# Equilibrium and Fluctuations in a Plasma Confined in a Pure Toroidal Field

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Equilibrium and confinement in currentless toroidal devices is explained in terms of a flow-fluctuation cycle. The initial limiter equilibrium is shown to get fortified via fluctuation driven poloidal rotation and ponderomotive force in a self-consistent manner. The relevance of this to the recently proposed *L-H* transition theories is briefly discussed. [S0031-9007(96)00596-0]

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It is well known that plasma cannot be confined in a pure toroidal magnetic field. This is because the curvature of field lines gives rise to an effective gravity in which plasma experiences a free fall along the major radius ($R$). In a typical moderate size machine the free fall time is about a few $\mu$s. Surprisingly, however, experimental observations have shown that the typical confinement time in such currentless toroidal devices (CTD) is about 2 orders of magnitude larger than the free fall time. This interesting observation has attracted considerable attention in the field.

Yoshikawa *et al.* [1] suggested that the space charge on a plasma surface can be shorted out by a conducting limiter and provide what is known as the "limiter equilibrium." However, in this equilibrium the space charge within the plasma is uncompensated and attempts have been made to find additional mechanisms of rotational transform which may operate within the plasma. The fluctuation driven cross-field current [2], or poloidal rotation which is generally observed in CTD [3], has been invoked for this purpose.

The mechanism involving poloidal rotation and improved confinement finds strong support in a recent experiment by Jain [4] where the rotation is enhanced by a biased probe driven radial electric field. The upshot of all this is that the fluctuation driven and/or externally assisted poloidal rotation or radial electric field is one of the primary reasons for enhanced confinement in CTD. However, this explanation is incomplete. It is well known that flows affect the characteristics of many types of fluctuations [5,6]. Thus not only fluctuations affect (or generate) flows but also flows affect fluctuations. This flow-fluctuation synergy is the central theme of recent theories by Diamond and co-workers [5–8] regarding high confinement transitions in tokamaks, e.g., *L-H* transition, CH transition, and recently discovered enhanced reverse shear (ERS) transition. A self-consistent explanation must thus take into account not only the generation of flows due to fluctuations but also the backreaction of flows on fluctuations.

In this Letter, we provide such an explanation of equilibrium and confinement in CTD in terms of the flow-fluctuation cycle. Briefly, our model is as follows. Initially, the limiter provides the "seed equilibrium." In this equilibrium, fluctuations which are generally due to Rayleigh-Taylor (RT) instability [9–12] grow to a significant level. These fluctuations provide rotational transform in two ways. Firstly, fluctuations directly drive a radial current $J_r$ [2] or poloidal rotation which improves the limiter equilibrium. Secondly, the flow back reacts on fluctuations to modify the rms level profile. The mean ponderomotive force (PF) due to these fluctuations then opposes the free fall of the plasma and further fortifies the limiter equilibrium. The two new elements of our model are as follows. Firstly, we show the generation of rotational transform through the fluctuation driven PF. In some cases, this mechanism is found to be more efficient than the one involving a radial electric field. Secondly, and perhaps more importantly, our model enlarges the role of the flow-fluctuation cycle in the context of confinement physics. In tokamaks, the rotational transform is provided by the toroidal current, while the flow-fluctuation cycle provides an access to an equilibrium with quenched fluctuations, i.e., $H$ mode. In CTD, on the other hand, this cycle provides the basic rotational transform itself required for the equilibrium. In this sense, the role of the flow-fluctuation cycle in confinement physics becomes even more important.

We begin by describing the axisymmetric plasma equilibrium in the poloidal plane using cylindrical coordinates $\hat{R}$, $\hat{\zeta}$, $\hat{Z}$ ($\hat{R}$ is along the major radius, $\hat{\zeta}$ is the toroidal angle). Within the single fluid magnetohydrodynamics, the equation of motion is given by [13]

$$\rho\left[\frac{\partial \vec{V}}{\partial t} + \left(\vec{V}\cdot\vec{\nabla}\right)\vec{V}\right] = \frac{\vec{J}\times\vec{B}}{c} + \frac{2P}{R}\hat{R} - \rho\vec{V}\nu_{in}, \quad (1)$$

where $\rho$, $\vec{V}$, and $\vec{J}$ are plasma mass density, velocity and current density, respectively. This equation is a model equation where the effective gravity due to field line curvature is represented by the term $(2P/R)\hat{R}$ and suffices for the compact derivation of the results given here. Earlier, this equation has been used by Coppi *et al.* [13] and Rosenbluth *et al.* [13] for studying curvature driven modes. The term $\nu_{in}$ represents ion-neutral collision frequency. The $R$ component of this equation describes the free fall along $\hat{R}$ in the presence of various forces

$$\frac{\partial V_R}{\partial t} = -[(\vec{V}\cdot\vec{\nabla})\vec{V}]_R + \frac{2c_s^2}{R} - \nu_{in}V_R, \quad (2)$$

 

where $\vec{B} = |B|\hat{\zeta}$ and $c_s$ is the ion sound speed. In Eq. (2) we have neglected the $\hat{R}$ component of the term $\vec{J} \times \vec{B}$. Using Possion's equation, the continuity equation for charges, and Ohm's law, it can be shown that the ratio of this term to the inertial term on the left hand side of Eq. (2) is $V_A^2/c^2$, which is small ($V_A$ is the Alfvén speed). We next consider PF due to fluctuation given by $[(\tilde{v} \cdot \vec{\nabla})\tilde{v}]_R \sim (1/2)\partial \tilde{v}_R^2/\partial R$ ($\tilde{v}_R$ is the perturbed velocity). From Eq. (2), without the fluctuation effect, one gets the terminal velocity of the fall $V_R \sim 2c_s^2/\nu_{in}R$. The ion-neutral collisions impede the radial acceleration of plasma due to effective gravity. It is interesting to note that a significant improved confinement can be observed in such devices if $\nu_{in} > (2c_s^2/Ra)^{1/2}$. The terminal velocity of the fall in the presence of fluctuations is given by

$$V_R = \frac{1}{\nu_{in}}\left[\frac{2c_s^2}{R} - \frac{1}{2}\frac{\partial}{\partial R}\tilde{v}_R^2\right]. \qquad (3)$$

If $\partial \tilde{v}_R^2/\partial R > 0$, the PF due to fluctuations opposes the free fall. For static equilibrium, the estimated critical fluctuation level is $e\tilde{\varphi}_c/T_e \approx (2/k_\theta a_s)\sqrt{a/R}$ where $\tilde{\varphi}$ is the potential fluctuation, $\tilde{v}_R = -i(c/B)k_\theta \tilde{\varphi}$, $k_\theta$ is the poloidal wave number, and $a_s$ is the Larmor radius with ion sound speed. It is interesting to note that the plasma confinement time $a/V_R$, where $a$ is the minor radius, can be extended by (1) exciting the shorter wavelength fluctuations, (2) going to higher atomic mass, and (3) increasing neutral pressure.

We now present some results from a device called BETA [4] which are consistent with this theoretical study. For typical BETA parameters of the argon plasma $R = 45$ cm, $a = 15$ cm, plasma radius $r = 10$ cm, $T_e \sim 5$ eV, $T_i \sim 0.2$ eV, $B = 200$ G, plasma density $N \sim 5 \times 10^{10}$ cm$^{-3}$, neutral density $N_n \sim 1.5 \times 10^{13}$ cm$^{-3}$, $k_\theta \sim 0.3$ cm$^{-1}$, and equilibrium scale lengths of density, potential, and mean rms fluctuation level are $\sim$15 cm, we get $c_s \sim 3.5 \times 10^5$ cm/s, $a_s \sim 7$ cm, $\nu_{in} \sim 10^4$ s$^{-1}$, and $e\tilde{\varphi}_c/T_e \approx 55\%$. In Fig. 1 we show the experimentally measured rms fluctuation level profile at $B = 800$ and 200 G. The profile has an off-axis maxima with a value $e\tilde{\varphi}/T_e \approx 45\%$ (and $\partial|\tilde{v}_R^2|/\partial R > 0$ on the outboard) which within the experimental uncertainties is sufficient to arrest the free fall on the outboard. The figure also shows significant level on the inboard. The implication of this will be discussed at the end.

We would now like to describe the theory of fluctuations in a toroidal device (like BETA). Since such devices typically always have self-consistent poloidal flows, we shall investigate the theory in the presence of such flows. In the last part of our calculation, we will make our analysis self-consistent by calculating poloidal flow generation due to these fluctuations via Reynolds stress.

As has been identified earlier, the fluctuations in CTD are generally due to RT mode [9–12]. We investigate the effect of poloidal flow on linear RT mode. This problem has been investigated by a number of authors earlier [14].

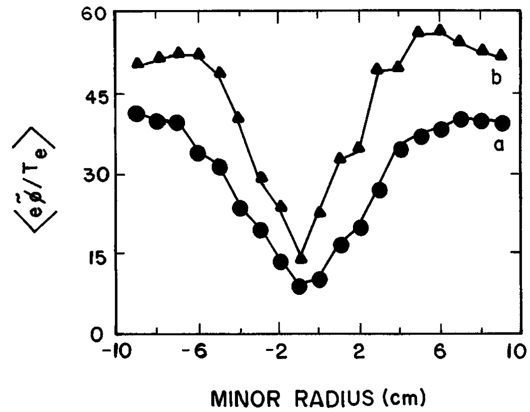


FIG. 1. Radial distribution of potential fluctuation amplitude for (a) $B = 200$ G and (b) $B = 800$ G.

Since the wavelength of these modes is small as compared to $a$, we use slab approximation ($x$, $y$, $z$). The plasma is in a vacuum field which decreases as $1/R$. We assume that the linear eigenfunction is slowly varying in $y$ (or $\theta$) so that the poloidal variation of the mode amplitude is neglected. We thus consider an eigenfunction of the form $\tilde{\varphi} = \varphi(x)\exp[-i(\omega t - k_y y)]$. The electron continuity and $\vec{\nabla} \cdot \vec{J} = 0$ equations give two coupled equations for $\tilde{n}$ and $\tilde{\varphi}$:

$$\frac{d\tilde{n}}{dt} + V_{*e}\frac{\partial \tilde{\varphi}}{\partial y} + \epsilon_n V_{*e}\frac{\partial}{\partial y}(\tilde{n} - \tilde{\varphi}) = 0, \qquad (4)$$

$$a_s^2\left[\frac{d}{dt} - KV_{*e}\frac{\partial}{\partial y}\right]\left(\frac{\partial^2}{\partial x^2} - k_y^2\right)\tilde{\varphi} - a_s^2\left(-\frac{V_E'}{L_n} + V_E''\right)\frac{\partial \tilde{\varphi}}{\partial y} = -\epsilon_n V_{*e}\frac{\partial \tilde{n}}{\partial y}, \qquad (5)$$

where the dependent variables are normalized as $\tilde{n} = n/N$, $\tilde{\varphi} = e\varphi/T_e$, $N$ is the equilibrium density, $n$ and $\varphi$ are the perturbed density and potential fluctuations, $d/dt = \partial/\partial t + \vec{V}_E \cdot \vec{\nabla}$, $V_{*e} = a_s c_s/L_n$ is the diamagnetic drift velocity, $a_s = c_s/\Omega_i$, $c_s = \sqrt{T_e/m_i}$, $L_n^{-1} = -d\ \ln N/dx$, $K = T_i/T_e$, $\Omega_i$ is the ion cyclotron frequency, $\epsilon_n = 2L_n/R$, $V_E'$ and $V_E''$ are the first and second derivatives of $\vec{E} \times \vec{B}$ in the $x$ direction, $V_E = -cE_x/B$. Here we have ignored the electron and ion temperature perturbations.

We now solve the coupled Eqs. (4) and (5) in weak shear limit. The effect of the strong velocity shear on the RT mode will be investigated in a future paper. We choose $E_x = E_x(x_0) + E_x'(x_0)(x - x_0) + E_x''(x_0)(x - x_0)^2/2$ to separate the shear $[E_x'(x_0)]$ and curvature $[E_x''(x_0)]$ effects around the local radial point $x = x_0$. Eliminating $\tilde{n}$ from Eqs. (4) and (5), we get the eigenvalue equation in the weak shear limit as

$$\frac{\partial^2 \tilde{\varphi}}{\partial \tau^2} = \left[k_y^2 a_s^2 - \frac{2\hat{V}_E'' - a_s\hat{V}_E'/L_n}{\overline{\omega} + K} + \frac{\epsilon_n(1 - \epsilon_n)}{(\overline{\omega} + K)(\overline{\omega} - \epsilon_n)} - \frac{\alpha^2}{4\beta} + \beta\tau^2\right]\tilde{\varphi}, \qquad (6)$$

where $\tau = (x - x_0)/a_s + \alpha/2\beta$, $\overline{\omega} = [\omega - kV_E(x_0)]/\omega_{*e}$ is the normalized Doppler shifted eigenvalue, $\hat{V}_E' = L_n V_E'/c_s$, $\hat{V}_E'' = (a_s L_n/2)V_E''/c_s$ evaluated at $x = x_0$, and

$$\alpha = \frac{2a_s\hat{V}_E''}{L_n(\overline{\omega} + K)} + \frac{\epsilon_n(1 - \epsilon_n)(2\overline{\omega} + K - \epsilon_n)}{(\overline{\omega} + K)^2(\overline{\omega} - \epsilon_n)^2}\hat{V}_E',$$

$$\beta = \frac{\epsilon_n(1 - \epsilon_n)(2\overline{\omega} + K - \epsilon_n)}{(\overline{\omega} + K)^2(\overline{\omega} - \epsilon_n)^2}\hat{V}_E''. \tag{7}$$

Equation (6) is a standard Hermite equation. For $l = 0$ ($l$ is the radial mode number) case, the eigenfunction of Eq. (6) $\tilde{\varphi} \sim \exp\{-\beta^{1/2}[(x - x_0)/a_s + (\alpha/2\beta)]^2\}$, while the linear dispersion relation is expressed as

$$k_y^2 a_s^2 - \frac{2\hat{V}_E'' - a_s\hat{V}_E'/L_n}{\overline{\omega} + K} + \frac{\epsilon_n(1 - \epsilon_n)}{(\overline{\omega} + K)(\overline{\omega} - \epsilon_n)} - \frac{\alpha^2}{4\beta} = -\beta^{1/2}. \tag{8}$$

We now take the effect of electric field shear and curvature perturbatively in Eq. (8). For $T_i/T_e \ll 1$, the result is

$$\overline{\omega} = \frac{\epsilon_n}{2} + i\overline{\gamma}_0 - \frac{\hat{V}_E'^2}{4\hat{V}_E''} + \frac{\hat{V}_E''}{k_y^2 a_s^2}\left[1 + \frac{1}{4k_y^2 L_n^2}\left(\frac{\epsilon_n^2}{4\overline{\gamma}_0^2} - 1\right)\right] + i\frac{\hat{V}_E''}{k_y^2 a_s^2}\frac{\epsilon_n}{2\overline{\gamma}_0}\left(1 - \frac{2}{k_y^2 L_n^2}\right) + (1 - i)\frac{1}{2k_y^2 a_s^2}\left[\frac{\epsilon_n(1 - \epsilon_n)}{\overline{\gamma}_0}\hat{V}_E''\right]^{1/2}, \tag{9}$$

with the constraint $\mathrm{Re}(\beta^{1/2}) > 0$ for the growing mode so that the eigenfunction is spatially bounded. The first two terms on the right hand side (r.h.s.) of Eq. (9) are the standard response of the RT mode (where $\overline{\gamma}_0 = [\epsilon_n(1 - \epsilon_n)]^{1/2}/k_y a_s$). The third, fourth, and the real part of the last term on the r.h.s. contribute to the Doppler shifted real frequency. The growth due to the fifth term is neglected in comparison to that of the last term since $(\epsilon_n\hat{V}_E''/\overline{\gamma}_0)^{1/2} < 1$. Thus the imaginary part of the last term indicates the destabilization of the modes due to electric-field curvature for $\hat{V}_E''(x_0) < 0$. For $\hat{V}_E''(x_0) > 0$, the mode stabilizes. The behavior of the eigenfunction $\tilde{\varphi}$ at $x \to \infty$ is bounded and the typical radial mode width $\Delta = a_s/2^{1/2}\beta^{1/4} \sim (0.92/k_y)[\epsilon_n/8\overline{\gamma}_0\hat{V}_E''(x_0)]^{1/4}$. For typical BETA parameters, $\Delta \sim 1$ cm. Here, it is interesting to note that the electric-field shear affects only the real part of frequency ($\overline{\omega}$).

In the last part of our calculations, we make the analysis self-consistent by calculating the mean poloidal flow generation in the core due to Rayleigh-Taylor fluctuations. As shown by Diamond and Kim [15], the mean poloidal flow in the core at the radial location $x = x_0$, due to Reynolds stress, is given by $\partial\vec{V}_y(x_0)/\partial t = -\partial\overline{\pi}_{xy}(x_0)/\partial x_0 - \nu_{in}V_y(x_0)$ where the last term describes the damping of flows due to ion-neutral collisions, $\overline{\pi}_{xy}(x_0)$ is the mean Reynold's stress at $x_0$, and is defined as $\overline{\pi}_{xy}(x_0) = \int_{-\infty}^{\infty}[\tilde{v}_x\tilde{v}_y^* + \tilde{v}_x^*\tilde{v}_y]dx$, $\vec{\tilde{v}} = (c/B)\hat{z} \times \vec{\nabla}\tilde{\varphi}$. (It should be noted that now $x_0$ is the radial coordinate with the origin at $x_0 = 0$.) The condition for $\overline{\pi}_{xy}(x_0) \neq 0$ is that the linear eigenfunction $\tilde{\varphi}$ should be radially inhomogeneous and asymmetric [15]. In the present case, the radial asymmetry comes from $V_E'$. This gives rise to a partial differential equation for $V_y(x_0, t)$ which can be solved in the steady state to obtain $V_y(x_0)$ [$= V_E(x_0)$]. With the eigenfunction $\tilde{\varphi}$, the flow equation $\partial\overline{\pi}_{xy}(x_0)/\partial x_0 + \nu_{in}V_y(x_0) = 0$ gives the following equation involving $|\tilde{\varphi}_0|^2$ and $V_y(x_0)$:

$$a_s\frac{d}{dx_0}[G_1|\tilde{\varphi}_0|^2\hat{V}_E''^{1/4}(x_0)] = \frac{1}{k_y a_s}\frac{\nu_{in}}{\Omega_i}\frac{V_E(x_0)}{c_s}, \tag{10}$$

where $G_1 = (8\pi)^{1/2}F\exp[F^2\hat{V}_E''^{1/2}(x_0)]$ and $F = \epsilon^{3/4}/[2k_y L_n(1 - \epsilon_n)^{1/2}\gamma_0^{-3/4}]$. It should be noted that Eq. (10) is still not an equation for $V_E(x_0)$ because $|\tilde{\varphi}_0|^2$, which is the rms level of fluctuation, is still a function of $x_0$ which is not yet determined. This should be determined by some other physical arguments. In the present case, we use the following physical arguments based on the linear theory described earlier. The radial mode width of the mode $\Delta$ is small, i.e., $\Delta/a \ll 1$. In this case, the mode will sense only the local velocity shear and curvature rather than the average global profile. Now, the condition $V_E''(x_0) < 0$ is satisfied around the maximum of $V_E(x_0)$ and, according to the linear theory, the mode is destabilized here; hence, the fluctuation level $\tilde{\varphi}_0^2(x_0)$ is expected to be maximum at this location. Similarly, the fluctuation level will be a minimum at locations where $V_E''(x_0) > 0$. It is thus reasonable to look for a $|\tilde{\varphi}_0|^2$ profile which follows the $V_E(x_0)$ profile, i.e., $|\tilde{\varphi}_0|^2 = G_2\hat{V}_E(x_0)$, where $\hat{V}_E(x_0) = V_E(x_0)/c_s$ and $G_2$ is a constant. Using this ansatz in Eq. (10), we obtain the following differential equation for poloidal flow generation in the core:

$$\hat{V}_E''^{1/4}(x_0)\hat{V}_E'(x_0) = \frac{1}{k_y a_s}\frac{1}{G_1G_2}\frac{\nu_{in}}{\Omega_i}\frac{L_n}{a_s}\hat{V}_E(x_0). \tag{11}$$

Here, third and higher order derivatives of $\hat{V}_E(x_0)$ have been neglected, and $G_1$ and $F$ are assumed to be weak functions of $x_0$. With the boundary conditions $\hat{V}_E(x_0)$ and $\hat{V}_E'(x_0) = 0$ at $x_0 = 0$, the solution of Eq. (11) is given by $|\tilde{\varphi}|^2 \propto V_E(x_0) = \alpha^4 x_0^6$, where $\alpha = 0.07\nu_{in}L_n/k_y a_s^2 G_1G_2\Omega_i$. This solution implies a monotonically increasing rms level in the core consistent with the experimental observations shown in Fig. 1. The magnitude of $V_E$ can be estimated from Eq. (1) as $\hat{V}_E \sim (k_y c_s a_s^2/\nu_{in}L_\phi\Delta_x)|\tilde{\varphi}|^2$ where $k_x = \Delta_x^{-1}$ and $L_\phi$ is the equilibrium scale length of rms fluctuation level. For BETA parameters, we get $V_E \sim c_s$. For poloidal rotation to effectively neutralize the charge accumulation in the core, the poloidal convective time $a/V_E$ should be smaller than the free fall time $a/V_R$, thus $V_E \geq V_R \sim 2c_s^2/\nu_{in}R$ for BETA parameters $V_E \sim V_R$, which implies that the fluctuation driven poloidal rotation is just about sufficient to neutralize the charge accumulation in the core. In this paper, we have discussed the poloidal flow generation via the Reynold stress $[(\tilde{v} \cdot \vec{\nabla})\tilde{v}]_\theta$. Equivalently, one can also consider poloidal flow generation using the term $J_r B_\zeta$ ($J_r$ is the radial current in the poloidal cross section) as is done by Rypdal *et al.* [2]. However, the two approaches are equivalent and give identical results. In a detailed paper, Diamond and Kim [15,16] have shown the equivalence for the drift wave fluctuations and, in general, for any class of electrostatic fluctuations.

We now return to the point regarding the fluctuations on the inboard (Fig. 1). One may surmise that these fluctuations could be due to drift waves or any other instability which does not depend on magnetic curvature [17,18]. However, on the basis of $\tilde{n}$ and $\tilde{\varphi}$ phase relationship, there is strong evidence in BETA that the nature of fluctuation on the inboard and outboard is nearly the same [9–11]. One could then argue that the poloidal rotation convects fluctuations in $\theta$, which are originally created on the outboard [17]. The condition for this is $2\pi\gamma_0 a/V_E < 1$ where $a/V_E$ is the time for poloidal convection for typical BETA parameters $2\pi\gamma_0 a/V_E \sim a/\sqrt{RL_N} \sim 1$. Hence fluctuations could be convected azimuthally by poloidal rotation. A consequence of the fluctuation on the inboard is that PF due to them is in the direction of the free fall (along $\hat{R}$). On the outboard, on the other hand, the PF effectively opposes the free fall. One could thus argue that equilibrium on the outboard is mainly due to PF while on the inboard it is mainly due to poloidal rotation. However, since the poloidal convection time is just about equal to the free fall time and also since the outward force along $\hat{R}$ is more (PF + gravity), the equilibrium on the inboard is not very good. As a result most of the plasma resides on the outboard. This is consistent with the experimentally measured profile of density in BETA [4,10] where profiles are more sharp in the inboard than in the outboard, which indicates an accumulation of the plasma on the outboard.

To summarize, we have developed a model of equilibrium and confinement in CTD based on the flow-fluctuation cycle. In the initial seed equilibrium via limiter, the Rayleigh-Taylor fluctuations grow. These fluctuations provide the rotational transform in two ways. Firstly, poloidal flow (radial electric field) required to neutralize the charge accumulation is generated via Reynold stress [Eq. (11)]. Secondly, the poloidal flow back reacts on fluctuations to modify the rms profile in such a way that the ponderomotive force opposes the gravity [Eqs. (2) and (9)]. This interplay of equilibrium, fluctuations, and poloidal flow is responsible for improving the confinement in CTD. We may now compare the role of flow-fluctuation cycle in CTD vis-à-vis tokamaks. In tokamaks, the rotational transform is due to the toroidal current. Next, according to the theories of Diamond and co-workers [5–8], the $L$ mode is a state of high fluctuation level and low poloidal rotation. When the external drive, say neutral beam power, exceeds a threshold, fluctuations generate flow which in turn back reacts on fluctuations to give rise to a new state of large rotation and a low level of fluctuations, i.e., the $H$ mode. In CTD, on the other hand, there is no equilibrium. However, we have shown that the flow-fluctuation cycle provides the basic rotational transform itself. Thus the role of the flow-fluctuation cycle in CTD is similar to that of the flow-fluctuation cycle [5–8] in a tokamak *albeit with important differences*.

The two crucial elements of the theories of Diamond *et al.*, i.e., flow generation and modification of fluctuation by flow, can be seen at work in CTD. Besides, the suppression of fluctuations by flows in CTD has been experimentally demonstrated by Jain [4]. It thus appears that the CTD can serve as a test bed for some of the hot ideas which are being currently discussed in high confinement tokamak physics.

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[1] S. Yoshikawa, W. L. Harries, and R. M. Sinclair, Phys. Fluids **6**, 1506 (1963).
[2] K. Rypdal *et al.*, Plasma Phys. Controlled Fusion **36**, 1099 (1994).
[3] R. Van Nieuwenhove, Plasma Phys. Controlled Fusion **34**, 873 (1992).
[4] K. K. Jain, Phys. Rev. Lett. **70**, 806 (1993).
[5] H. Biglari, P. H. Diamond, and P. W. Terry, Phys. Fluids B **2**, 1 (1989).
[6] P. H. Diamond *et al.*, Phys. Rev. Lett. **72**, 2565 (1994).
[7] B. A. Carreras *et al.*, Phys. Plasmas **1**, 4014 (1994).
[8] P. H. Diamond *et al.* (to be published).
[9] D. Bora, Phys. Lett. **139**, 308 (1989).
[10] G. Prasad *et al.*, Phys. Plasmas **1**, 1832 (1994).
[11] G. Prasad, D. Bora, and Y. C. Saxena, Geophys. Res. Lett. **19**, 241 (1992); **19**, 245 (1992).
[12] A. Komori, N. Sato, and Y. Hatta, Phys. Rev. Lett. **40**, 768 (1978).
[13] P. H. Furth *et al.*, in *Proceedings of the Conference on Plasma Physics and Controlled Nuclear Fusion Research, Culham Laboratory, Berkshire, England, 1965* (IAEA, Vienna, 1966), Vol. 1, p. 103; B. Coppi and M. N. Rosenbluth, *ibid.*, Vol. 1, p. 617.
[14] P. N. Guzdar *et al.*, Geophys. Res. Lett. **9**, 547 (1982); A. B. Hassam, Phys. Fluids B **4**, 485 (1992).
[15] P. H. Diamond and Y. Kim, Phys. Fluids B **3**, 1626 (1990).
[16] P. H. Diamond (private communication).
[17] P. K. Kaw (private communication).
[18] S. Jacob *et al.*, in *Proceedings of the 15th International Conference on Plasma Physics and Controlled Nuclear Fusion Research, Seville, Spain, 1994* (IAEA, Vienna, 1995).