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# BOOTSTRAP STATES OF THE Z-PINCH

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## ABSTRACT

Steady bootstrap states of a Z-pinch are investigated both in absence and in presence of an imposed axial magnetic field, in terms of MHD theory with classical resistivity. The results indicate that bootstrap operation should become possible for certain classes of plasma profiles and that such operation can lead to higher bootstrap currents in a Z-pinch without axial magnetic field than in a tokamak-like case under similar plasma conditions. The ratio between the latter and the former currents is of the order of the square root of the beta value in the tokamak-like case. A simple numerical example is given on boot-strap operation in the Z-pinch.

Neoclassical or anomalous diffusion will increase the diffusion velocity of the plasma but are not expected to affect the main physical features of the present results. This applies also to the kinetic effects in the weak-field region near the axis of the Z-pinch because these effects can largely be described by MHD-like equations for a steady equilibrium.

Bootstrap operation and the technical difficulty in realizing a volume distribution of particle sinks introduce certain constraints on the plasma and current profiles. This has to be taken into account in a stability analysis. The latter cannot only be performed in terms of MHD-like theory but has to be based on kinetic theory including large Larmor radius (LLR) effects.

## 1. Introduction

Steady operation becomes a technical advantage for a large class of magnetic confinement schemes aiming at controlled fusion. Current-drive based on high-frequency oscillations, beam injection, and bootstrap operation are some of the methods which have earlier been proposed for this purpose. In tokamak research current-drive and beam injection have been studied extensively during the last decade. After some recently reported experimental results on tokamaks the bootstrap mechanism may, however, become subject to a renewed interest. Thus, bootstrap currents reaching a substantial fraction of the induced current strength have been measured in JET [1] and TFTR [2]. In multipole devices bootstrap and Pfirsch-Schlüter currents have been reported by Prager [3].

Steady bootstrap operation has also been discussed in connection with the proposed Extrap scheme in which a Z-pinch is stabilized by a multipole field generated by a system of external conductors [4]. In this report the bootstrap equilibrium states of the Z-pinch and related systems will be analysed more in detail. Investigations on the stability of such states are on the other hand outside the frame of this context.

## 2. MHD Analysis of Bootstrap State

The present treatment starts with an MHD analysis of a pinch with an imposed axial magnetic field. The limiting case of a Z-pinch without such a field leads to a zero point at the magnetic axis, and to a corresponding weak-field region within which kinetic analysis becomes necessary. Discussions on the latter are postponed to Section 3.

The bootstrap state to be investigated here is characterized by three main features, namely

- the absence of an induced electric field component along the pinch axis;
- particle losses due to plasma diffusion across the magnetic field, being balanced by a spatial distribution of volume sources, e.g. through pellet injection or ionization of recirculated neutral gas;
- heat losses caused by such processes as diffusion, charge exchange and radiation, being balanced by various intrinsic and externally imposed heating mechanisms including heat producing reactions.

### 2.1. Plasma Balance of Cylindrically Symmetric State

A quasineutral magnetized hydrogen plasma is assumed to be kept in a steady state. The corresponding two-fluid MHD equations for the particle and momentum balance become

$$\mathrm{div}(n\underline{v}_{\nu}) = \rho \qquad (1)$$

$$nm_\nu(\underline{v}_\nu\cdot\underline{\nabla})\underline{v}_\nu = q_\nu n(\underline{E} + \underline{v}_\nu\times\underline{B}) - \underline{\nabla}p_\nu - q_\nu e n^2\eta(\underline{v}_\nu - \underline{v}_\mu) \qquad (2)$$

Here subscript ($\nu$) stands for ions ($\nu$=i) and electrons ($\nu$=e), $n_\nu$=n is the particle density, $\underline{v}_\nu$ the macroscopic fluid velocity, $\rho$ the source intensity of particle production per unit volume, $m_\nu$, $q_i$=e and $q_e$=-e the particle mass and charge, $p_\nu$=p/2=nKT with isotropic temperature $T=T_i=T_e$ of ions and electrons and K denoting Boltzmann's constant, $\underline{E}$ and $\underline{B}$ the electric and magnetic fields, $\eta$ being an isotropic approximation of the resistivity, and $\nu\neq\mu=(i,e)$. In this simplified approach plasma-neutral gas interaction is only included in the source intensity $\rho$ and in the heat balance equation which will not be treated explicitly here. Concerning the momentum balance, restriction is made to the fully ionized rather hot parts of a plasma being impermeable to neutral gas, i.e. where plasma-neutral gas interaction has a minor influence on this balance and on the effective resistivity (compare References [5,6]). The latter is therefore approximated by Spitzer's form $\eta=k_\eta/T^{3/2}$ where $k_\eta=129(\ell n\Lambda)$ with $\Lambda$ being the ratio between the Debye distance and the Coulomb impact parameter. Anomalous and neoclassical diffusion can be roughly simulated by increasing the value of $k_\eta$.

With the substitutions $\underline{v}=(m_i\underline{v}_i+m_e\underline{v}_e)/(m_i+m_e)$ and $\underline{j}=en(\underline{v}_i-\underline{v}_e)$ equations (1)-(2) become expressed in the one-fluid form

$$\mathrm{div}(n\underline{v})=\rho \qquad \mathrm{div}\underline{j}=0 \qquad (3)$$

$$nm(\underline{v}\cdot\underline{\nabla})\underline{v} = \underline{j}\times\underline{B}-\underline{\nabla}p \qquad (4)$$

$$\eta\underline{j} = \underline{E}+\underline{v}\times\underline{B}-(1/2en)\underline{\nabla}p-(m/e)(\underline{v}\cdot\underline{\nabla})\underline{v} \qquad (5)$$

where $m=m_i+m_e$. To this system are added Maxwell's equations

$$\mathrm{curl}\,\underline{B} = \mu_o \underline{j} \tag{6}$$

$$\underline{E} = -\underline{\nabla}\phi \tag{7}$$

In a cylindrically symmetric state, a frame (r, ,z) is introduced with $z$ along the pinch axis and where all field quantities depend on $r$ only. We use expressions (6)-(7) and write $\underline{E}=(-d\phi/dr,0,0)$, $\underline{B}=(0,B,B_o+b)$ where $B_o$ is a constant imposed axial magnetic field, and $\underline{j}=(0,j_\varphi,j_z)$ since $\mathrm{div}\underline{j}=0$ and no current is supposed to be extracted from the pinch surface. In the z-direction equation (4) further yields $dv_z/dr=0$, and since no constant macroscopic fluid motion is assumed in this direction, we have $\underline{v}=(v,0,0)$. Thus, equations (3)-(7) reduce to

$$\frac{d}{dr}(rnv) = r\rho>0 \tag{8}$$

$$nmv\frac{dv}{dr} = (B_o+b)j_\varphi - Bj_z - \frac{dp}{dr} \tag{9}$$

$$\frac{d\phi}{dr} = -(1/2en)\frac{dp}{dr} - (m/e)v\frac{dv}{dr} \tag{10}$$

$$j_\varphi = -(B_o+b)v/\eta = -(1/\mu_o)\frac{db}{dr} \tag{11}$$

$$j_z = Bv/\eta = (1/\mu_o r)\frac{d}{dr}(rB) \tag{12}$$

Here it should first be observed that eq. (10) determines the ambipolar electric field in terms of the plasma profiles and need not be considered in detail in this context. It is further notices that equation (8) can be interpreted as a relation which determines the source intensity $\rho$ in terms of the profile of rnv. The important condition $\rho>0$ has to be imposed here, because volume sources of particles are readily realized by means of injection and ionization, whereas a volume distribution of particle sinks becomes difficult to establish in a hot plasma. It is finally observed that the ratio between the dv/dr and dp/dr terms in equation (9) is of the order of the ratio between the diffusion velocity and the thermal velocity of the ions. This ratio usually becomes small, and the inertia force due to the acceleration of the diffusion velocity represented by the left-hand member of eq. (9) can therefore be neglected with good approximation.

Combination of equations (9), (11) and (12) now yields the bootstrap current in the axial direction

$$j_z = -B(dp/dr)/\left[B^2+(B_o+b)^2\right] \tag{13}$$

From equations (11) and (12) is seen that both currents $j_\varphi$ and $j_z$ are driven by the electric field $\underline{v}\mathrm{x}\underline{B}$ arising from the diffusion velocity v across the magnetic field. Two special cases are of interest here. The first is represented by the axial current

$$j_z^{(Z)} = -(dp/dr)/B^{(Z)} \tag{14}$$

in a Z-pinch without axial field B and without the current $j_\varphi$. In equation (14) the symbol $B^{(Z)}$ represents the poloidal magnetic field component in this case. The second case is represented by a tokamak-like state where $B_o^2 \gg B^2$ and $b^2$, and where the axial current becomes

$$j_z^{(T)} = -B^{(T)}(dp/dr)/B_o^2 \qquad (15)$$

with $B^{(T)}$ denoting the corresponding poloidal field. We now compare a Z-pinch configuration with a tokamak-like configuration having the same dp/dr and the same radius of the plasma cross section. Equations (14) and (15) then yield

$$j_z^{(Z)}/j_z^{(T)} = B_o^2/B^{(Z)} \cdot B^{(T)} \qquad (16)$$

Under the assumption that the pressure profile does not deviate too much from a parabolic shape, equations (12)-(14) can be used to estimate the profiles $B^{(Z)}$ and $B^{(T)}$. Relation (16) then leads to

$$j_z^{(Z)}/j_z^{(T)} \simeq (\beta_T)^{-1/2}; \qquad \beta_T = 4\mu_o n_o K T_o/B_o^2 \qquad (17)$$

where $\beta_T$ stands for the beta value of the tokamak-like case. Usually $\beta_T$ becomes much smaller than unity. It has to be stressed that the results (13)-(17) are independent of the magnitude of the resistivity. Therefore similar results are likely to hold also in the cases of neoclassical and anomalous diffusion. We thus expect the bootstrap current to become larger in a Z-pinch than in tokamak-like systems under similar plasma conditions.

## 2.2. Special Bootstrap States of the Z-Pinch

The following analysis is restricted to a more detailed investigation on the case $B_o$=0 and b=0 of the Z-pinch.

In particular, we shall investigate the constraint imposed by condition (8) on the profiles. For this purpose equations (9) and (12) are combined to

$$\frac{d}{dr}(rB)^2 = -4\mu_o Kr^2 \frac{d}{dr}(nT) \tag{18}$$

$$v = -2(k_\eta K/B^2T^{3/2})\frac{d}{dr}(nT) \tag{19}$$

The condition (8) for a positive source intensity $\rho$ becomes

$$\frac{d}{dr}\left[(r/B^2T^{5/2})\ \frac{d}{dr}(nT)^2\right] \leqslant 0 \tag{20}$$

where $B$ is determined from the profiles of $n$ and $T$ through equation (18).

2.2.1. Some Specific Plasma Profiles

A special class of equilibria is now considered for which the profiles are given by

$$nT = n_oT_o\left[1-(r/r_o)^{2\alpha}\right] \qquad n = n_o\left[1-(r/r_o)^{2\gamma}\right] \qquad (\alpha,\gamma>0) \tag{21}$$

Then equation (18) yields the magnetic field distribution

$$B = \left[4\alpha\mu_o n_o KT_o/(1+\alpha)\right]^{1/2}\cdot(r/r_o)^{\alpha} \tag{22}$$

and an axial current density

$$j_z = [4\alpha(1+\alpha)\mu_o n_o K T_o]^{1/2}\cdot(r/r_o)^{\alpha}\cdot(1/\mu_o r) \tag{23}$$

Equation (19) yields the velocity distribution

$$v=v_{oi}(r_i/r)[1-(r/r_o)^{2\gamma}]^{3/2}/[1-(r/r_o)^{2\alpha}]^{3/2};\quad v_{oi}r_i=(1+\alpha)k_\eta/\mu_o T_o^{3/2} \tag{24}$$

and the corresponding source intensity becomes

$$\begin{aligned} \rho &= n_o v_{oi}(r_i/r^2)F/[1-(r/r_o)^{2\alpha}]^{5/2} \\ F &= [1-(r/r_o)^{2\gamma}]^{3/2}\cdot\{3\alpha(r/r_o)^{2\alpha}[1-(r/r_o)^{2\gamma}]-5\gamma(r/r_o)^{2\gamma}[1-(r/r_o)^{2\alpha}]\} \end{aligned} \tag{25}$$

The condition $\rho \geqslant 0$ is here seen to be satisfied for $F \geqslant 0$. A numerical example which fulfills this requirement will be given at the end of this section.

The present MHD-solution is valid in a region $r > r_i$ the border $r_i$ of which is defined by the local ion radius of curvature being equal to the axial distance $r$. This limiting radius is approximately given by

$$r_i=(2m_i K T_o)^{1/2}/eB(r=r_i)=[(1+\alpha)m_i r_o^{2\alpha}/2e^2\alpha\mu_o n_o]^{1/2(1+\alpha)} \tag{26}$$

### 2.2.2. Alpha Particle Containment

The containment of alpha particles generated by the deuterium-tritium reaction is now considered. These particles have an initial energy corresponding to the potential $\phi_c=3.52\times10^6$V. When they are generated in the plasma core near the axis, the cut-off condition becomes

$$A_c^2 = 2m_p\phi_c/e \tag{27}$$

where $m_p$ is the proton mass and $A_c$ is the difference in magnetic vector potential between the axis and the boundary of the pinch.

From equation (22) we further obtain the total pinch current

$$J_z = 2\pi\left[4\alpha\mu_o n_o KT_o/(1+\alpha)\right]^{1/2} r_o/\mu_o \tag{28}$$

and

$$A_c = \mu_o J_o/2\pi(1+\alpha) \tag{29}$$

The cut-off condition (27) then yields the critical value

$$J_{zc}=(2\pi/\mu_o)(1+\alpha)(2m_p\phi_c/e)^{1/2}=1.36\times10^6(1+\alpha) \tag{30}$$

of the pinch current for alpha particle containment. From equations (28) and (30) we thus have

$$J_z/J_{zc} = \left[2\alpha\mu_o eKn_o T_o/m_p\phi_c(1+\alpha)^3\right]^{1/2} r_o \qquad (31)$$

The obtained result shows that tendencies of inward peaking of the pressure profile by decreasing values of $\rho$ facilitate alpha particle containment. This is understandable because decreasing values of $\alpha$ distribute the magnetic field strength (22) over a larger fraction of the pinch volume. A similar result would apply to more general profile shapes than that given by equation (21). With the special form (21) values of $\alpha$ below 1/2 would lead to infinite pressure gradients at the axis and must be excluded.

### 2.2.3. The Source Intensity and Equivalent Neutral Gas Density

It is here assumed that the source intensity $\rho$ of particle production is due to an ionization rate $\xi$ and an equivalent neutral gas density $n_n$. With plasma temperatures above $10^5$K the ionization rate can be approximated by its maximum value $\xi \simeq \xi_m \simeq 10^{-14}$ $m^3/s$. The concentration ratio of the equivalent neutral gas density then becomes

$$n_n/n = \rho/\xi_m n^2 \qquad (32)$$

This ratio should become much smaller than unity in a fully ionized plasma.

## 3. Considerations of a Kinetic Approach

With the notation $\underline{j}_\nu = q_\nu n \underline{v}_\nu$ equations (2) can be rewritten in the form

$$\underline{j}_\nu = \underline{B} \times [q_\nu n \underline{\nabla} \phi + \underline{\nabla} p_\nu + q_\nu n \eta (\underline{j}_i + \underline{j}_e)] / B^2 \qquad (33)$$

when the inertia force of the left-hand members is neglected. We now introduce the magnetic vector potential $\underline{A}$ defined by $\underline{B} = \mathrm{curl} \underline{A}$ and have $\underline{A} = (0,0,A)$ in the present cylindrically symmetric case of the Z-pinch. Since $A=A(r)$, equations (33) can then be replaced by the equivalent set

$$j_{\nu r} = -q\, n\eta (j_{iz} + j_{ez}) \frac{dr}{dA} \qquad (34)$$

$$j_{\nu z} = q_\nu n \frac{d\phi}{dA} + \frac{dp_\nu}{dA} \qquad (35)$$

because $j_{ir} + j_{er} = j_r = 0$.

In the limit $\eta = 0$ equations (34) and (35) become identical with those deduced by Ågren and Persson [7] from kinetic theory for the equilibrium state of a Z-pinch. As shown by Ågren and Persson, these equations are valid also in the weak-field region near the axis $r=0$, and for non-circular plasma cross sections. The contributions from ions which perform the "Meander-shaped" orbits earlier studied by Hellsten and Tennfors [8] are thus included in these balance equations which should be valid under the approximation of negligible macroscopic fluid accelerations.

For any physically relevant velocity distribution it is always possible to define the current densities $\underline{j}_\nu$ by integration over velocity space. Even in the weak-field region it is therefore plausible that an integrated macroscopic drag force arises between ions and electrons in presence of finite resistivity, as given by the last term of equation (33). It also becomes physically plausible that the bootstrap balance can be extended into the weak-field region $r<r_i$ by equations which yield largely the same result as those used in the MHD study of the previous Section 2. In the case $\alpha \gg 1$, $\gamma \gg 1$ of equations (21)-(26) only the diffusion velocity $v$ is seen to diverge at the axis $r=0$, but from the numerical values to be given at the end of Section 4 this velocity should only reach the thermal limit at very small distances from the axis $r=0$. The present description of the bootstrap state therefore appears to be physically relevant at least concerning its major features.

## 4. A Numerical Example

As a simple numerical example we finally choose $\alpha=1$, $T_o=4\times10^8$ K, $r_o=0.3$ m and $\gamma\to\infty$, $n=n_o=5\times10^{20}m^{-3}$. It is further assumed that the pinch is bounded by a cold-mantle and that the temperature at the interface between the mantle and the fully ionized plasma core is $T_b=10^5$ K. From equation (21) the core radius can then be defined by $T(r=a)=T_b$ which yields $a/r_o=1-5\times10^{-4}$.

These data result in a limiting radius $r_i=0{\cdot}19r_o=0.056$ m, a constant axial current density $j_z=14\times10^6$ A/m$^2$, a total pinch current $J_z=4.0\times10^6$ A, and a magnetic field strength $B(r=a)=2.6$ Tesla at the pinch surface.

The source intensity further becomes positive within the entire plasma core, and it varies from $\rho(r=0)=3.3\times10^{18}$ $m^{-3}{\cdot}s^{-1}$ to $\rho(r=0.95r_o)=3.3\times10^{21}$ $m^{-3}{\cdot}s^{-1}$ in the range $0<r<0.95\ r_o$. Close to the boundary of the core, i.e. when $r$ approaches $a\simeq r_o$, the source intensity would reach the excessively high value $\rho(r=a)=3.3\times10^{42}m^{-3}{\cdot}s^{-1}$ according to equation (25) but we could as well limit the source intensity to $\rho=\rho(0.95r_o)$ in the range $0.95r_o<r<a$ without affecting the equilibrium solution within the grandpart $0<r<0.95r_o$ of the plasma. In a more detailed theory the layers close to $r=a$ will in any case have to be included within a cold-mantle model. According to equation (32) the neutral gas concentration ratio becomes $n_n/n\simeq1.3\times10^{-9}$ at $r=0$.

The classical diffusion velocity finally becomes $v(r_i)\equiv v_{oi}=6.6\times10^{-3}$ m/s at the limiting radius and $v(r=0.95r_o)=0.045$m/s near the boundary. Both these values are much smaller than the ion thermal velocity. According to equation (23) the diffusion velocity would then reach the thermal value $v=3.6\times10^6$ m/s first at the extremely small distance $r\simeq10^{-9}$ m from the axis $r=0$. Even if anomalous diffusion is allowed to increase the values of $v_{oi}$ and $\rho$ by orders of magnitude, this does not seem to lead to physically unacceptable results with the present data. Near $r=0$ there are also negligible electromagnetic forces, thus leading to a pure gas-dynamic plasma balance. Consequently, the simple example given here indicates that physically relevant conditions can be realized for bootstrap operation in the Z-pinch, also under reactor-like conditions. A total pinch current of 4 MA would become sufficient for alpha particle containment. The present example is illustrated by Fig.1.

## 5. Conclusions and Discussions

The simple equilibrium analysis presented here indicates that steady bootstrap operation is at least under certain conditions possible in Extrap and other Z-pinch type systems. Such operation is expected to become more efficient and to lead to larger bootstrap currents in the Z-pinch case than in a tokamak-like case under similar plasma conditions. Moderately large deviations from a circular plasma cross section should not affect the main results. Alpha particle containment appears to be possible.

The analysis has been carried out in terms of classical diffusion based on MHD theory with the Spitzer resistivity. As far as equilibrium bootstrap states is concerned, the general physical features of such states are not expected to become drastically changed by neoclassical or anomalous diffusion. The latter effects can, in a first approximation, be simulated by increasing the resistivity and the diffusion velocity accordingly. Also the kinetic effects in the weak-field region near the pinch axis appear to be describable by MHD-like equations for the bootstrap equilibrium state.

Due to the deduced results. bootstrap operation becomes restricted to certain classes of density, temperature and current profiles. This restriction is caused by the fact that particle sources can be introduced in the plasma volume by various experimental methods, whereas distributed plasma sinks do not seem to be realizable.

The stability of possible bootstrap states and profiles has not been analysed here. Due to the presence of the magnetic weak-field region and the large Larmor radius (LLR) effects of a high-beta pinch state, MHD-analysis of time-dependent instability modes becomes an incomplete and questionable method [9]. A rigorous analysis has to be combined with a supplementary kinetic approach.

## 6. Acknowledgement

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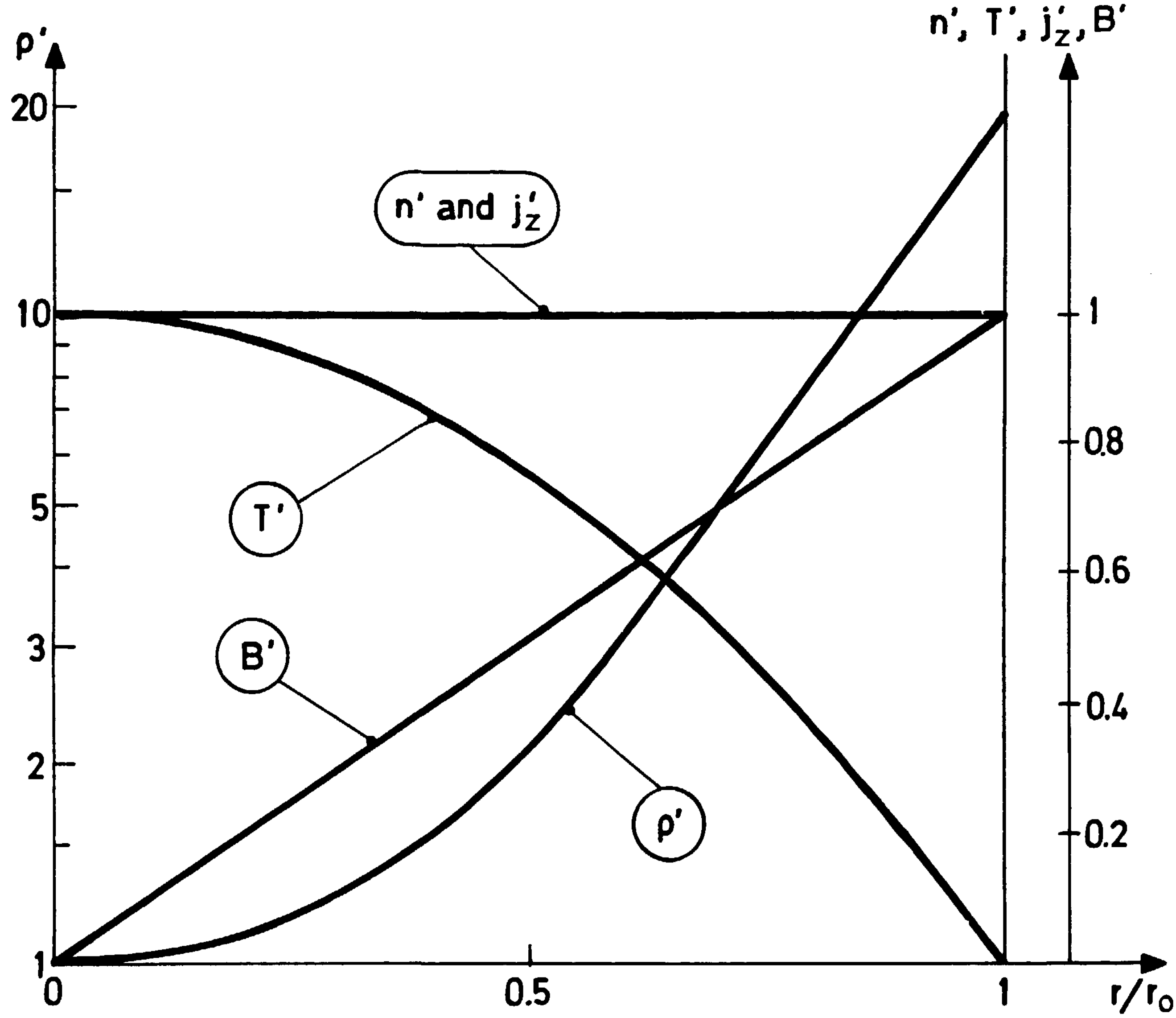
Figure Caption

Fig.1. The numerical example given in Section 4. Figure shows the normalized ion density $n'=n/n(r=0)$, current density $j_z'=j_z/j_z(r=0)$, magnetic field strength $B'=B/B(r=r_0)$, temperature $T'=T/T(r=0)$ and source intensity $\rho'=\rho/\rho(r=0)$ as functions of $r/r_0$ where $r_0 \simeq a$ is the pinch radius.

Fig. 1



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Key words: Magnetized plasma, steady equilibrium.