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AUTHOR(S): RICHARD A. GERWIN

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University of California

LOS ALAMOS SCIENTIFIC LABORATORY
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GLOBAL PROPERTIES OF OHMICALLY HEATED REVERSED-FIELD PINCHES*

Richard A. Gerwin

Los Alamos Scientific Laboratory
Los Alamos, New Mexico 87544 USA

ABSTRACT

The simultaneous requirements of power balance and pressure balance have been considered. The treatment generalizes the Pease-Braginskii pinch current limit by including toroidal magnetic field, anomalous resistivity, nonradiative losses, and time-dependent fields. The rise of the temperature to a state of power balance proves to be amenable to a very simple and unified description. Finally, the practical parameter windows implied by the joint action of power balance and pressure balance are displayed.

ohmic heating, reversed-field pinch, power balance

INTRODUCTION

Interest in the stabilized Z-pinch (toroidal field within the pinch, poloidal field outside the pinch) was originally motivated by its ideal MHD stability properties In a sharp boundary model, only the kink mode is susceptible to growth, provided beta is less than unity. If the cylindrical pinch is surrounded by a nearby concentric conducting shell, then even the kink mode can be stabilized.[1]

Because of the use of wall stabilization of the kink mode as opposed to the $q > 1$ stabilization condition employed by tokamaks, the important potential advantage of the stabilized Z-pinch appears, namely, that it may be resistively heated to ignition, by its own internal currents. more easily than tokamaks, which are thought to require auxiliary heating.

Diffuse profile considerations show that a monotone-decreasing pitch profile can suppress current-driven modes. The associated magnetic shear also has a stabilizing effect on the pressure-driven modes (Suydam modes). Moreover, this shear probably stabilizes many microinstabilities that might otherwise be worrisome.

A reversal in the toroidal field profile near the edge of the pinch helps to preserve the monotone behavior of the pitch profile and the concomitant desirable magnetic shear; hence the name, Reversed-Field Pinch, or RFP, arose.

More recently, there has been speculation that this reversed field configuration

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arises in a natural way, perhaps through resistive MHD instabilities, or by MHD turbulence.[2a,b]

In this paper, we do not dwell on questions of MHD stability, or details of the profiles. Instead, we assume that reasonable diffuse profiles of the stabilized Z-pinch type are somehow obtained, and we then inquire into the resulting global properties to be expected under the conditions of ohmic heating and power balance. A small amount of reversal in the toroidal field is not expected to alter our general conclusions, nor are other details such as a moderately hollow pressure profile.

Only subignition conditions ($\lesssim$ 5 keV) are considered here, so that alpha-particle heating is largely ignored. A subignition state of power balance must be attained, maintained, and finally augmented by alpha-particle heating, in a progression of experiments leading to a reactor. Thus, the achievement by ohmic heating of a subignition state of power balance in the RFP constitutes an important milestone whose properties need to be displayed and understood. That is the motivation of this paper.

EARLIER WORK

Previous work on power balance in pure Z-pinches has been referenced in a Z-pinch paper by J. Shearer.[3] Extensions to the RFP configuration were made by Butt and Pease,[4] and by Robinson.[5]

This earlier work usually involved exactly steady-state assumptions (vanishing azimuthal electric field, uniform axial electric field) as well as the assumption that the resistivities are classical. In the present paper, these restrictions are lifted. It is important to do so not only because anomalous resistivities may arise in the pinch, due to microturbulence, but also because the desired toroidal field reversal cannot be maintained in a cylindrically symmetric ohmic steady state.

In addition, the temperature rise to the state of global power balance is examined, and the practical parameter windows allowed by the final subignition state are displayed.

A simplifying assumption that allows the global results to be easily obtained is that the temperature profile is uniform. We shall discuss this in more detail later.

Another assumption made here is that the RFP pinch is heated up at constant radius, which is maintained by minor adjustment of the toroidal current. In justification, the pinch must be "fat" (near to the wall) to maintain a semblance of stability to global ideal and resistive modes, so it probably won't have much room to expand.

THE NEED FOR A "LOW-BETA" RFP

The ideal MHD stability of diffuse linear pinches has been considered by D. C. Robinson,[6] who found that RFP profiles can be stable to beta values up to 0.5. The ideal MHD stability of diffuse toroidal RFP configurations has been studied by Baker and Mann,[7] with about the same upper limit on beta.

Since beta is proportional to temperature (at fixed radius), we conclude that $\beta \approx 0.5$ at a reactor temperature of $\gtrsim 5$ keV, say, corresponds to $\beta \approx 0.01$ at a postformation temperature of ~ 100 eV. Thus, for ideal MHD stability reasons, we need to learn how to form very low-beta RFP configurations, in order that they be heatable. (Moreover, impurities should not be relied upon to achieve the low-beta).

Another reason to form low-beta configurations comes from a comparison of time scales for ohmic heating, and resistive evolution and diffusion of the profiles. The former is roughly beta times the latter, so that if $\beta \sim 1$, the profiles will change drastically (and the plasma will diffuse away across the fields) during the time required to heat the plasma. Conversely, if $\beta << 1$, the plasma can be heated by a large factor with scarcely any change in the field or pressure profiles. The latter possibility has indeed been verified by a large transport code.[8] Thus, low-beta allows the plasma to be heated without resistive loss of profiles.

Finally, resistive MHD stability considerations may restrict the maximum beta-values even more than is mentioned. In the following, we assume that the low-beta configuration has been formed, and inquire into its ohmic heating properties.

TIME SCALES

If we combine Ohm's law (in the moving frame) in the form

$$\vec{E}' = \eta_\perp \vec{J}_\perp + \eta_\parallel \vec{J}_\parallel$$

with Ampere's law and Faraday's law, and note that $J_\perp \sim \beta J_\parallel$, then simple arguments can be used to obtain a time for resistive evolution of the profiles. The result is, for each field component, B_z and B_θ,

$$t_\eta \sim a^2/\{D_\parallel \; [1 + O(\beta \; \eta_\perp/\eta_\parallel)]\}, \tag{1}$$

where "a" is the pinch radius, and $D_\parallel \equiv (c^2/4\pi)\eta_\parallel$. In the following we suppose that β is so small that $\beta \; \eta_\perp/\eta_\parallel << 1$. (This does not require $\eta_\perp$ to be classical.) Here the symbols $\perp$ and $\parallel$ are with reference to the local magnetic field direction.

Under these conditions, the ohmic heating will be due primarily to $\eta_\parallel J_\parallel^{\,2}$, and we shall suppose that $\eta_\parallel$ is classical. Then, simple arguments based on the energy equation and Ampere's law yield the characteristic ohmic heating time as

$$t_{O.H.} \sim \beta \; a^2/D_\parallel \; . \tag{2}$$

The use of Ohm's law and pressure balance yields a characteristic time for plasma to diffuse across the magnetic field. This is

$$t_D \sim \beta^{-1} \; a^2/D_\perp \tag{3}$$

where $D_\perp \equiv (c^2/4\pi) \; \eta_\perp$.

Thus, if $\beta \sim 1$, all three times are comparable (even if $\eta_\perp$ is classical), so that ohmic heating is impossible without loss of the profiles and of the plasma. Conversely, if $\beta << 1$, ohmic heating becomes possible even with moderately anomalous values of perpendicular resistivity.

Other times of interest are the electron-ion equipartition time, teq., and the classical cross-field ion thermal conduction time, $t_{\perp i}$. A simple estimate provides

$$\frac{t_{eq.}}{t_{O.H.}} \sim \frac{2}{9}\,\beta^{-1}\,\frac{c^2}{\omega_{pi}^{\;2}\,a^2} \tag{4}$$

where ω_{pi} is the ion plasma frequency. For most conditions of interest, we find that $t_{eq.} \ll t_{O.H.}$ even for very low beta. Thus, it is appropriate to take equal temperatures during the heating phase, $T_e \approx T_i$.

Another simple estimate provides

$$\frac{t_{\perp i}}{t_{O.H.}} \sim \frac{3}{\beta^2}\sqrt{\frac{m_e}{m_i}}\,. \tag{5}$$

For $\beta \ll (m_e/m_i)^{1/4}$, the temperature does not remain uniform during the heating, and the uniform T of our model must be regarded as some average value over the profile. (We assume that the plasma does not contact the wall.) This situation can be mitigated if the postformation temperature profile is uniform. Also, for $\beta \gtrsim 20\%$, the ratio Eq. (5) becomes a value that favors uniformity.

We conclude that the low-beta pinch retains its postformation profiles during the heating phase. The steady state profiles used in earlier work we believe to be irrelevant.

SIMPLE SCALING RESULTS

Relations that describe a state of power balance may be obtained for the RFP in a very simple way except for numerical factors. Consider local power balance (not used in our final treatment) in the form where ohmic heating is balanced by bremsstrahlung and other losses,

$$\eta_{\parallel}\,J_{\parallel}^{\;2} = b_r\,n^2\,T^{1/2} + 3\,nT/t_E\;, \tag{6}$$

at some representative point in the profile, in which the t_E term is meant to include all nonradiative losses, and b_r is the bremsstrahlung constant with T in ergs.

If we set $\eta_{\parallel}^{-1} \equiv S_{\parallel}\,T^{3/2}$, define a characteristic time for energy loss due to bremsstrahlung by $t_b \equiv (3\,T^{1/2})/b_r\,n)$, and use Ampere's law in the crude form $J_{\parallel} \sim (c/4\pi)\,(B/a)$, Eq. (6) can easily be reduced to

$$\frac{c}{\sqrt{S_{\parallel}\,b_r}} = \frac{4\pi a\,n\,T}{B}\sqrt{1 + \frac{t_b}{t_E}}\;. \tag{7}$$

If we ignore the distinction between local beta, $\beta = 16\,\pi\,n\,T/B^2$ and poloidal beta, and bring in the RFP field profiles with the observation that B, the field magnitude, is not very nonuniform, so that $2\pi aB \sim 4\pi I/c$, with I being the total current, then Eq. (7) can immediately be written as

$$\beta \frac{I}{c} = \frac{2A}{\sqrt{1 + t_b/t_E}} \tag{8}$$

where A is a universal constant given by $c/\sqrt{S_{\parallel}\ b_r}$. (We are in c.g.s. units and c is the speed of light.)

Except for a numerical factor, Eq. (8) is the Pease-current result generalized to include axial magnetic fields and nonradiative losses. Note that the perpendicular resistivity does not enter the picture for $\beta \ll 1$, and it may therefore be anomalous without consequence.

The importance of Eq. (8) is that it determines the character of the RFP device, namely, that it probably needs to be a high-current machine.

It is possible to write Eq. (7) or Eq. (8) in several other useful ways. If we square Eq. (7) and define the line density $N \sim \pi a^2 n$ (to within . numerical factor), then we immediately find

$$\beta NT = A^2/(1 + t_b/t_E). \tag{9}$$

Moreover, we can solve for β and write, to within a numerical factor,

$$\beta \left(1 + \frac{t_b}{t_E}\right) = \frac{c^2}{S_{\parallel}\ b_r} \frac{1}{\pi a^2 nT} = \left(\frac{c^2}{S_{\parallel}\ T^{3/2}}\right)\left(\frac{3T^{1/2}}{b_r n}\right) \frac{1}{3\pi a^2}$$

$$= \frac{c^2\ \eta_{\parallel}}{3\pi} \frac{t_b}{a^2} \approx \frac{D_{\parallel}}{a^2} t_b\ , \quad \text{or}$$

$$\beta = \frac{t_b/t_\eta}{1 + t_b/t_E}\ , \tag{10}$$

up to a numerical factor. Thus, the power-balance β-value is directly related to the ratios of the various fundamental time-scales.

Finally, if we multiply the numerator and denominator of the right side of Eq. (7) by $2\pi a$, and use Ampere's law and $N \sim n\pi a^2$, then we find for the total axial current in power balance,

$$\frac{I}{c} = \frac{2NT}{A} \sqrt{1 + \frac{t_b}{t_E}}\ , \tag{11}$$

up to a numerical factor. Here, unlike Eq. (8), the β is not specified.

All of these relations have been verified by a more rigorous treatment, in which, for example, for example, t_b is a well-defined profile-independent constant.

SOME CONSEQUENCES

We shall later present plots of the required features of a power-balanced RFP configuration for given values of beta and temperature, because these are the quantities of primary interest from the standpoints of stability, reactor economics, and approach to ignition.

To see how the magnetic field strength behaves under such conditions, we use Eqs. (8) and (9) to write

$$B \sim \frac{aB}{a} \sim \frac{\beta I}{\beta a} \sim \frac{1}{\beta a\sqrt{1 + t_b/t_E}} \sim \frac{n^{1/2}\,T^{1/2}}{\beta^{1/2}} ,$$

which also follows directly from the definition of β. These dependences are all verified by the numerical results.

More interesting is the scaling of the characteristic ohmic heating time. As discussed later, this time is found to be given by $\beta\, a^2/D_{\|}$ (to within a fixed numerical factor), practically regardless of the ratio t_b/t_E, and in accordance with the simple discussion in the TIME SCALES section. To see how it scales, we use Eq. (9) and write

$$t_{O.H.} \sim \beta \frac{a^2}{D_{\|}} \sim \beta\, a^2\, T^{3/2} \sim (\beta NT)\, \frac{T^{1/2}}{n} \sim \frac{T^{1/2}}{n(1 + t_b/t_E)} .$$

These dependences are also all verified with the more careful theory. We note that, perhaps contrary to off-hand intuition, the higher densities imply shorter heating times to reach a given power balance temperature. Still more surprising, the shorter the t_E-loss time is, the shorter the ohmic heating time to reach a given temperature. These results are all verified by the more rigorous theory, and can be made palatable with the help of Eq. (11).

OUTLINE OF MORE RIGOROUS DERIVATION

We start with "reasonable" postformation profiles, namely parabolic density, $n = n_o\,(1 - r^2/a^2)$, and parabolic axial field, $B_z = B_o\,(1 - r^2/a^2)$, with vacuum for $r > a$. Then $B_\theta(r)$ can be computed from pressure balance. We emphasize that these stabilized Z-pinch profiles are not steady-state profiles since B_z has a reversal at $r = a$. From the fields, the current densities are obtained. Then the ohmic heating can be calculated, and used in the power balance relation.

Pressure balance at $r = a$ proves useful. It can be written as

$$B_\theta^{\,2}\,(a) = 8\pi n_o T + B_o^{\,2}/3 . \qquad (12)$$

If one defines the poloidal beta to be the average plasma pressure divided by the external poloidal field pressure

$$\bar{\beta}_\theta = 8\pi n_o T/B_\theta^2\,(a) = 4NT/(I/c)^2\,, \tag{13}$$

then Eq. (12) provides a relation to the local beta on axis, $\beta_o = 16\pi n_o T/B_o^2$, namely

$$1/\bar{\beta}_\theta = 1 + 2/(3\beta_o). \tag{14}$$

Thus $\bar{\beta}_\theta = 1$ for a pure Z-pinch.

Another useful form of Eq. (12) is

$$I\sqrt{1 - \bar{\beta}_\theta} = (c/\sqrt{3})\ (\Phi_z/\pi a) \tag{15}$$

where Φ_z is the toroidal flux trapped within the pinch during the heating phase. Thus as $\bar{\beta}_\theta$ increases during the heating, the toroidal current should be increased so that the pinch radius, a, stays fixed.

Finally, Eq. (12) can also be written as

$$B_o/B_\theta\,(a) = \sqrt{2\,\bar{\beta}_\theta/\beta_o}\,, \tag{16}$$

which is used later on. The ohmic heating rate per unit length is found from integrating $\eta_\parallel J_\parallel^2$ over the profile. It is found to be

$$\text{O.H.} = g\,(\beta_o)\,D_\parallel\,B_o^2\,, \tag{17}$$

where $g\,(\beta_o)$ is a very complicated function of β_o which is, however, very close to unity for $\beta_o \lesssim 30\%$. For comparison, the pure Z-pinch result is $(\text{O.H.})_z = D_\perp B_\theta^2\,(a)$. In Eq. (17), we have used some average or bulk value of $\eta_\parallel$.

Now we assume global power balance,

$$g\,D_\parallel\,B_o^2 = b_r\,T^{1/2}\int_o^a n^2\,2\pi r dr + 3NT/t_E, \tag{18}$$

with a simple nonradiative energy loss term, with T everywhere in ergs, write $\eta_\parallel^{-1} = 2\,S_\perp\,T^{3/2}$, and make use of Ampere's law, Eq. (14), and Eq. (16), to obtain

$$\bar{\beta}_\theta\,I_{amps} = \left\{\left[\frac{g(\beta_o)}{1 + 1.5\,\beta_o}\right]^{1/2}\frac{60\,c}{\sqrt{2\,b_r\,S_\perp}}\right\}\frac{1}{\sqrt{1 + t_b/t_E}} \tag{19}$$

with

$$t_b = (9/2)\,T^{1/2}/b_r\,n_o\,. \tag{20}$$

The function $g(\beta)/(1 + 1.5\ \beta)$ is very close to 1 for $\beta \lesssim 30\%$. Here, we also used $I_{amps} = 10\ I_{cgs}/c$. The factor in braces is about 2×10^6 amps. If the average parallel resistivity is anomalous, say η_*, one has to replace $S_\perp^{-1/2}$ by $\eta_*^{1/2}\ T^{3/4}$

Equation (19) is the Pease current result we seek.

Combining Eqs. (13) and (19), one finds

$$\bar{\beta}_\theta\ N\ (cm^{-1})\ T\ (keV) = 6.3 \times 10^{18}/(1 + t_b/t_E). \tag{21}$$

Combining Eqs. (19) and (21) yields

$$I_{(amps)} = N\ (cm^{-1})\ T\ (keV)\ \sqrt{1 + t_b/t_E}\ (3.1 \times 10^{-13}). \tag{22}$$

As Eq. (19) determines, in large part, the character of the machine, it is important to understand the foundations upon which it rests. Here, we did not assume steady state conditions, included nonradiative energy losses, and noted that an anomalously large value of perpendicular resistivity need not affect the results provided $\beta^2\ \eta_\perp \ll \eta_\parallel$ and $\beta\ \eta_\perp \ll \eta_\parallel$.

NUMERICAL RESULTS AT POWER BALANCE CONDITIONS

Given values of beta and temperature, and t_b/t_E, Eq. (14) determines poloidal beta, and Eq. (21) determines the line density N. Then the radius, a, is determined as a function of peak density n_o. From the radius, and the current as fixed by Eq. (19), one finds $B_\theta(a)$. Then B_o, the peak axial magnetic field is determined by Eq. (16). As we shall discuss later, the ohmic heating time, $t_{O.H.}$, is given to good approximation by $(3/32)\beta_o\ a^2/D_\parallel$, which is now completely determined. Thus all global quantities of interest are now known and can be graphically displayed.

At this point, two remarks are needed about the ohmic heating. Firstly, a state of power balance really is asymptotic. The $t_{O.H.}$ given here is the time required to reach 95% of the final temperature. Secondly, a careful examination of Poynting's theorem shows that the ohmic heating is taking place, in a low-beta RFP pinch, by depletion of the magnetic energy in and around the pinch. This would occur even if the surrounding discharge tube were fully crowbarred. The point is that the instantaneous power input from the circuit will generally be much less than the ohmic heating power. The point is that we do not have a true steady state with $\nabla \times \vec{E} = 0$. The pinch is slowly resistively evolving during its heating phase.

We now display the power-balance relations between density, radius, and field strength, with given beta and temperature, and ohmic heating time listed along the curves. The figures are self-explanatory. The following qualitative points are emphasized here.

Firstly, with no losses, $t_E = \infty$, it is difficult to envisage an experimental device with densities less than $10^{14}\ cm^{-3}$ because the required minor diameters become impractically large (several meters). This difficulty is exacerbated by alpha-particle energy deposition, see Fig. 6 for t_E negative. On the other hand, it is difficult to envisage running the device with $n_o \to 10^{16}\ cm^{-3}$ because the required

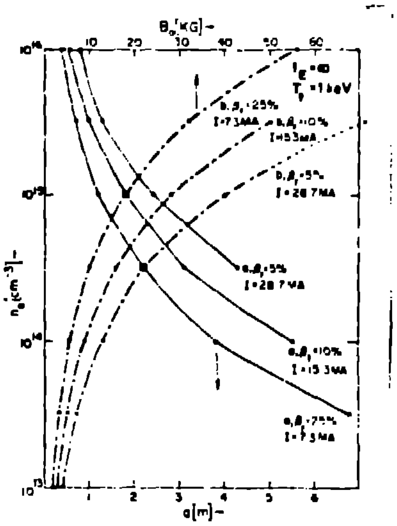

Fig. 1. Density vs radius and field with beta as the parameter and no nonradiative losses.

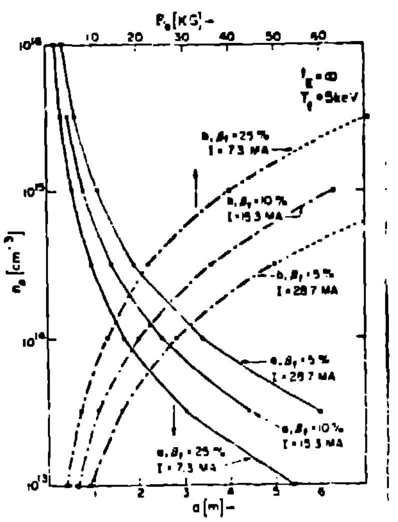

Fig. 2. Density vs radius and field with beta as the parameter and nonradiative losses.

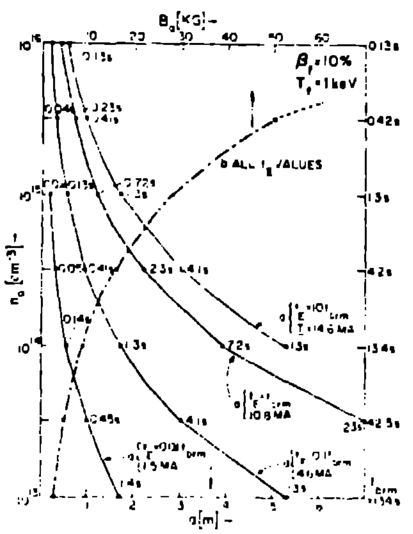

Fig. 3. Density vs radius and field with nonradiative loss as the parameter.

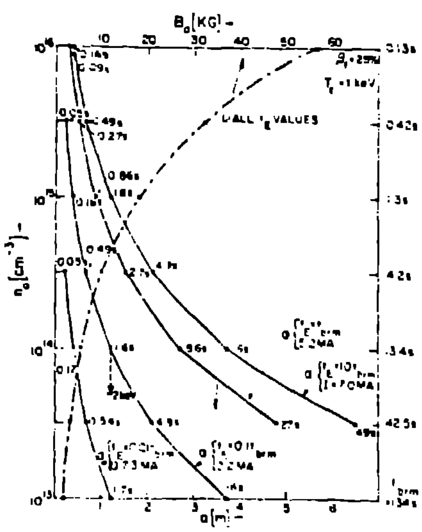

Fig. 4. Density vs radius and field with nonradiative loss as the parameter. Note the use of these figures to obtain a new temperature and density at fixed radius, i.e., the vertical dashed curve.

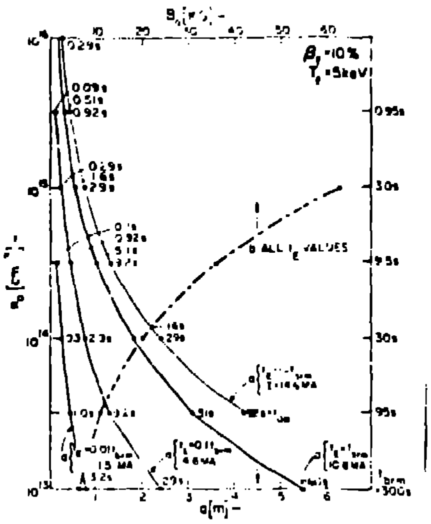

Fig. 5. Density vs. radius and field with nonradiative loss as the parameter.

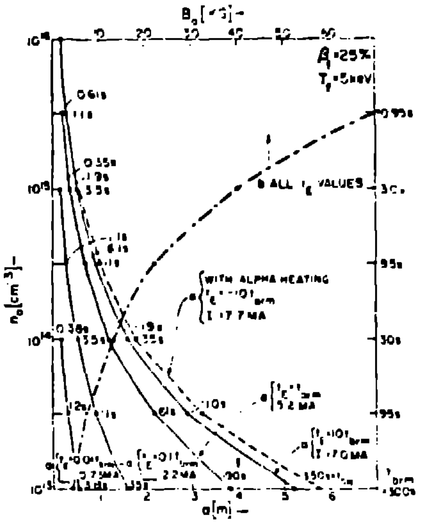

Fig. 6. Density vs. radius and field with nonradiative loss as the parameter. Note the effect of alpha-particle energy deposition as represented by negative t_E.

magnetic field strengths (at a fixed reasonable beta) become impractically large (50 to 100 KG).

Secondly, the addition of nonradiative energy losses mitigates the above-mentioned effects. If one is able to observe a state of global power balance with moderate betas, moderate diameter discharge tubes, moderate field strengths, and moderate toroidal currents, then one can suspect that the nonradiative energy losses are exceeding the bremsstrahlung loss.

THE RISE OF TEMPERATURE TO POWER BALANCE

We write down the equation for energy balance per unit length, where ohmic heating, radiation, and other losses appear on the right,

$$3N \frac{dT}{dt} = g\,(\mathcal{E}_o)\, D_{\parallel}\, B_o^2 - \frac{4}{3\pi a^2}\, b_r\, N^2\, T^{1/2} - \frac{3NT}{t_E}\ , \qquad (23)$$

and in which the line density and radius remain fixed, but $\mathcal{E}_o$ and T are time-dependent. There is no adiabatic compression term in a constant-radius condition.

If we set dT/dt = 0 at T = T_f, the final power-balance temperature, then we can show from Eq. (23), with the help of the various power balance and pressure balance relations, that

$$\beta_{of} = \frac{32}{3} g\ (\beta_{of})\ \frac{t_{bf}/t_{\eta f}}{1 + t_{bf}/t_E} \tag{24}$$

which is equivalent to Eqs. (19) and (21), and is a more precise representation of Eq. (10).

Now if we normalize the temperature T to T_f, and normalize the time t to $t_u = (3/32)\ \beta_{of}\ t_{\eta f}$, set $g(\beta) = 1$, and use Eq. (24), then the equation for the temperature rise can be reduced to

$$\frac{dT}{dt} = T^{-3/2} - \frac{T^{1/2}}{1 + t_b/t_E} - \frac{t_b/t_E}{1 + t_b/t_E} T \ . \tag{25}$$

Numerically obtained solutions to this equation are plotted in Fig. 7. They have been checked analytically in the limits $t_E = 0$ and $t_E = \infty$. It is to be noted that the solutions are insensitive to the ratio t_b/t_E and to the initial condition T_i/T_f. The time required to heat to 95% T_f is always very close to the time unit t_u as defined above.

CONCLUSION

The foundations of the Pease current relation for the RFP have been clarified. An absolute steady state is not required, nor is classical cross field resistivity.

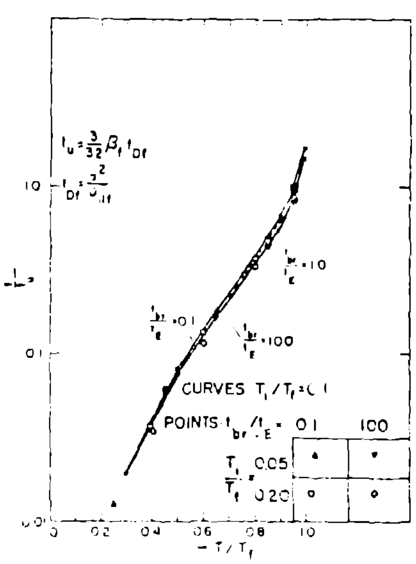

Fig. 7. Normalized time vs. normalized temperature with t_b/t_E and T_i/T_f as parameters.

Moreover, nonradiative energy losses have been included. Eq. (19) has been found to be accurate to within ~20. In comparison to a large transport code, even though the net power flow is still a significant fraction of the heating power.[9]

Finally we showed that the temperature rise due to ohmic heating, properly normalized constitutes an almost universal curve practically independent of loss rates and initial conditions. The final temperature is obtained (i.e., 95%) in a characteristic time that is $(3/32)\,\beta_{0f}\,t_{nf}$, and is insensitive to the parameters characterizing nonradiative losses or initial conditions, thus verifying the scaling given by Eq. (2).

REFERENCES

1. Glasstone, Samuel and Lovberg, Ralph H. (1960). In Controlled Thermonuclear Reactions, D. Van Nostrand & Co., Inc., Prinxcton, NJ.

2a. Montgomery, David, Turner, Leaf, and Vahala, George (1978). Phys. Fluids 21, 757.

2b. Sykes, A. and Wesson, J. A. (1978). Controlled Fusion and Physics (Proc. 8th Europ. Conf., Prague, 1978). p. 80.

3. Shearer, J. W., (1976). Phys. Fluids 19, 1426.

4. Butt, E. P. and Pease, R. S. (1965). Culham Laboratory report CLM-R30.

5. Robinson, D. C. (1978). Preliminary unpublished work presented at the 1978 RFP Workshop in Padua, Italy.

6. Robinson, D.C. (1971). Plasma Physics 13, 439.

7. Baker, D. A. and Mann, L. W. (1972). Proc. 2nd Top. Conf. on Pulsed High Beta Plasmas, Garching 1972, Paper B7, p. 69.

8. Christian, Richard S., private communication (1979).

9. Moses, R. W., Jr., and Nebel, R. private communication (1979).

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