EQUILIBRIUM AND STABILITY OF A z PINCH IN A MULTIPOLE MAGNETIC FIELD

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In the theoretical investigation of the dynamic stabilization of a current-carrying plasma filament by a high-frequency multipole magnetic field it is usually assumed that the cross section of the filament has a circular form in equilibrium [1, 2]. This considerably simplifies the calculations but it does not correspond to reality, since the surface of the plasma must be fluted in the multipole field. An attempt to estimate the influence of the deformation of the filament cross section on its stability against bending in the special case of quadrupole field was made in [3], in which the parameters were determined of the elliptical cross section corresponding to a plasma filament with current in a quadrupole field and an expression was found for the electrodynamic force acting on the filament in the case of long-wavelength kink perturbations. However, this force was calculated incorrectly in [3]. In the present paper a study is made of the equilibrium and stability of a current-carrying plasma filament against kink perturbations in the general case of a multipole stabilizing field. Under the assumption that the flute depth is small, the equilibrium form of the cross section of the current-carrying plasma filament in the multipole magnetic field is found and the components of the force exerted by the field on the perturbed filament are calculated. It is shown that the external field interacts with the current in the perturbed filament only in the case of a quadrupole field. The results are discussed in connection with the problem of multipole dynamical stabilization of a z pinch against kink perturbations.

1. Equilibrium of Current-Carrying Plasma Filament in a Multipole Magnetic Field

We consider a perfectly conducting plasma cylinder with longitudinal surface current I_0 in an external multipole magnetic field. The axis of the conductor coincides with the z axis of a cylindrical coordinate system r, θ, z, and its surface is described by a function $r = r(\theta)$, whose form is determined by the magnetic field.

The magnetic field **B** outside the plasma filament is a superposition of the field $\mathbf{B}_0$ of the current I_0 and the external multipole field $\mathbf{B}_n$, which can be produced, for example, by passing currents I_n in n pairs of linear conductors arranged parallel to the axis of the system, the currents flowing in opposite directions in neighboring conductors. The vector potential A of this field has only the single component $A_z = A_z(r, \theta)$, which satisfies the equation

$$\Delta A_z = 0 \tag{1.1}$$

and appropriate boundary conditions.

If it is assumed, as in [1, 2], that the plasma conductor has circular cross section of radius a, the solution of Eq. (1.1) can be written in the form

$$A_z = A_{z0}(r) + A_{zn}^{\circ}(r, \theta)$$
$$A_{z0} = -A_0 \ln (r / a), \quad A_{zn}^{\bullet} = -A_n [(r / a)^n - (r / a)^{-n}] \cos n\theta \tag{1.2}$$

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Here A_0 and A_n are constants, which can be expressed in terms of the components of the field on the surface of the cylinder. The field $\mathbf{B} = \text{rot}\,\mathbf{A}$ near the cylinder has the components

$$B_r = {}^1/_2 B_n\,[(r/a)^{n-1} - (r/a)^{-n-1}]\sin n\theta$$
$$B_\theta = B_0 (r/a)^{-1} + {}^1/_2 B_n\,[(r/a)^{n-1} + (r/a)^{-n-1}]\cos n\theta \tag{1.3}$$

where $B_n = 2nA_n/a$, $B_0 = A_0/a = 2I_0/ca$, and c is the velocity of light in vacuum.

In the special case when the multipole field is produced by currents I_n in n pairs of linear conductors parallel to the z axis but removed from it by the distance b,

$$B_n = 4n\,(a/b)^n (2I_n/ca) \tag{1.4}$$

Let us consider a more realistic case when the filament section takes the form determined by the condition of constancy of the pressure on the plasma surface. Assuming that the deformation of the filament cross section due to the multipole field is small, we shall look for the equation of the filament surface in the form

$$r(\theta) = a + \delta(\theta) = a(1 - \delta_n \cos n\theta), \quad \delta_n \ll 1 \tag{1.5}$$

The flute depth δ_n is determined by the ratio of the fields $\mathbf{B}_n$ and $\mathbf{B}_0$ near the surface of the conductor. Assuming that

$$|B_n/B_0|\,|_{r=a} \ll 1 \tag{1.6}$$

we shall look for the potential A_z in the form

$$A_z = A_{z0}(r) + A_{zn}^{\delta}(r, \theta)$$
$$A_{z0} = -A_0 \ln(r/a), \quad A_{zn}^{\delta} = -[A_n (r/a)^n + A_{-n}(r/a)^{-n}]\cos n\theta \tag{1.7}$$

[A_0 and A_n are the same as in (1.2)] subject to the additional conditions

$$A_z = A_{z0} + A_{zn}^{\delta} = \text{const} \tag{1.8}$$
$$\mathbf{B}^2 = (\mathbf{B}_0 + \mathbf{B}_n)^2 = \text{const} \tag{1.9}$$

on the surface of the plasma (1.5). The condition (1.8) is a conseqeunce of the translational symmetry of the system; (1.9), of the equilibrium condition of the plasma. Expanding A_z and B^2 in a Taylor series, we obtain in the linear approximation in δ

$$A_z|_{r=a+\delta} \approx (A_{z0} + A_{zn}^{\delta} + \delta\,\partial A_{z0}/\partial r)|_{r=a} = \text{const} = A_{z0}(a) = 0 \tag{1.10}$$

$$\mathbf{B}^2|_{r=a+\delta} \approx (B_{\theta 0}^2 + 2B_{\theta 0} B_{\theta n} + 2B_{\theta 0}\delta\,\partial B_{\theta 0}/\partial r)|_{r=a} = \text{const} = B_{\theta 0}^2(a) = B_0^2 \tag{1.11}$$

Equations (1.10) and (1.11) are obvious if one remembers that the second and third terms in the brackets depend on θ. From (1.10) and (1.11)

$$B_{\theta n} = -\delta\,(\partial B_{\theta 0}/\partial r)|_{r=a}, \quad A_{zn}^{\delta} = -\delta\,(\partial A_{z0}/\partial r)|_{r=a} \tag{1.12}$$

Using (1.7) and the equation $\mathbf{B}_n = \text{rot}\,\mathbf{A}_n$, we find from (1.12) for n > 1 (when n = 1 equilibrium is impossible)

$$A_{-n} = \frac{n+1}{n-1} A_n, \quad \delta_n = \frac{2n}{n-1}\frac{A_n}{A_0}$$
$$A_{zn}^{\delta} = -A_n\left[\left(\frac{r}{a}\right)^n + \frac{n+1}{n-1}\left(\frac{r}{a}\right)^{-n}\right]\cos n\theta \tag{1.13}$$

The magnetic field in the neighborhood of the plasma filament then has the components

$$B_r = \frac{1}{2} B_n\left[\left(\frac{r}{a}\right)^{n-1} + \frac{n+1}{n-1}\left(\frac{r}{a}\right)^{-n-1}\right]\sin n\theta$$
$$B_\theta = B_0\left(\frac{r}{a}\right)^{-1} + \frac{1}{2} B_n\left[\left(\frac{r}{a}\right)^{n-1} - \frac{n+1}{n-1}\left(\frac{r}{a}\right)^{-n-1}\right]\cos n\theta \tag{1.14}$$

Let us consider the behavior of the multipole field near the axis of the system under different conditions. It follows from the expressions (1.3) and (1.14) that when $r = a$

$$B_{\theta n} = {}^1/_2 B_n \cos n\theta, \quad B_{\theta n} = B_n \cos n\theta,$$
$$B_{\theta n} = -[B_n / (n-1)] \cos n\theta$$

in the absence of a plasma, in the case of a circular plasma filament, and in the case of a fluted cross section, respectively.

It can be seen that the multipole field near the axis of the system changes appreciably. In the case of a fluted section even the sign of the field changes, and this must reverse the stabilization effect (destabilization). The change of the multipole field is due to the diamagnetic currents in the plasma conductor induced by the primary external field.

Using the expression (1.13) for δ_n, we can write

$$B_n = (n-1) B_0 \delta_n \tag{1.15}$$

This is a formal equation from which one cannot conclude that $B_n = 0$ when $\delta_n = 0$, since B_n and δ_n were determined above for the two different cases of a circular and fluted section. The relationship between B_n and δ_n arises because for given current I_0 they are determined by the multipole field B_n.

It follows from Eqs. (1.14) and (1.15) that when $r = a$

$$B_{zn} \sim n\delta_n B_0, \quad B_{\theta n} \sim \delta_n B_0, \quad |\mathbf{B}_n| \sim n\delta_n |\mathbf{B}_0|$$

and therefore the condition (1.6) of applicability of the treatment is satisfied when $n\delta_n \ll 1$.

2. Magnetic Field Near the Perburbed Filament

Suppose the plasma conductor undergoes a small perturbation of the $m = 1$ mode under which its surface is described by the equation

$$r(\theta, z, t) = r_0(\theta) + \xi(\theta, z, t) = r_0(t) + a\xi_1(t)\cos(\theta \pm kz), \; \xi_1 \ll 1 \tag{2.1}$$

where $r_0(\theta) = a + \delta(\theta)$ is the equilibrium surface of the conductor and $\xi(\theta, z, t)$ is the displacement of the surface due to the perturbation. The field $\mathbf{B}$ near the perturbed filament is determined from the equations

$$\mathbf{B} = \nabla\Phi, \quad \Delta\Phi = 0 \tag{2.2}$$

with allowance for the decrease of the perturbation of the field at infinity and the boundary condition

$$(\mathbf{Bn}) = 0 \tag{2.3}$$

on the plasma surface (2.1). In (2.3), $\mathbf{n}$ is the unit vector of the external normal to the plasma surface. After expansion in a Taylor series and allowance for the condition $(\mathbf{B}^{(0)}\mathbf{n}^{(0)})|_{r=r_0} = 0$, Eq. (2.3) becomes

$$[\mathbf{B}^{(0)}\mathbf{n}^{(1)} + \mathbf{B}^{(1)}\mathbf{n}^{(0)} + \xi\partial(\mathbf{B}^{(0)}\mathbf{n}^{(0)})/\partial r]|_{r=r_0} = 0 \tag{2.4}$$

Here and below the superscript 0 is appended to the equilibrium quantities, and 1 to their increments in the perturbation. In particular, the fields obtained in Sec. 1 must have the superscript 0.

The vectors $\mathbf{n}^{(0)}$ and $\mathbf{n}^{(1)}$ on the equilibrium surface r_0 have the components (in the linear approximation in δ and ξ)

$$\mathbf{n}^{(0)}\left\{1, -\frac{1}{a}\frac{\partial\delta}{\partial\theta}, -\frac{\partial\delta}{\partial z}\right\} = \mathbf{n}^{(0)}\{1, -n\delta_n \sin n\theta, 0\}$$
$$\mathbf{n}^{(1)}\left\{0, -\frac{1}{r_0}\frac{\partial\xi}{\partial\theta}, -\frac{\partial\xi}{\partial z}\right\} = \mathbf{n}^{(1)}\{0, (1+\delta_n\cos n\theta)\xi_1\sin(\theta-\theta_0), \pm ka\xi_1\sin(\theta-\theta_0)\} \tag{2.5}$$

where $\theta_0 = \pm k\mathbf{z}$ is the angle between the x axis ($\theta = 0$) and the direction of the displacement of the axial line of the filament in the cross section with coordinate z.

Using Eqs. (1.16) and (2.5), we obtain from the boundary condition (2.4) in the linear approximation in δ and ξ

$$-B_0 n\delta_n \xi_1 \sin n\theta \cos(\theta-\theta_0) + B_0\xi_1(1+\delta_n \cos n\theta)\sin(\theta-\theta_0) + B_r^{(1)}|_{r=a} - B_\theta^{(1)}|_{r=a} n\delta_n \sin n\theta = 0 \qquad (2.6)$$

We shall seek the scalar potential $\Phi^{(1)}$ in the form

$$\Phi^{(1)} = \sum_{j=-1}^{1} C_j K_{1+jn}(kr)\sin[(1+jn)\theta-\theta_0] \qquad (2.7)$$

where C_j are constant coefficients and $K_i(x)$ are Macdonald functions. In this form, $\Phi^{(1)}$ satisfies Eq. (2.2) and the boundary condition at infinity. Substituting $B_r^{(1)} = \partial\Phi^{(1)}/\partial r$ and $B_\theta^{(1)} = (1/r)\partial\Phi^{(1)}/\partial\theta$ into (2.6), we obtain

$$B_0\xi_1\sin(\theta-\theta_0) + {}^1/_2 B_0\delta_n\xi_1 \sum_{j=-1,1}(1-jn)\sin[(1+jn)\theta-\theta_0] + \sum_{j=-1}^{1}\Big\{C_j k K'_{1+jn}(ka)\sin[(1+jn)\theta-\theta_0] - C_j a^{-1} K_{1+jn}(ka) n\delta_n \sin n\theta\cos[(1+jn)\theta-\theta_0]\Big\} = 0 \qquad (2.8)$$

The prime denotes differentiation with respect to the argument.

Equating to zero the total coefficients of $\sin[(1+jn)\theta-\theta_0]$ for $j = 0, \pm 1$, we obtain the system of equations

$$B_0\xi_1 + C_0 k K_1'(ka) = 0, \quad j = 0$$
$${}^1/_2(1-jn)B_0\delta_n\xi_1 + C_j k K'_{1+jn}(ka) - {}^1/_2 jn\delta_n C_0 a^{-1} K_1(ka) = 0$$
$$j = \pm 1$$

from which we find the coefficients C_j:

$$C_0 = -\frac{B_0\xi_1}{kK_1'(ka)}, \quad C_{j=\pm1} = -\frac{B_0\delta_n\xi_1}{2kK'_{1+jn}}\left[1 - jn\left(1 - \frac{K_1(ka)}{kaK_1'(ka)}\right)\right] \qquad (2.9)$$

Since the coefficients C_j for $j \neq 0$ are proportional to δ_n, the terms remaining in Eq. (2.8) that contain $\sin n\theta \cos[(1+jn)\theta-\theta_0]$ are of second order in δ_n, which we here ignore. Thus, in the linear approximation in δ, Eq. (2.2) and the boundary conditions at infinity and on the plasma surface can be satisfied by choosing the scalar potential $\Phi^{(1)}$ in the form (2.7) with the coefficients C_j (2.9).

3. Force Acting on a Perturbed Plasma Filament in the Multipole Field

One of the possible ways of investigating the stability of a plasma conductor with a skin current against m = 1 mode perturbations under which the surface of the conductor is described by Eq. (2.1) uses the so-called "model of a flexible filament." In the linear approximation in ξ one calculates the force **F** that the magnetic field exerts on the perturbed conductor and one then investigates the equation of motion of an infinitesimal length of the conductor, $M d^2\xi/dt^2 = \mathbf{F}$, where M is the relevant mass.

As follows from its derivation, this equation actually describes the transverse (relative to the axis) oscillations of thin homogeneous disks that are displaced as a whole and moved relative to one another in the aximuthal direction through the angle $\Delta\theta = \pm k\Delta z$.

This crude model must satisfactorily describe the system if the perturbed plasma behaves as an incompressible fluid and the force F is purely transverse. These conditions are satisfied adequately if, respectively $kv_s \gg \Omega$ and $k\xi \ll 1$, where v_s is the velocity of sound in the plasma and Ω is the frequency of oscillations (growth rate) of the system. The first inequality is the condition of incompressibility, while the second, which means that one must consider only "smooth" perturbations, is obtained from the condition $|n_\perp| \sim 1 \gg |n_\parallel| \sim k\xi$.

The transverse force $F_\perp$ can be found from the formula

$$\mathbf{F}_\perp = -(1/8\pi)\oint \mathbf{B}^2 n_\perp \, dl \qquad (3.1)$$

where the integration is around the bounding contour of the conductor cross section. The contour arc element

$$dl = \{[r(\theta)]^2 + [dr/d\theta]^2\}^{1/2}$$

in the considered case, when $r(\theta)$ is given by Eq. (2.1), is equal in the linear approximation in δ and ξ to

$$dl = [a + \delta + \xi + (1/a)(\delta\xi + \delta'\xi')]d\theta$$

Taking into account the expressions (2.5) for $\mathbf{n}$, one can obtain the components of the vector $n_\perp dl$

$$n_x dl = \Big\{\cos\theta + \xi_1\cos(2\theta-\theta_0) + n\delta_n\xi_1\sin\theta_0\sin n\theta - {}^1/_2\delta_n \sum_{j=-1,1} [(1+jn)\cos(1+jn)\theta + \xi_1\cos\theta\cos[(1+jn)\theta-\theta_0]]\Big\}\, a\, d\theta$$

$$n_y dl = \Big\{\sin\theta + \xi_1\sin(2\theta-\theta_0) - n\delta_n\xi_1\cos\theta_0\sin n\theta - {}^1/_2\delta_n \sum_{j=-1,1} [(1+jn)\sin(1+jn)\theta + \xi_1\sin\theta\cos[(1+jn)\theta-\theta_0]]\Big\}\, a\, d\theta \qquad (3.2)$$

The magnetic pressure on the surface of the perturbed filament in the linear approximation in δ and ξ is

$$(1/8\pi)\mathbf{B}^2|_{r=r_0+\xi} = (1/8\pi)\Big\{\mathbf{B}^{(0)2} + 2\mathbf{B}^{(0)}\mathbf{B}^{(1)} + \xi\partial\mathbf{B}^{(0)2}/\partial r + \delta\frac{\partial}{\partial r}[\mathbf{B}^{(0)2} + 2\mathbf{B}^{(0)}\mathbf{B}^{(1)} + \xi\partial\mathbf{B}^{(0)2}/\partial r]\Big\}\Big|_{r=a}$$

After the calculations we obtain

$$\frac{1}{8\pi}\mathbf{B}^2|_{r=r_0+\xi} = \frac{1}{8\pi}B_0^2\Big\{1 + \sum_{j=-1}^{1} b_j\cos[(1+jn)\theta-\theta_0]\Big\}$$
$$b_0 = 2(\Psi_1 - 1)\xi_1, \quad \Psi_i = -i^2K_i(ka)/kaK_i'(ka) \qquad (3.3)$$
$$b_{j=\pm1} = \{\Psi_1 + n^2 + jn + \Psi_{1+jn}(1+jn)^{-1}[1 - jn(1+\Psi_1)]\}\delta_n\xi_1$$

Analysis of (3.2) and (3.3) shows that the terms with $j = 1$, $n \neq 2$, and also the third and fourth terms in (3.2) do not contribute to the integral (3.1). The expressions for the components of the force $F_\perp$ are

$$F_{x,y} = {}^1/_4\{B_0^2(1-\Psi_1) \pm B_0B_2(\Psi_1^2 + {}^1/_2\Psi_1 - {}^1/_4)\}\xi_{x,y} \qquad (3.4)$$

where $\xi_x = a\xi_1\cos\theta$, $\xi_y = a\xi_1\sin\theta$, and the upper sign in front of the term with B_0B_2 corresponds to the x component, the lower to the y component. In (3.4) we have used the relation $\delta_2 = B_2/B_0$ in accordance with (1.13).

If $ka \ll 1$, then $\Psi_1 \approx 1$, $1-\Psi_1 \approx (ka)^2\ln(2/\eta ka)$ ($\ln\eta = 0.577\ldots$ is Euler's constant) and (3.4) yields for $F_{x,y}$

$$F_{x,y} = {}^1/_4\{B_0^2(ka)^2\ln(2/\eta ka) \pm {}^5/_4B_0B_2\}\xi_{x,y} \qquad (3.5)$$

If the quadrupole field is generated by currents I_2 in two pairs of rods that are connected in antiphase and are situated at a distance b from the axis of the system, then it follows from (1.4) that $B_2 = 8(a/b)^2(2I_2/ca)$ and the second term in the curly brackets of (3.5) takes the form $\pm 40I_0I_2/b^2c^2$. This expression is half that obtained in [3], in which a simplified scheme for calculating F was used.

We give without calculations the result obtained by calculating $F_{x,y}$ for $\delta = 0$. To terms quadratic in B_n, which determine the "minimum B" effect

$$F_{x,y} = {}^1/_4\{B_0^2(1-\Psi_1) \mp B_0B_2\Psi_1 + {}^1/_2B_n^2[1 - {}^1/_2(\Psi_{1-n} + \Psi_{1+n})]\}\xi_{x,y} \qquad (3.6)$$

$$F_{x,y} = {}^1/_4\{B_0^2(ka)^2\ln(2/\eta ka) \mp B_0B_2 + {}^1/_2(1-n)B_n^2\}\xi_{x,y}, \quad ka \ll 1 \qquad (3.7)$$

Comparing (3.4) and (3.5) with (3.6) and (3.7), we can separate out the correction to $F_{x,y}$ linear in B_n due to the fluting of the filament surface:

$$\pm{}^1/_4B_0B_2(\Psi_1^2 + {}^3/_2\Psi_1 - {}^1/_4)\xi_{x,y} \approx \pm{}^9/_{16}B_0B_2\xi_{x,y}, \quad ka \ll 1$$

which exceeds in absolute magnitude by a factor 9/4 the force of the interaction of the quadrupole field with the current I_0 in the case $\delta = 0$ and has the opposite sign.

Inspection of (3.4)-(3.7) shows that the interaction of the multipole field with the current I_0 in the plasma conductor occurs only when $n = 2$ (quadrupole field). The deformation of the surface of the plasma filament under the influence of the quadrupole field leads in the region of the long-wave perturbations to reversal of the effect of the interaction between the quadrupole field and the current in the filament: whereas perturbations in the zx plane (for the field (2.1) considered here) are stabilized and those in the zy plane are destabilized when the filament has a circular cross section [1, 2], the situation is reversed as regards the stabilization of long-wave kink perturbations of a filament with elliptical cross section.

4. Multipole Dynamical Stabilization of a z Pinch

The interaction of the quadrupole field with the current in the plasma filament depends on the direction in which the filament is displaced in the perturbation: in one plane (zy, for example) the filament is stabilized; in the other (zx), destabilized. Periodic variation with the time of the quadrupole field or the current in the filament leads to a dynamic stabilizing effect [1-4], which does not depend on the direction of displacement. For a multipole field of higher order ($n > 2$) a similar effect cannot occur according to the results of Sec. 3.

Experiments on the multipole dynamic stabilization of a z pinch were made with quadrupole [5, 6] and hexagonal ($n = 3$) [7] high-frequency fields. In both cases a stabilizing effect was observed, not only in the region of long-wave perturbations, as was predicted by the theory of [1, 2], but quite generally for kinks of any wavelength. This extension of the range of multipole dynamic stabilization ($n > 2$, short-wave perturbations) indicates that some additional mechanisms must be invoked to explain the stabilization effect (in addition to the interaction between the multipole field and the current in the filament considered here and earlier in [1, 2]). One of them may be associated with rapid oscillations of the plasma surface under the influence of the alternating multipole field. This suggestion is based on the results of [8], in which a general investigation was made of the influence of high-frequency oscillations of the plasma surface on its stability against slowly increasing perturbations.

In the multipole field (1.14) with $B_n = B_{n1} \cos \omega t$ the boundary of the plasma must follow the magnetic "wall" so that the plasma cross section takes the fluted form (1.5), which oscillates on account of (1.15) with the frequency ω of the field if $v_s \gg \omega \delta_n a$. This condition is usually satisfied experimentally, since a relatively small quantity $\sim 10^7$ cm/sec ($\omega \sim 10^7$ sec^{-1}, $\delta_n a \sim 1$ cm) occurs on the right-hand side. Ellipticity of the filament cross section in an alternating quadrupole field was observed in [6]. Rapid oscillations of the filament surface in an alternating multipole field for any n must give a stabilizing effect that is stronger in the region of small-scale short-wave perturbations.

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