0460-X

I-17374

THEORY AND SIMULATION OF "FISHBONE"-TYPE INSTABILITIES
IN BEAM-HEATED TOKAMAKS

By

L. Chen et al.

SEPTEMBER 1984

DISCLAIMER

This report was prepared as an account of work sponsored by an agency of the United States Government. Neither the United States Government nor any agency thereof, nor any of their employees, makes any warranty, express or implied, or assumes any legal liability or responsibility for the accuracy, completeness, or usefulness of any information, apparatus, product, or process disclosed, or represents that its use would not infringe privately owned rights. References herein to any specific commercial product, process, or service by trade name, trademark, manufacturer, or otherwise does not necessarily constitute or imply its endorsement, recommendation, or favoring by the United States Government or any agency thereof. The views and opinions of authors expressed herein do not necessarily state or reflect those of the United States Government or any agency thereof.

PLASMA
PHYSICS
LABORATORY

PRINCETON UNIVERSITY
PRINCETON, NEW JERSEY

PREPARED FOR THE U.S. DEPARTMENT OF ENERGY,
UNDER CONTRACT DE-AC02-76-CBO-3073.

DISTRIBUTION OF THIS DOCUMENT IS UNLIMITED

THEORY AND SIMULATION OF "FISHBONE"-TYPE INSTABILITIES

IN BEAM-HEATED TOKAMAKS

PPPL--2136

DESS 002128

by

L. Chen, R. B. White, C. Z. Cheng, F. Romanelli,[†] J. Weiland,^{††} and R. Hay

Plasma Physics Laboratory, Princeton University

P.O. Box 451

Princeton, NJ 08544

J. W. Van Dam, D. C. Barnes, M. N. Rosenbluth, and S. T. Tsai^{†††}

Institute of Fusion Studies, University of Texas

Austin, TX 78712

Abstract

Energetic trapped particles are shown to introduce a new unstable solution to the internal kink and ballooning modes in tokamaks. Both the real frequencies and growth rates of the instabilities are comparable to the trapped-particle precession frequency. Simulations including the excitation and particle-loss mechanisms of the internal kink mode are found to reproduce essential features of the "fishbones". Furthermore, the energetic trapped particle-induced ballooning modes are shown to be consistent with the associated high-frequency oscillations observed experimentally. Several possible stabilizing schemes are considered.

[†]Centro Ricerche Energia Frascati, Rome, Italy

^{††}Institute for Electromagnetic Field Theory and Plasma Physics,
Chalmers University of Technology, Gothenburg, Sweden

^{†††}Institute of Physics, Chinese Academy of Sciences, Beijing, PRC

MASTER

I. INTRODUCTION

The role of energetic trapped particles in tokamaks has recently been dramatized by their significant effect on magnetohydrodynamic (MHD) stability properties such as the maximum achievable β ($=8\pi nT/B^2$) values. In particular, "fishbone" instabilities (~ 20 kHz) observed in tokamaks during high-power, nearly-perpendicular neutral beam injection^[1] are theoretically explained as the internal kink mode excited by trapped beam ions.^[2] The instability mechanism is due to a resonance between the MHD mode and hot trapped particle precessional drifts. This resonance has also been found to cause rapid ejection of the trapped particles.^[3] Combining the instability and ejection processes, we can then simulate the complete "fishbone" event. Since this destabilizing mechanism is rather general, we then apply it to the ideal MHD ballooning mode^[4,5] in order to interpret the high-frequency (~ 100 kHz) oscillations also observed experimentally.^[6] In addition to neutral beam injection, other sources of energetic trapped particles can be considered, such as cyclotron resonant heating and fusion alpha particles.^[7] Finally, highly energetic trapped particles whose precession frequency exceeds the ideal MHD growth rate can enhance stabilization and possibly provide access to second stability in tokamaks.^[8]

In the next section, analytical theory of the "fishbone" internal kink instability is first described. We then present numerical simulations of the full "fishbone" cycle. In Sec. III, trapped particle excitation of the ideal MHD ballooning mode is considered. Several schemes for stabilization are presented in Sec. IV.

II. INTERNAL KINK ("FISHBONE") INSTABILITIES

A. Stability analysis

Consider a large-aspect-ratio ($\epsilon \equiv a/R \ll 1$) tokamak plasma consisting of core (c) and hot (h) components. Consistent with the experimental identification^[1] of the predominantly $m=1$, $n=1$ internal kink perturbations (m and n are, respectively, the poloidal and toroidal mode numbers), we adopt the following formal orderings: $\beta_{pc} \sim 1$, $\beta_{ph} \sim \epsilon$, $T_c/T_h \sim \epsilon^2$ and $|\omega/\omega_A| \sim |\bar{\omega}_{dh}/\omega_A| \sim \epsilon^2$. Here $\bar{\omega}_{dh}$ is the hot trapped particle precession frequency, and $\omega_A = V_A/qR$ with q the safety factor and V_A the Alfvén velocity. The core plasma is treated as an ideal MHD fluid, while the hot component is described by the gyrokinetic equation.

Summing the collisionless equations of motion for each species, we obtain

$$-\omega^2 \rho_m \xi = (\delta \mathcal{L} \times \mathcal{B} + \mathcal{L} \times \delta \mathcal{B})/c - \nabla \delta p_c - \nabla \cdot \delta \mathcal{P}_h \quad (1)$$

where ξ is the fluid displacement, $\rho_m = N_{01} M_1$, $\delta \mathcal{B} = \mathcal{L} \times (\xi \times \mathcal{B})$, $\delta \mathcal{L} = c \mathcal{L} \times \delta \mathcal{B} / 4\pi$, and $\delta p_c = -[\xi \cdot \nabla p_c + \gamma p_c (\mathcal{L} \cdot \xi)]$. Assuming ω is much less than the hot-particle transit and bounce frequencies, the hot particle pressure tensor $\delta \mathcal{P}_h$ is given by

$$\delta \mathcal{P}_h = -\xi_{\perp} \cdot \nabla [P_{\perp} \mathcal{I} + (P_{\parallel} - P_{\perp}) \mathcal{E}_{\parallel} \mathcal{E}_{\parallel}]_h + \delta \hat{P}_{\perp} \mathcal{I} + (\delta \hat{P}_{\parallel} - \delta \hat{P}_{\perp}) \mathcal{E}_{\parallel} \mathcal{E}_{\parallel},$$

where

$$\left\{ \begin{array}{l} \delta p_{\perp} \\ \delta p_{\parallel} \end{array} \right\} = 2^{7/2} \pi M_h B \int_{B_{\max}}^{B^{-1}} d\alpha (1-\alpha B)^{1/2} \int_0^{\infty} dE \frac{E^{5/2} Q}{\omega - \bar{\omega}_{dh}} \bar{J} \left\{ \begin{array}{l} \alpha B/2 (1-\alpha B) \\ 1 \end{array} \right\}, \quad (2)$$

$\bar{a}_{\parallel} = B/B$, $\alpha B = v_{\perp}^2/v^2$, $E = v^2/2$, $\bar{A} \equiv (\oint Ad\ell/|v_{\parallel}|)/\oint d\ell/|v_{\parallel}|$ denotes bounce averaging, $Q = (\omega \partial/\partial E + \omega_{*h}) F_{oh}$, $\omega_{*h} = -(i/\omega_c) (\bar{a}_{\parallel} \times \nabla \ln F_{oh})$, F_{oh} is the hot particle distribution function, $\bar{J} = (\alpha B/2) \nabla \cdot \xi_{\perp} - (1-3\alpha B/2) \xi_{\perp} \cdot \xi$, and $\xi = \bar{a}_{\parallel} \cdot \nabla \bar{a}_{\parallel}$. Equations (1) and (2) constitute a complete normal-mode description in terms of ξ .

Assuming a fixed conducting boundary, we can derive a dispersion relation variationally. The dispersion functional (quadratic form) is $D[\xi] = \delta W_{MHD} + \delta W_k + \delta I$. Here δW_{MHD} is the usual ideal MHD δW involving the total fluid pressure $P = P_c + (P_{\parallel} + P_{\perp})_h/2$; $\delta I = -(\omega^2/2) \int d^3x \rho_m |\xi|^2$ is the inertial energy, and

$$\delta W_k = -2^{9/2} \pi^3 M_h \int_{RB} r dr \int_{B_{\max}}^{B_{\min}^{-1}} d\alpha \int_0^{\infty} dE E^{5/2} K_b \bar{J}^* \frac{Q}{\omega - \bar{\omega}_{dh}} \bar{J} \quad (3)$$

with $K_b = \oint (d\theta/2\pi) (1 - \alpha B)^{-1/2}$. Note that $\delta W_k^{(2)} \sim \epsilon^2 (B^2 |\xi/R|^2 V) \sim \delta I^{(2)}$, where V is the volume and superscripts denote orderings. The variational scheme is then to find a ξ_k which minimizes D to $O(\epsilon^3)$. This minimizing procedure is facilitated by separating the regions outside and inside the singular surface located at $q_s = q(r_s) = 1$. For circular cross sections, we obtain the following dispersion relation^[2]:

$$\frac{-i\omega}{\omega_A} + \delta\hat{w}_{fc} + \delta\hat{w}_k = 0, \quad (4)$$

where

$\hat{w}_A = v_A / [(1 + 2q_s^2)^{1/2} R_0 s]$, $s = r_s q' (r_s) / q_s$, and $\delta\hat{w}_{fc} = \pi (r_s / R_0)^2 \times \delta\tilde{w}_{TC}$. Also $\delta\tilde{w}_{TC}$ has been given previously^[9] and contains only the core-plasma contribution, and

$$\delta\hat{w}_k = \frac{2^{3/2} M_h \pi^2}{B_0^2} \left\langle \int_0^{1+r/R} \frac{d(\alpha B)}{1-r/R} \int_0^\infty dE E^{5/2} \frac{K_2^2}{K_b} \frac{\omega}{\bar{\omega}_{dh} - \omega} \left(\frac{\partial}{\partial E} + \frac{\bar{\omega}_{*h}}{\bar{\omega}_{dh}} \right) F_{oh} \right\rangle \quad (5)$$

Here $\langle y \rangle \equiv (2/r_s^2) \int_0^{r_s} y \, r dr$ is the average of y within the $q = 1$ surface, $K_2 = \phi(d\theta/2\pi) \cos\theta (1 - \alpha B)^{-1/2}$, and ($m = 1$, $n = 1$) perturbations are assumed.

The stability properties are modified qualitatively by the fact that $\delta\hat{w}_k$ contains the mode-precessional drift resonance. A Nyquist analysis reveals the following general characteristics which are relatively insensitive to the F_{oh} one assumes: (i) There exist two branches of the solution. One is the usual MHD branch driven by $\delta\hat{w}_{fc}$, and the other is a new branch induced by the energetic trapped particle. (ii) Even when the MHD branch is stable ($\delta\hat{w}_{fc} > 0$), the trapped-particle branch can become unstable if $\langle \beta_{h,t} \hat{i} \rangle$ exceeds $\langle \beta_{h,t} \hat{i} \rangle_{crit}$ where

$$\langle \beta_{h,t} \hat{I} \rangle_{\text{crit}} = \begin{cases} \bar{\omega}_{dh}/\hat{\omega}_A & \text{for } \hat{\omega}_{fc} < \bar{\omega}_{dh}/\hat{\omega}_A \\ \hat{\omega}_{fc} & \text{for } \hat{\omega}_{fc} > \bar{\omega}_{dh}/\hat{\omega}_A \end{cases}, \quad (6)$$

Here $\hat{I} \sim (K_2/\epsilon K_D)$ is a trapped particle geometrical factor typically of $O(1/\epsilon)$ and, thus, $\langle \beta_{h,t} \rangle_{\text{crit}} \sim O(10^{-3}-10^{-2})$. (iii) Both the real frequencies and growth rates of this trapped particle-induced branch are typically comparable to $\bar{\omega}_{dh}$.

The above features are consistent with the experimental observations. Specific calculations for slowing-down and Maxwellian F_{oh} have been performed. [2,7]

B. Simulation of the "fishbone" cycle.

MHD perturbations with ω in resonance with $\bar{\omega}_{dh}$ can cause secular radial drift of energetic trapped particles. To see this, consider a perturbation δB_r rotating toroidally with $\chi = \omega R \frac{r}{\phi}$. In the frame moving with the perturbation there is an electric field $\delta E = \chi \times \delta B$, producing a δE_ϕ . This in turn produces a $\delta E_\phi \times B$ radial drift of particles. If the trapped particles and the perturbation remain in phase ($\omega = \bar{\omega}_{dh}$), this motion is secular. A Monte-Carlo code based on a Hamiltonian guiding center drift orbit formalism [3] has been used to study the resonant loss process and to simulate the full "fishbone" cycle. The mode structure was obtained from a nonlinear MHD initial value code which reproduced the two principal diagnostic signals, due to soft X-ray detectors and Mirnov coils. The dispersion relation was calculated numerically using the instantaneous particle distribution. The initial particle distributions were given by beam injection and evolution

codes.

The results of a typical simulation using PDX parameters are shown in Fig. 1. As soon as $\langle \beta_{h,t} \rangle$ exceeds $\langle \beta_{h,t} \rangle_{crit}$ the mode begins to grow. At some point the mode reaches a significant amplitude ($\delta B/B \approx 10^{-4}$) and begins to affect particle motion, ejecting particles in a beam rotating toroidally with frequency $\omega \approx \langle \bar{\omega}_{dh} \rangle$. This loss process is rapid, and $\langle \beta_{h,t} \rangle$ drops to $\langle \beta_{h,t} \rangle_{crit}$ typically in hundreds of microseconds. The loss begins with particles of highest energy, and since $\bar{\omega}_{dh} \sim E$, this produces a whistling downward of the "fishbone" frequency (~ 20 kHz for PDX) during the loss process, which is also observed experimentally. The maximum mode amplitude is reached when $\langle \beta_{h,t} \rangle$ has dropped to $\langle \beta_{h,t} \rangle_{crit}$ ($Im\omega=0$), and the mode continues to eject particles, driving $\langle \beta_{h,t} \rangle$ well below $\langle \beta_{h,t} \rangle_{crit}$. The mode then decays to a point where it no longer affects particle motion, and the cycle begins over. The beam deposition rate is typically much lower than the loss rate, so the time between bursts of MHD activity is longer than the bursts themselves.

Growth rates, frequencies, and other properties of the "fishbone" cycle depend on the beam injection parameters, and the nature of the equilibrium. In particular, high shear can make the particle precession rate a strong function of position, limiting resonance with the mode as a particle moves outwards.

IV. TRAPPED PARTICLE-INDUCED MHD BALLOONING INSTABILITY

By analogy with the theory of "fishbones," one may expect that energetic trapped particles can also destabilize ideal MHD ballooning modes by introducing a new unstable branch at $|\omega| \approx \bar{\omega}_{dh}$. The major difference is that whereas the $m=1=n$ internal kink-"fishbones" involve radially global

quantities, the high-mode-number ballooning instabilities will be localized near a flux surface.

As before, adopt an MHD fluid description for the warm core plasma but use the gyrokinetic equation for the hot component where $\omega \sim \bar{\omega}_{dh}$. With the orderings $\epsilon\beta_{pc} \sim 1$, $\beta_{ph} \sim 1$, and $T_c/T_h \sim \epsilon^2$, the ballooning eigenmode equation can be derived^[4] as

$$\frac{d}{d\chi} (1 + s^2 \chi^2) \frac{d\phi}{d\chi} + \Omega^2 (1 + s^2 \chi^2) \phi + g(\chi) [\epsilon\beta_{pc} \phi + \epsilon\beta_{ph} \phi] = 0 \quad (7)$$

where $s = (r/q) dq/dr$, $\Omega = \omega/\hat{\omega}_A$, ϕ is the perturbed electrostatic potential, χ is the extended poloidal angle, and $g(\chi) = \cos\chi + (s\chi - \Delta_c \sin\chi) \sin\chi - \epsilon(1 - q^{-2})$. The trapped particle term is

$$\epsilon\beta_{ph} \phi = \frac{2^{3/2} \pi^2 q^2 M_h}{B^2} \int_{B_{\max}^{-1}}^{B^{-1}} \frac{B d\alpha}{(1 - \alpha_B)^{1/2}} \int_0^\infty \frac{dE E^{5/2}}{(\omega - \bar{\omega}_{dh})} Q \overline{g(\chi) \phi}, \quad (8)$$

with notations the same as in Sec. II. If small shear, $\hat{s} < 1$, is assumed, Eq. (8) can be analyzed in terms of the short and long spatial scales, χ and $z = \hat{s}\chi$. Introducing the subsidiary ordering $\Omega^2 \sim \Delta_c^4 \sim \Delta_h^4 \sim \epsilon(1 - q^{-2}) \sim \hat{s}^2$ and then expanding and solving for $\phi = \phi_0(z) + \phi_1(\chi, z) + \dots$ up to fourth order in powers of $\hat{s}^{1/2}$, one obtains an averaged ballooning equation for the function

$\psi(z) = \phi_0(z) (1+z^2)^{1/2}$. Now match the asymptotic solution $\psi \sim \exp [i (\Omega/\hat{s}) |z|]$ for $|z| \gg 1$ with the logarithmic derivative of the inner solution for $|z| \sim 1$, assuming $|\Omega/\hat{s}| < 1$ and constant- ψ in the inner region. This yields a dispersion relation of the same form as that of the internal kink mode, Eq. (4), with appropriate replacements of $\delta \hat{W}_{fc}$ and $\langle \beta_{h,t} \hat{I} \rangle$ in terms of surface quantities. The stability properties are, thus, formally similar to those discussed in Sec. II. Note that the $\hat{s} < 1$ limitation on the preceding analysis can be somewhat relaxed if the constant- ψ approximation is replaced by a weakly varying trial function which provides better estimates of $\delta \hat{W}_{fc}$ and $\langle \beta_{h,t} \hat{I} \rangle$ for $\hat{s} \lesssim 1$. Otherwise, the stability properties remain unchanged.

The restrictions on $\hat{s}$ and β can be further relaxed in the following way. [5] Assume that the energetic particles are deeply trapped so that the $\epsilon \beta_{ph} \phi$ term in Eq. (7) is localized near $\chi=0$, where it couples the zero-frequency even and odd solutions ϕ_e and ϕ_o . For $\hat{s} \chi_s > 1$, each has the asymptotic form $\phi_{e0} \sim M_{e,0} \chi^\nu + N_{e,0} \chi^{-\nu}$, where the four Mercier coefficients $M_{e,0}$ and $N_{e,0}$, and the exponent ν can be computed numerically for a given equilibrium. With this information, the total solution $\phi_{in} = \phi_e(\chi) - \phi_o(\chi)$ [$d\chi \epsilon \beta_{ph} g(\chi)$ can be matched to the small argument form of the Hankel function solution $\phi_{out} = \chi^{1/2} H_\nu^{(1)}(\Omega \chi)$ in the outer, inertial region where $\Omega \chi \sim 1$. This procedure yields a complex dispersion relation, which can be analyzed by the Nyquist technique.

Figure 2 presents the results of such an analysis for a large-aspect-ratio model equilibrium. The trapped particle-induced branch is unstable for trapped particle beta values, $\epsilon \beta_{ph}$, to the right of the solid curves for various $\bar{w}_{dh}/\omega_A$ values. Note that these curves approximately have the generic features given in Eq. (6). For $\bar{w}_{dh}/\omega_A < 1$, the trapped particle branch restricts stability significantly. However, at very high energies, we can

have $\bar{\omega}_{dh}/\omega_A \gg 1$. MHD ballooning can then be analyzed^[8] in the limit where $|\omega|/\bar{\omega}_{dh}$ is negligibly small, i.e., the hot particles behave so as to conserve the flux adiabatic invariant, which leads to enhanced compressional stabilization. The dashed curve in Fig. 2 is the marginal stability boundary for this zero-frequency MHD mode. Thus, for large $\bar{\omega}_{dh}/\omega_A$, the introduction of highly energetic particles is able to provide direct stable access to the second stability regime at high core plasma beta values, $\epsilon\beta_{pc}$. In general, however, ballooning stability in the presence of hot trapped particles requires that both the zero- and precession-frequency unstable branches be avoided.

This trapped particle-induced ballooning instability appears to be a likely candidate to explain the high-frequency (50 ~ 150 KHz) MHD-type oscillations observed in the PDX tokamak during high-power perpendicular injection,^[6] often as a "precursor" to the "fishbone" oscillations: (i) The ratio of the "precursor" frequency to the frequency of the n=1 "fishbones" is close to the measured toroidal mode number $n \approx 2-6$ of the "precursor." Theory predicts $\omega_r \approx \hat{\omega}_{dm}$, in agreement with experiment. (ii) While the high-frequency oscillations tend to occur at high values of beam power and $\epsilon\beta_p$, the "fishbones" can occur at lower values. Likewise, theory finds a higher $\langle\beta_{h,t}\rangle$ instability threshold than that for the "fishbones." (iii) Experimentally, the high-frequency oscillations are localized near the q=1 surface. Since the q profile flattens there, shear $\hat{s}$ is reduced along with the threshold β value, consistent with theory. (iv) When the high-frequency oscillations are present along with "fishbones," they occur in the beginning stage of the "fishbones." Since q on axis drops during heating, ballooning modes, which require $q(0) \gg 1$, would occur before the internal kink instability sets in at $q(0) < 1$. (v) In cases where only the high-frequency oscillations are

observed, sawteething generally is either absent or has small amplitude, which is consistent with the idea that ballooning modes occur with $q(0) \gtrsim 1$.

IV. STABILIZATION SCHEMES

We have demonstrated that "fishbones" and their high-frequency "precursors" are excited via the resonant coupling between the curvature drift precessional motion of energetic trapped particles and ideal MHD modes (internal kink and ballooning, respectively) of the core plasma. To suppress instabilities of this type and the associated loss of hot ions, we propose the following stabilization schemes: (i) Apply either tangential or not directly perpendicular neutral beam heating, in order to reduce the trapped hot ion beta value below its threshold value; (ii) Employ shaping of the plasma cross section, e.g., bean shape^[10], to make δW_{fc} more positive (i.e., raise the $\langle \beta_{h,t} \rangle$ threshold) and make the MHD branch even more stable; (iii) Introduce highly energetic ($\sim$ MeV) beams,^[8] in order to raise the threshold trapped particle beta value, which is proportional to the injection energy through the ratio $\bar{\omega}_{dh}/\omega_A$.

ACKNOWLEDGMENTS

This work was supported by the U.S. DoE Contract No. DE-AC02-76-CHO-3073.

REFERENCES

- [1] MCGUIRE, K., et al., Phys. Rev. Lett. 50, 891 (1983)
- [2] CHEN, L., WHITE, R.B. and ROSENBLUTH, M.N., Phys. Rev. Lett. 52, 1122 (1984).

- [3] WHITE, R.B., et al., Phys. Fluids 26, 2958 (1983).
- [4] WEILAND, J. and CHEN, L., Annual Controlled Fusion Theory Conference (1984) paper 2P23
- [5] VAN DAM, J.W. and ROSENBLUTH, M.N., ibid.
- [6] MCGUIRE, K. et al., Bull. Am. Phys. Soc. 28, 1173 (1983).
- [7] WHITE, R.B., CHEN, L., ROMANELLI, F. and HAY, R., Princeton Plasma Physics Laboratory Report PPPL-2115 (1984).
- [8] ROSENBLUTH, M.N., TSAI, S.T., VAN DAM, J.W., and ENGQUIST, M.G., Phys. Rev. Lett. 51, 1967 (1983).
- [9] BUSSAC, M. N., PELLAT, R., EDERY, D., and SCALE, J.L., Phys. Rev. Lett. 35, 638 (1975).
- [10] CHANCE, M.S., JARDIN, S.C., and STIX, T.H., Phys. Rev. Lett. 51, 1963 (1983).

FIGURE CAPTIONS

Fig. 1. Results of a "fishbone" simulation, for typical PDX parameters and $\langle \beta_{h,t} \rangle_{\text{crit}} \approx 2 \times 10^{-3}$. Shown here are $\langle \beta_{h,t} \rangle$ and the Mirnov signal $A = \delta B/B$ with maximum amplitude 5×10^{-4} .

FIG. 2. Stability boundaries for the trapped particle-induced ballooning mode (solid curves) and the zero-frequency MHD ballooning mode (dashed curve). The parameters are $q = 1.5$, $\hat{s} = 0.5$ and θ_0 (magnetic turning point) $= \pi/4$.

8470136

Fig. 1

Fig. 2

EXTERNAL DISTRIBUTION IN ADDITION TO TIC UC-20

Plasma Res Lab, Austre Nat'l Univ, AUSTRALIA
 Dr. Frank J. Paoloni, Univ of Wollongong, AUSTRALIA
 Prof. I.R. Jones, Flinders Univ., AUSTRALIA
 Prof. M.H. Brennan, Univ Sydney, AUSTRALIA
 Prof. F. Cap, Inst Theo Phys, AUSTRIA
 Prof. Frank Verheest, Inst theoretische, BELGIUM
 Dr. D. Palumbo, Dg XII fusion Prog, BELGIUM
 Ecole Royale Militaire, Lab de Phys Plasmas, BELGIUM
 Dr. P.H. Sakaneke, Univ Estadual, BRAZIL
 Dr. C.R. James, Univ of Alberta, CANADA
 Prof. J. Teichmann, Univ of Montreal, CANADA
 Dr. H.M. Skarsgard, Univ of Saskatchewan, CANADA
 Prof. S.R. Sreenivasan, University of Calgary, CANADA
 Prof. Tudor W. Johnston, INRS-Energie, CANADA
 Dr. Hannes Bernard, Univ British Columbia, CANADA
 Dr. M.P. Bachynski, MFB Technologies, Inc., CANADA
 Zhenggu Li, SM Inst Physics, CHINA
 Library, Tsing Hua University, CHINA
 Librarian, Institute of Physics, CHINA
 Inst Plasma Phys, Academia Sinica, CHINA
 Dr. Peter Lukac, Komenského Univ, CZECHOSLOVAKIA
 The Librarian, Culham Laboratory, ENGLAND
 Prof. Schatzman, Observatoire de Nice, FRANCE
 J. Redet, CEN-SP6, FRANCE
 AM Dupas Library, AM Dupas Library, FRANCE
 Dr. Tom Muai, Academy Bibliographic, HONG KONG
 Preprint Library, Cent Res Inst Phys, HUNGARY
 Dr. S.K. Trehan, Panjab University, INDIA
 Dr. Indira, Mohan Lal Des, Banaras Hindu Univ, INDIA
 Dr. L.K. Chavda, South Gujarat Univ, INDIA
 Dr. R.K. Chhajlani, Var Ruchi Marg, INDIA
 P. Kew, Physical Research Lab, INDIA
 Dr. Phillip Rosenau, Israel Inst Tech, ISRAEL
 Prof. S. Superman, Tel Aviv University, ISRAEL
 Prof. G. Rostagni, Univ Di Padova, ITALY
 Librarian, Int'l Ctr Theo Phys, ITALY
 Miss Giella De Palo, Assoc EURATOM-CNEN, ITALY
 Biblioteca, del CNR EURATOM, ITALY
 Dr. H. Yamato, Toshiba Res & Dev, JAPAN
 Prof. M. Yoshikawa, JAERI, Tokai Res Est, JAPAN
 Prof. T. Uchida, University of Tokyo, JAPAN
 Research Info Center, Nagoya University, JAPAN
 Prof. Kyoji Nishikawa, Univ of Hiroshima, JAPAN
 Prof. Sigeru Mori, JAERI, JAPAN
 Library, Kyoto University, JAPAN
 Prof. Ichiro Kawakami, Nihon Univ, JAPAN
 Prof. Setoshi Itoh, Kyushu University, JAPAN
 Tech Info Division, Korea Atomic Energy, KOREA
 Dr. R. England, Ciudad Universitaria, MEXICO
 Bibliothek, For-Inst Voor Plasma, NETHERLANDS
 Prof. E.S. Lilley, University of Waikato, NEW ZEALAND
 Dr. Suresh C. Sharma, Univ of Calabar, NIGERIA
 Prof. J.A.C. Cebrai, Inst Superior Tech, PORTUGAL
 Dr. Octavian Petrus, ALI CUZA University, ROMANIA
 Prof. M.A. Mellberg, University of Natal, SC AFRICA
 Dr. Johan de Villiers, Atomic Energy Bd, SO AFRICA
 Fusion Div, Library, JEN, SPAIN
 Prof. Hans Wilhelmson, Chalmers Univ Tech, SWEDEN
 Dr. Lennart Stenflo, University of UMEA, SWEDEN
 Library, Royal Inst Tech, SWEDEN
 Dr. Erik T. Karlson, Uppsala Universitet, SWEDEN
 Centre de Recherches, Ecole Polytech Fed, SWITZERLAND
 Dr. W.L. Weise, Nat'l Bur Stand, USA
 Dr. W.M. Stacey, Georg Inst Tech, USA
 Dr. S.T. Wu, Univ Alabama, USA
 Prof. Norman L. Olsson, Univ S Florida, USA
 Dr. Benjamin Mo, Iowa State Univ, USA
 Prof. Magne Kristiansen, Texas Tech Univ, USA
 Dr. Raymond Askew, Auburn Univ, USA
 Dr. V.T. Tolok, Kharkov Phys Tech Ins, USSR
 Dr. D.D. Ryutov, Siberian Acad Sci, USSR
 Dr. G.A. Eliseev, Kurchatov Institute, USSR
 Dr. V.A. Glukhikh, Inst Electro-Physical, USSR
 Institute Gen. Physics, USSR
 Prof. T.J. Boyd, Univ College N Wales, WALES
 Dr. K. Schindler, Ruhr Universität, W. GERMANY
 Nuclear Res Estab, Jülich Ltd, W. GERMANY
 Librarian, Max-Planck Institut, W. GERMANY
 Dr. H.J. Kaeppeler, University Stuttgart, W. GERMANY
 Bibliothek, Inst Plasmaforschung, W. GERMANY