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L-H Transition in Tokamaks and Stellarators

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Abstract It is shown that the bifurcation dynamics in the H-mode theory based on the ion orbit loss and nonlinear viscosity is not sensitive to the radial pressure and temperature gradients after bifurcation. This allows the poloidal flow velocity to be opposite to that of the poloidal $E \times B$ velocity in the H-mode phase if the pressure gradient is steep. Here, $E(B)$ is the electric (magnetic) field. The theory is also extended to stellarator configurations to demonstrate that the tokamak H-mode theory can be tested in a controlled manner in those devices. A set of poloidal and toroidal momentum evolution equations is derived in which the radial current density outside the last closed flux surface is expressed in terms of the radial integral of the poloidal current density into the divertor plates. This facilitates the calculation of the radial electric field profile across the separatrix.

I. Introduction

Experimental measurements and theoretical models have shown that the L-H transition observed in tokamaks is caused by the sudden change in the poloidal $E \times B$ velocity U_E (TAYLOR, et al. 1989; GROEBNER, et al. 1990; IDA, et al., 1990; WEYNANTS, et al., 1990; ITOH and ITOH, 1988; SHAING and CRUME, 1989; HASSAM, et al. 1991). Here, we will discuss some aspects of the incompressible theory developed in (SHAING AND CRUME, 1989).

The sensitivity of the bifurcation dynamics to the radial pressure and temperature gradients is studied to investigate the poloidal rotation deep in the H-mode phase. It has been found that the poloidal $E \times B$ velocity (actually $U_{p,m}$ as defined in SHAING & CRUME 1989) is not sensitive to these gradients. This allows the poloidal rotation velocity to be opposite to the poloidal $E \times B$ velocity if the pressure gradient is steep (BURRELL, 1993). The theory is also extended to the stellarator devices. It has been found that the poloidal plasma viscosity can have multiple local maxima in the plateau-Pfirsch-Schlüter regime. The position of the local maximum for each (m, n) mode in the magnetic field $|B|$ spectrum is approximately the poloidal $E \times B$ Mach number M_p at $|m - nq|/m$. Here, $m(n)$ is the poloidal (toroidal) mode number and q is the safety factor. However, the particular local maximum of the (m, n) mode may or may not be observed, depending on the amplitude ϵ_{mn} of the mode. This offers an opportunity to test the H-mode theory in a controlled manner by changing the $|B|$ spectrum of a given stellarator or by performing experiments on several stellarators with different $|B|$ spectra. A set of poloidal and toroidal evolution equations is derived in which the radial current density is expressed in terms of the radial integral of the poloidal current density into the divertor plates outside the last closed flux surface ψ_s . The constitutive relations of these equations can be calculated with the results

of the shock theory which includes the effects of the ion diamagnetic flow. It is found that the strength of the poloidal variation of plasma density (TAYLOR, 1991) depends on the direction of the poloidal flow.

II. Sensitivity of Bifurcation Dynamics to Pressure Gradient

The plasma pressure gradient become steep deep into the H-mode phase. The poloidal rotation velocity U_p can be opposite to the poloidal $\mathbf{E} \times \mathbf{B}$ velocity as observed in DIII-D (BURRELL, 1993). To understand this phenomenon in terms of the incompressible bifurcation dynamics in (SHAING & CRUME, 1989), we study the sensitivity of the solution to the pressure gradient based on the equation

$$\langle \mathbf{B} \cdot \nabla \cdot \vec{\pi}_i \rangle + \langle \mathbf{B} \cdot \nabla \cdot \vec{\pi}_e \rangle = -\langle \mathbf{B} \cdot \nabla \cdot \vec{\pi}_i \rangle_{\text{orbit}} \quad (1)$$

where $\vec{\pi}_i$ ($\vec{\pi}_e$) is the ion (electron) viscous tensor. The part of $\vec{\pi}_i$ that is responsible for the orbit loss is separated to $\langle \mathbf{B} \cdot \nabla \cdot \vec{\pi}_i \rangle_{\text{orbit}} = (e/c) \Gamma_{\text{orbit}} \times \mathbf{B} \cdot \mathbf{B}_p$ where Γ_{orbit} is the orbit loss flux, e is the ion charge, c is the speed of light, and $\mathbf{B}_p$ is the poloidal magnetic field. We also include the electron viscosity in Eq. (1) which can be important deep in the H-mode phase because ion viscosity is reduced after bifurcation. Ion and electron viscosity can be written as

$$\langle \mathbf{B} \cdot \nabla \cdot \vec{\pi}_j \rangle = \frac{\sqrt{\pi}}{4} \frac{\epsilon^2}{r} N M_j v_{tj} B (I_{pj} U_{pj} + I_{Tj} U_{pj,o}) \quad (2)$$

in a large aspect ratio tokamak where the subscript j indicates electrons (e) or ions (i), M_j is the mass, v_{tj} is the thermal speed, U_{pj} is the poloidal rotation speed, $U_{pj,o} = \text{sgn}(j) \rho_j v_{tj} (dT_j/dr)/(2T_j)$, T_j is the temperature, and ρ_j is the gyroradius, $\text{sgn}_j = (-1)$ for ions and $(+1)$ for electrons, ϵ is the inverse aspect ratio and r is the local minor radius. To obtain Eq. (2), we assume the parallel electron and the ion heat flows can be neglected. The integrals I_{pj} and I_{Tj} are defined as

$$\begin{aligned} \left\{ \begin{array}{l} I_{pi} \\ I_{Tj} \end{array} \right\} &= \frac{1}{\pi} \int_0^A dx D \left\{ \left(\frac{5}{2} - x \right) \right\} x^2 e^{-x} \int_{-1}^1 dy (1 - 3y^2)^2 \\ &\times \frac{\nu_{*j} \epsilon^{3/2} (\nu_{Tj}/\nu_j \sqrt{x})}{(y + U_{p,m}/\sqrt{x})^2 + [\nu_{*j} \epsilon^{3/2} (\nu_{Tj}/\nu_j \sqrt{x})]^2} \end{aligned} \quad (3)$$

where ν_{*j} is collisionality, $A = \nu_{*i}^{1/2}$ for ions, $A = \infty$ for electrons, $D = (2.5\nu_{*e}/x^2)/[1 + 2.5\nu_{*e}/x^2]$ for electrons, $D = 1$ for ions, $U_{p,m} = 0$ for electrons because v_{te} is large, $U_{p,m} = U_p B / (v_{ti} B_p) - \rho_{pi} (dp_i/dr)/(2p_i)$ for ions, p_i is ion pressure, and ν_{Tj} is the collision frequency for anisotropy relaxation defined in (HIRSHMAN & SIGMAR, 1981).

The sensitivity of Eq. (1) to the pressure gradient can be studied by examining the normalized viscosity $\pi_N = (\langle \mathbf{B} \cdot \nabla \cdot \vec{\pi}_i \rangle + \langle \mathbf{B} \cdot \nabla \cdot \vec{\pi}_e \rangle) (B/v_{ti} B_p) / (\sqrt{\pi}/4) (\epsilon^2/r) N M_i v_{ti} B$ and is

$$\begin{aligned} \pi_N &= \left(I_{pi} + \sqrt{\frac{M_e}{M_i}} \sqrt{\frac{T_e}{T_i}} I_{pe} \right) U_{p,m} - I_p \left(-\frac{1}{2} \frac{\rho_{pi}}{p_i} \frac{dp_i}{dr} \right) + \sqrt{\frac{M_e}{M_i}} \sqrt{\frac{T_e}{T_i}} I_{pe} \left(-\frac{1}{2} \frac{\rho_{pe}}{p_e} \frac{dp_e}{dr} \right) \\ &+ I_{Ti} \frac{U_{pi,o} B}{v_{ti} B_p} + \sqrt{\frac{M_e}{M_i}} \sqrt{\frac{T_e}{T_i}} I_{Te} \frac{U_{pe,o} B}{v_{ti} B_p} - \frac{T_e}{T_i} I_{pe} \frac{J_{||}}{N e v_{te}} \end{aligned} \quad (4)$$

where $J_{||}$ is the parallel current density which includes the Ohmic current and the bootstrap current. The functional dependance on dp_i/dr and dT_i/dr is calculated for two cases shown to simulate the results from the transport equations in Figure (1a) for $\epsilon = 0.25$, $\nu_{*i} = 2$ and neglecting the electron viscosity. For one case, $-\rho_{pi} (dp_i/dr)/(2p_i) = -\rho_{pi} (dT_i/dr)/(2T_i) = 0.1$ and in the other case, $-\rho_{pi} (dp_i/dr)/(2p_i) = 4\{0.1 + (U_{pm}/2)^4/[1 + (U_{pm}/2)^4]\}$ and

Fig. 1a

Fig. 1b

Figure 1 Normalized plasma viscosity versus $U_{p,m}$ for constant pressure and temperature gradients (solid line) and $U_{p,m}$ dependent gradients (dashed line) for $\nu_{*i} = 2$, $\epsilon = 0.25$. In a) only ion viscosity is calculated. In b), electron viscosity is also included with $T_e/T_i = 0.5$.

$-\rho_{pi}(dT_i/dr)/(2T_i) = (1.25/4)[- \rho_{pi}(dp_i/dr)/(2p_i)]$. The complicated U_{pm} dependence for the pressure and temperature gradients is to simulate the steepening of these gradients after bifurcation. As can be seen, the difference in π_N for these two cases decreases as $U_{p,m}$ becomes much larger than unity, even though there is a large difference in dp_i/dr and dT_i/dr . This is because I_p and I_T decrease when $U_{p,m} > 1$. The contribution from dp_i/dr and dT_i/dr thus decreases. This implies that the bifurcated value of U_{pm} is not sensitive to dp_i/dr and dT_i/dr . The poloidal rotation velocity can be opposite to the poloidal $\mathbf{E} \times \mathbf{B}$ velocity if the pressure gradient is very steep. It is, therefore, proper to view Eq. (1) as an equation for U_{pm} . The pressure and temperature gradients are determined from transport equations. The poloidal rotation is calculated from U_{pm} and dp_i/dr from the relation $U_p B / (v_t B_p) = U_{pm} + (\rho_{pi} dp_i/dr) / (2p_i)$. Note that the electron viscosity will also weaken the dp_i/dr and dT_i/dr dependences in π_N if I_{pe} becomes comparable to I_{pi} after bifurcation as shown in Figure (1b) with $T_e/T_i = 0.5$, $\nu_{*e} = 2.0$, and all the other gradients are the same as the ions in Figure (1a).

III. Extension to Stellarators

The incompressible bifurcation theory can be extended to stellarators with the $|\mathbf{B}|$ spectrum given by $B = B_0[1 + \sum_{m,n \neq 0} \epsilon_{mn} \cos(m\theta - n\zeta)]$ with θ the poloidal angle and ζ the toroidal angle. The nonlinear ion viscosity is calculated in Hamada coordinates (SHAING, 1993)

$$\begin{aligned} \langle \mathbf{B}_p \cdot \nabla \cdot \vec{\pi} \rangle &= \frac{\sqrt{\pi}}{4} N M v_{ti} B \sum_{m,n} [I_{mn} \epsilon_{mn}^2 m (m \mathbf{U} \cdot \nabla \theta - n \mathbf{U} \cdot \nabla \zeta) \\ &+ L_{mn} \epsilon_{mn}^2 m \frac{2}{5p} (m \mathbf{q} \cdot \nabla \theta - n \mathbf{q} \cdot \nabla \zeta)]. \end{aligned} \quad (5)$$

where

$$\left\{ \begin{array}{l} I_{mn} \\ L_{mn} \end{array} \right\} = \frac{1}{\pi} \int_0^\infty dx x^2 e^{-x} \left\{ \begin{array}{l} 1 \\ x - \frac{5}{2} \end{array} \right\} \int_{-1}^1 dy (1 - 3y^2)^2 \left(\frac{v \chi'}{B} \right) R_{mn},$$

$R_{mn} = \nu_k / [(m\omega_\theta - n\omega_\zeta)^2 + \nu_k^2]$, $\omega_\theta = (u_{||} + U_{||})\chi'/B + \mathbf{U}_E \cdot \nabla \theta$, $\omega_\zeta = (u_{||} + U_{||})\psi'/B + \mathbf{U}_E \cdot \nabla \zeta$, $\chi' = \mathbf{B} \cdot \nabla \theta$, $\psi' = \mathbf{B} \cdot \nabla \zeta$, ν_k is chosen to be $\nu_k = 3\nu_D + \nu_E$, ν_D and ν_E are defined in (HIRSHMAN and SIGMAR, 1981), $U_{||}$ is the parallel mass flow, $\mathbf{U}_E$ is the $\mathbf{E} \times \mathbf{B}$ velocity, $v_{||}/v = y$ with $v_{||}$ the parallel particle speed, and $\mathbf{q}$ is the heat flow. Assuming the parallel ion mass flow and the heat flow are small in the plateau-Pfirsch-Schlüter regime, the poloidal viscosity can have multiple maxima as shown in Figure (2a) for a simple model

$B = B_0[1 - \epsilon_t \cos \theta - \epsilon_h \cos(m\theta - n\zeta)]$ with $\epsilon_t = 0.1$, $\epsilon_h = 0.1$, $m = 2$, $n = 5$, and $q = 2$ at $\nu_{*i} = 8$. The normalized poloidal viscosity here is defined as $\pi_{p,N} = \langle \mathbf{B}_p \cdot \nabla \cdot \vec{\pi}_i \rangle / (NM v_{ti}^2 \chi' \sqrt{\pi}/4)$. One local maximum is at $M_p \sim 1$, similar to that of a tokamak and the other at $M_p \sim |m - nq|/m$. If $\epsilon_h \gg \epsilon_t$, the tokamak local maximum can be made to disappear, as shown in Figure (2b), with $\epsilon_t = 0.025$, $\epsilon_h = 0.1$ and $\nu_{*i} = 64$. This property can be employed to test the tokamak H-mode theory. Bifurcation can occur either at $M_p \sim 1$ or at $M_p \sim |m - nq|/m$ depending on the poloidal momentum source. However, if bifurcation occurs at $M_p \sim 1$, the neoclassical heat loss, although smaller than that before bifurcation, is still larger than that of an equivalent tokamak which undergoes a similar bifurcation. This is because the contribution of the helical component of the $|\mathbf{B}|$ field spectrum to plasma viscosity and heat viscosity $\langle \mathbf{B}_p \cdot \nabla \cdot \vec{\Theta} \rangle$ and $\langle \mathbf{B}_t \cdot \nabla \cdot \vec{\Theta} \rangle$ still exists even though the contribution of the tokamak component of $|\mathbf{B}|$ spectrum to plasma viscosity and heat viscosity is reduced after bifurcation. Thus, the stellarator H-mode may not be as good as that of an equivalent tokamak unless it bifurcates over all the local maxima. One way to accomplish this is to make $|m - nq|/m$ as close to unity as possible, and the other way, of course, is to drive the poloidal $\mathbf{E} \times \mathbf{B}$ rotation to a value of $M_p > |m - nq|/m$. The L-H transition rate is also slower in stellarators because they usually have larger aspect ratio than tokamaks.

Figure 2 Normalized poloidal viscosity versus poloidal $\mathbf{E} \times \mathbf{B}$ Mach number M_p for a $m = 2$, $n = 5$, and $q = 2$. In a), $\epsilon_t = 0.1$, $\epsilon_h = 0.1$, and $\nu_{*i} = 8$. In b), $\epsilon_t = 0.025$, $\epsilon_h = 0.1$, and $\nu_{*i} = 64$.

IV. Toroidal and Poloidal Momentum Equation

The toroidal and poloidal momentum evolution equation in tokamaks can be obtained by taking the $R^2 \nabla \zeta$ and $\nabla \zeta \times \nabla \psi$ components of the momentum equation

$$\frac{\partial}{\partial t} \langle R^2 \nabla \zeta \cdot N \mathbf{U} \rangle = \frac{1}{cM} \langle \mathbf{J} \cdot \nabla \psi \rangle - \langle R^2 \nabla \zeta \cdot N \mathbf{U} \cdot \nabla \mathbf{U} \rangle - \frac{1}{M} \langle R^2 \nabla \zeta \cdot \nabla \cdot \vec{\pi} \rangle, \quad (6)$$

$$\begin{aligned} \left\langle \frac{R^2 N}{I} \right\rangle \left(\left\langle \frac{B^2}{N} \right\rangle - \frac{I^2}{\langle R^2 N \rangle} \right) \frac{\partial K}{\partial t} &= -\frac{1}{cM} \langle \mathbf{J} \cdot \nabla \psi \rangle + \langle R^2 \nabla \zeta \cdot N \mathbf{U} \cdot \nabla \mathbf{U} \rangle \\ &- \left\langle \frac{R^2 N}{I} \right\rangle \langle \mathbf{B} \cdot \mathbf{U} \cdot \nabla \mathbf{U} \rangle + \frac{1}{M} \langle R^2 \nabla \zeta \cdot \nabla \cdot \vec{\pi} \rangle \\ &- \frac{1}{M \langle R^2 N / I \rangle} \left\langle \frac{\mathbf{B} \cdot \nabla \cdot \vec{\pi}}{N} \right\rangle - \left\langle \frac{R^2 N}{I} \right\rangle \left\langle \frac{\mathbf{B} \cdot \nabla (p_e + p_i)}{NM} \right\rangle, \quad (7) \end{aligned}$$

where $K = N \mathbf{U} \cdot \nabla \theta / \mathbf{B} \cdot \nabla \theta$; $\vec{\pi} = \vec{\pi}_{||} + \vec{\pi}_{\text{orbit}} + \vec{\pi}_{\perp}$ with $\vec{\pi}_{||}$ the CGL viscosity, $\vec{\pi}_{\text{orbit}}$ the part of $\vec{\pi}_{||}$ that is driven by the orbit loss, and $\vec{\pi}_{\perp}$ the gyroviscosity; M is the ion mass with

the subscript i neglected, U is the mass flow velocity, and J is the current density. Note that here, $\vec{\pi}$ includes both the ion and the electron viscosity in general. The flux coordinates employed are $B = I\nabla\zeta + \nabla\zeta \times \nabla\psi$ with $I = R^2 B \cdot \nabla\zeta$, R is the major radius and ψ is the poloidal flux. When the ion diamagnetic flow is of the order of U_E , $\vec{\pi}_\perp$ becomes important and is included. Note that $K = K(\psi)$ if the radial flux is small compared with the poloidal flux. The electron inertia term is neglected. The time evolution of plasma density, pressure, temperature, and their radial gradients are also neglected in Eqs. (6) and (7) for simplicity.

The angular brackets in Eqs. (6) and (7) denote flux surface average on the closed flux surface. It is well-known that $\langle J \cdot \nabla\psi \rangle = 0$ on closed flux surfaces if the V_A^2/c^2 term is ignored where V_A is the Alfvén speed. On open field lines, the angular brackets denote "segment average", defined as $\langle A \rangle = \int_{\theta_1}^{\theta_2} (d\theta/B \cdot \nabla\theta) A / \int_{\theta_1}^{\theta_2} (d\theta/B \cdot \nabla\theta)$ where θ_1 and θ_2 are the end points of the field line at the divertor plates or the toroidally symmetric limiters. The radial current density $\langle J \cdot \nabla\psi \rangle$, in general does not vanish on open field lines. From $\nabla \cdot J = 0$ and the condition that $\langle J \cdot \nabla\psi \rangle = 0$ at the last closed flux surface ψ_s , $\langle J \cdot \nabla\psi \rangle$ can be related to the poloidal current density flow to the divertor plates

$$\langle J \cdot \nabla\psi \rangle = \int_{\psi_s}^{\psi} d\psi' \left\langle B \cdot \nabla \left(\frac{J \cdot \nabla\theta}{B \cdot \nabla\theta} \right) \right\rangle \text{ for } \psi > \psi_s \quad (8)$$

where θ is the poloidal angle. By solving Eqs. (6) and (7), one can determine the toroidal rotation speed and U_E across the separatrix. However, to solve Eqs. (1) and (2), one must express the convective inertia terms and $\vec{\pi}$ in terms of flows. This can be accomplished by solving the kinetic equation to obtain $\vec{\pi}_\parallel$ and the fluid equations to express the convective inertia terms and $\vec{\pi}_\perp$ in terms of flows. The following shock theory is a typical example of such calculations.

V. Effects of Compressibility - Shock

When $M_p \sim 1$, the effects of compressibility can become important and a shock occurs at $M_p \sim 1$ (WINSOR, et al., 1970; HAZELTINE, et al., 1971; GREENE, et al., 1971; GREENE AND WINSOR, 1973; SHAIN, et al., 1992; TANIUTI, et al., 1992). Even if a shock exists, the parallel viscosity still has a local maximum and bifurcation persists. Because the poloidal ion diamagnetic flow, V_{dia} , is not small compared with U_E , one must include V_{dia} in the shock theory. It is found that when U_E and V_{dia} tend to cancel each other to make K small, the effects of compressibility become weaker than the situation when U_E and V_{dia} are in the same direction. However, when the effects of the radial gradient of the parallel flow are also included, the definition of the proper poloidal Mach number becomes complicated. One could conclude (crudely) that it is the poloidal mass flow that is relevant to the shock theory. The characteristic feature of the shock is to generate an up-down asymmetric component of the density and thus, the electrostatic potential within a flux surface. The position of the shock moves in the direction of the "proper" poloidal rotation which is defined in (SHAIN and HSU, 1993). Even when the poloidal rotation speed is not large, the up-down asymmetric component of the density can still contribute to the convective inertia and gyroviscosity terms. As an example, in the linear regime, i.e., when the poloidal Mach number is either much smaller or much larger than unity, the parallel viscosity for ions is,

$$\begin{aligned} \left\langle \frac{B \cdot \nabla \vec{\pi}_\parallel}{N} \right\rangle &= \sqrt{\pi} M v_{ti} B \left\langle \frac{B \cdot \nabla\theta}{N} \right\rangle \left\{ I_1 K \left[\left(\epsilon + \frac{2}{3} \chi_{N,c} \right)^2 + \left(\frac{2}{3} \chi_{N,s} \right)^2 \right] \right. \\ &+ \frac{2}{5} I_2 \frac{q \cdot \nabla\theta}{T_i B \cdot \nabla\theta} \epsilon \left(\epsilon + \frac{2}{3} \chi_{N,c} \right) + \frac{2}{5} (I_2 - I_4) K \left[\left(-\frac{2}{3} \chi_{N,c} + \chi_{T,c} \right) \left(\epsilon + \frac{2}{3} \chi_{N,c} \right) \right. \\ &- \frac{2}{3} \chi_{N,s} \left(\frac{2}{3} \chi_{N,s} - \chi_{T,s} \right) \left. \right] - \frac{2}{3} (I_3 - I_4) \frac{I_{cp_i}}{e B^2} (\ln \bar{N})' [\chi_{T,c} \left(\epsilon + \frac{2}{3} \chi_{N,c} \right) \\ &+ \frac{2}{3} \chi_{T,s} \chi_{N,s}] + \frac{2}{3} (I_3 - I_4) \frac{I_{cp_i}}{d B^2} (\ln \bar{T}_i)' [\chi_{N,c} \left(\epsilon + \frac{2}{3} \chi_{N,c} \right) + \frac{2}{3} \chi_{N,s}^2] \left. \right\}, \quad (9) \end{aligned}$$

where $\chi_N = \ln(N/\bar{N})$, $\chi_T = \ln(T/\bar{T})$, $\bar{N}(\bar{T})$ is normalized plasma density (temperature), $\chi_N = \chi_{N,s} \sin \theta + \chi_{N,c} \cos \theta$, $\chi_T = \chi_{T,s} \sin \theta + \chi_{T,c} \cos \theta$, prime denotes $d/d\psi$, and I_1 , I_2 , I_3 , and I_4 are defined in (SHAING and HSU, 1993). The compressible viscosity, namely the contribution to plasma viscosity from the region away from the shock angle when shock occurs, is

$$\left\langle \frac{\mathbf{B} \cdot \nabla \cdot \vec{\pi}_{\parallel}}{N} \right\rangle_{\text{com}} = \frac{\sqrt{\pi}}{3} M v_{ti} B \left\langle \frac{\mathbf{B} \cdot \nabla \theta}{N} \right\rangle \epsilon \frac{G}{A'} \left\{ \frac{2}{3} I_1 K - \frac{2}{3} (I_3 - I_4) \frac{I_{cp_i}}{e B^2} \times \left[\frac{2}{3} (\ln \bar{N})' - (\ln \bar{T})' \right] \right\}, \quad (10)$$

for adiabatic ions, and is

$$\left\langle \frac{\mathbf{B} \cdot \nabla \cdot \vec{\pi}_{\parallel}}{N} \right\rangle_{\text{com}} = \frac{\sqrt{\pi}}{3} M v_{ti} B \left\langle \frac{\mathbf{B} \cdot \nabla \theta}{N} \right\rangle \epsilon \frac{G}{A'} \left[\frac{2}{3} I_1 K - \frac{4}{15} (I_2 - I_4) K - \frac{2}{3} (I_3 - I_4) \frac{I_{cp_i}}{e B^2} (\ln \bar{T})' \right], \quad (11)$$

for isothermal ions, where A' is defined in (SHAING and HSU, 1993).

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