13} \div 10^{14} \text{ cm}.$$

It follows from radioastronomical data that for the radiogalaxy Cygnus A

$$H_0 \simeq 2 \cdot 10^{-4}$$

4. - Exploding Galactic Nuclei and Quasars.

For these objects, the density and mainly the regions where generation of particles takes place, are unknown. Taking according to [21],

$$n_0 \sim 10^5 \div 10^{15}, \quad T \sim 10^4 \div 10^5,$$

we shall obtain

$$H_0 \gtrsim 6 \cdot 10^{-2} \div 4 \cdot 10^3$$

Fields of such magnitude may perfectly emerge as a result of gravitational contraction of the initial magnetized gas cloud (see, for example, [21].

5. - Earth's Magnetosphere and its Magnetic Train

Assuming $n \sim 10$, $T \sim 3 \cdot 10^3$, we obtain the field required for the acceleration of particles

$$H_0 \gtrsim 6 \cdot 10^{-4} = 60\gamma.$$

In the gi en case the charac eristic dimension of a is determined by the flow pattern of solar wind past the Earth, and it may be $\ll \lambda$. The obtained values of H_0 are greater than the values observed in the train [4] which are $H_0 \lesssim 20\gamma$, particularly if we take into account that the temperature of the plasma may be higher: $T \simeq 10^4$ °K. That is why, we cannot expect here the appearance of relativistic particles. Alongside with this, by the strength of the weaker criterion (35), the conditions for total escape of electrons are already realized at $H_0 \gtrsim 20\gamma$. It is quite natural to expect that the observed appearance in the train of fast electrons [5], particularly during magnetosphere disturbances, is conditioned by the dynamic dissipation of the field near the neutral surface in the train. At the same time, the required shift of field "sources" is induced by the perturbation of the train by solar corpuscular streams.

6. - Laboratory Plasma.

For a density $n_0 \sim 10^{15}$ and $T \sim 10^5$, the dynamic dissipation is realized in a field $H_0 \gtrsim 4 \cdot 10^3$ at $a \sim \lambda \sim 0.2$ cm. The observed appearance in strong discharges of neutrons and of hard radiation may be explained by dynamic dissipation of the magnetic field at rapprochement of parallel sectors of the current in an unsteady, strongly twisted string, and also at deflections of discharge's cord from axial symmetry.

* * *

CONCLUSION

As may be seen from the preceding, the dynamic dissipation mechanism constitutes in itself the required link, which allows the connection of the observed rapid changeover of magnetic field with the appearance of accelerated particles in a nonstationary plasma. Given above were only the rough estimates of the basic characteristics of the process; They require refining.

Numerous questions still remain insufficiently clarified. In the first place here belongs the description of motion in the region $\sim r_s$ near the neutral point, and also the description of subsequent stages of the process, which is linked with the solution of quite complex a magnetohydrodynamic problem.

Let us enumerate here the circle of problems that emerge in connection with the dynamic dissipation event :

1) strict solution of the nonstationary magnetohydrodynamic problem;

2) solution of the electrodynamic problem in the region of "freeze-in" disruption;

3) distribution of accelerated particles by energies;

4) accounting of dynamic dissipation in the magnetohydrodynamic turbulence theory;

5) nuclear content of accelerated particles in a multicomponent plasma.

The solution of these questions will require numerous efforts. Nevertheless, the fundamental conclusion of the fact that regions with large field gradients may emerge in nonstationary plasma with a simultaneous high plasma rarefaction (and this may precisely serve as a cause of fast particle generation) will allow, as one may hope, to clarify at this time already a series of questions of magnetohydrodynamics, physics of cosmic rays and of space in general.

*** THE END ***

Contract No. NAS-5-9299
Consultants & Designers, Inc.

Translated by ANDRE L. BRICHANT
on 1 — 3 March 1966

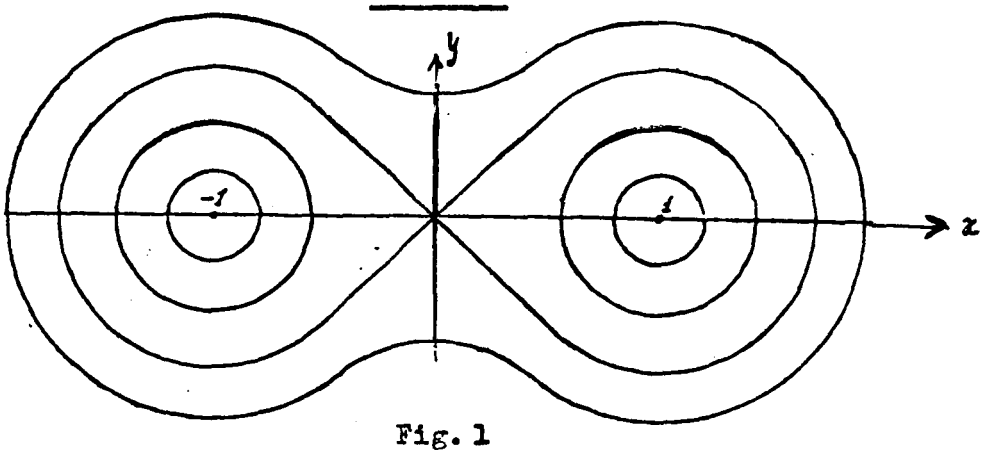

Fig. 1

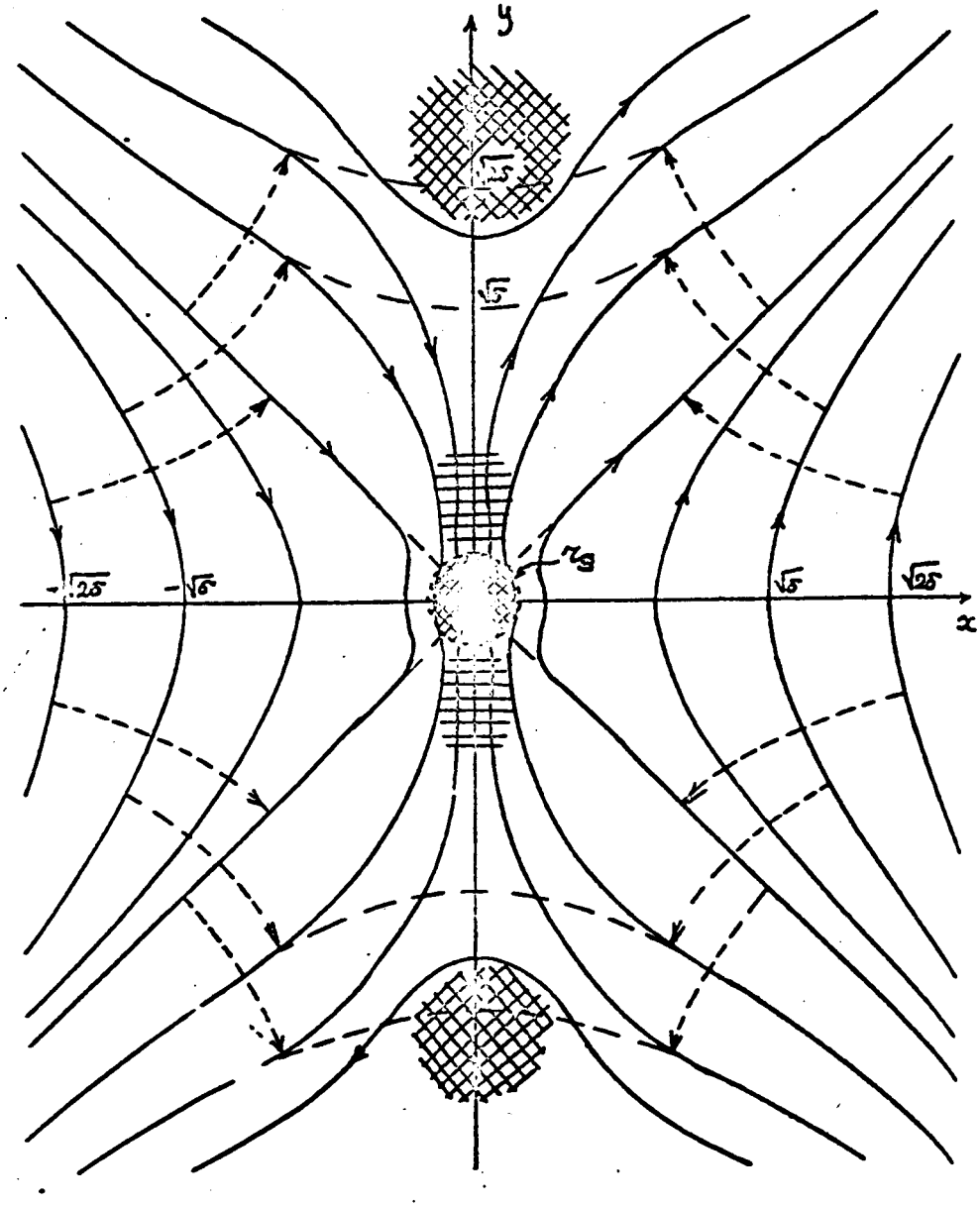

Fig. 2

REFERENCES

1. V. L. GINZBURG, S. I. SYROVATSKIY. Proiskhozhdeniye kosmicheskikh luchey. Izd-vo A. N. SSSR, 1963.
2. F. N. BYRNE, M. A. ELLISON, J. H. REID. Sp. Sci. Rev., 3, 319-41, 1964.
3. Y. G. ROEDERER. Sp. Sci. Rev., 3, 487-511, 1964.
4. N. F. NESS. J. Geophys. Res., 70, 2989, 1965.
5. K. A. ANDERSON. J. Geophys. Res., 70, 4741-63, 1965.
6. L. A. ARTSIMOVICH. Upravlyaemyye termoyadernyye reaktsii. Fizmatgiz, Moskva, 1961.
7. E. N. PARKER. Astrophys. J. Suppl., 8, No. 77, 177-212, 1963.
8. K. DE YAGER. Stroeniye i dinamika atmosfery Solntsa. IL, str. 183-87, 1962.
9. DZH. DANZHI. Kosmicheskaya elektrodinamika. Gosatomizdat, str. 112, 141, 1962.
10. A. B. SEVERNYY. Astron. zh., 35, 335, 1958; Izv. Krymskoy astrofiz. obs. 20, 22, 1958.
11. S. I. SYROVATSKIY. Astron. zh., 39, 987, 1962.
13. L. M. SHUL'MAN. V sb. "Problemy magnitnoy gidrodinamiki i kosmicheskoy gazodinamiki. Izd-vo NAUKA, str. 70, 1964.
12. A. B. SEVERNYY. Astron. zh., 39, 990, 1962.
14. P. A. SWEET. Intern. Astron. Union Symposium, 6, 123-24, 1958; Suppl. Nuovo Cim., 8, No. 2, 188-96, 1958.
15. H. E. PETCHEK. - Annihilation of the Magnetic Field, Preprint, 1964.
16. R. G. GIOVANELLI. - Mo. Not. T. 107, 378, 1947; 108, 163, 1948.
17. T. W. SPEISER. Doklad na mezhdunarodnoy konferentsii po kosmicheskim lucham, London 1965.
18. D. V. SIVUKHIN. V sb. Voprosy teorii plazmy, 4 Atomizdat s. 164, 1964.
19. S. I. SYROVATSKIY. ZHETF, 40, 1788, 1961.
20. S. I. SYROVATSKIY. UFN, 62, 247, 1957
21. V. L. GINZBURG, L. M. OZERNOY. ZhETF, 47, 1030, 1964.

Received for preprinting on 8 Nov. 1965
150 pamphlets

DISTRIBUTION

GODDARD SPACE F.C.

100	CLARK, TOWNSEND
110	STROUD
400	BOURDEAU
610	MEREDITH, SEDDON
611	McDONALD
	ABRAHAM, BOLDT
	GUSS, KNIFFEN
	CLINE, FICHTEL
	VKB, WILLIAMS
	KONRADI, HOFFMAN
	SERLEMITSOS
612	HEPPNER
	NESS
	CAIN
	SUGIURA, SCEARCE
613	KUPPERIAN
	BOGESS III, MILLIGAN
	MELCHER, DONN
614	WHITE
	HALLAM
	FROST, WOLFF
	BEHRING
615	BAUER
	GOLDBERG, HERMAN
	MAIER, STONE
640	HESS
	MEAD, MAEDA
	NORTHROP, NAKADA
	CAMERON, SPEISER
	STERN, JONES
	BURLEY
630	GI for SS [3]
620	SPENCER
252	LIBRARY
256	FREAS

NASA HQS

SS	NEWELL, NAUGLE
	MITCHELL
	SCHARDT, SCHMERLING
	DUBIN
	ROMAN, SMITH
SL	LIDDEL
	BRUNK
	FELLOWS
	HIPSHER-HOROWITZ
SM	FOSTER
	GILL
RR	KURZWEG
RTR	NEILL
ATSS	ROBBINS
WX	SWEET

Ames RC

SONETT
LIBRARY

La R C

160 ADAMSON
185 WEATHERWAX

JPL

SNYDER, NEUGEBAUER
NEWBURN
WYCKOFF

UN. IOWA

VAN ALLEN

UCLA

COLEMAN

SWCAS

JOHNSON