Theory of Shear Suppression of Edge Turbulence by Externally Driven Radio-Frequency Waves

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Here, we propose and analyze a technique for active suppression of tokamak edge turbulence. Suppression occurs due to the effects of a sheared radial electric field generated by externally driven radio-frequency waves. Plasma flow is induced by radially varying wave-driven Reynolds and magnetic stresses, and opposed by neoclassical damping. For Alfvénic flow drive, the predicted shear flow profile is determined by ion inertia and electron dissipation effects. Results indicate that a modest amount of absorbed power is required for edge-turbulence suppression. More generally, several novel results in the theory of momentum transport by electromagnetic fluctuations are presented.

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The understanding and control of turbulent transport of heat and particles are the principal obstacles confronting controlled fusion research today [1]. Progress toward this end has been aided by the discovery of various improved confinement regimes. Foremost of these is the *H* mode [2], in which a transport barrier spontaneously forms at the plasma periphery or edge, resulting in the development of steep density and, occasionally, temperature gradients. Prior to the formation of this barrier, i.e., during the *L* mode, this periphery region is characterized by large electromagnetic and density fluctuation levels, indicating a region of strong edge turbulence [3]. Recently, the $L \rightarrow H$ transition has been observed to coincide with the onset of large poloidal flows in the edge region [4]. Such flows are thought to be symptomatic of a strongly sheared inward electric field, which in turn can suppress turbulence and transport via shear-enhanced decorrelation [5]. The transport barrier then forms as a consequence of the turbulence suppression. Indeed, recent studies of edge shear layer structure in non-*H*-mode discharges indicate that shear decorrelation is likely to play an important role in regulating edge transport in *all* confinement regimes [6].

In this Letter, we discuss the suppression of turbulence using sheared electric fields generated by externally driven radio-frequency (rf) waves. This technique for turbulence suppression has several potential advantages over the spontaneous $L \rightarrow H$ transition. First, it allows enhanced confinement with $E_r > 0$, thus preventing the undesirable impurity accumulation characteristic of spontaneous *H* mode (with $E_r < 0$), and facilitating ash removal. Second, the proposed technique can render edge localized mode (ELM) [7] control feasible, via confinement control. Third, rf flow drive permits controlled, perturbative experiments on the $L \rightarrow H$ transition and angular momentum transport. Finally, rf flow drive does not suffer from the limitations intrinsic to electrode-driven flows [8] or from the disadvantages of wall-sputtering associated with flow drive by neutral-beam injection [9].

While rf flow drive is possible in a broad range of frequencies, we focus here on Alfvén-wave flow drive. This frequency range of flow drive has inherent basic physics interest outside the realm of magnetic fusion. Indeed, flow drive processes are analogous to wave-driven dynamos [10] or wave helicity injection. Therefore, wave-driven flows could be present in regions of wave-driven dynamo activity. In this sense, the theory of Alfvén-wave-driven flows can be applied to resonant absorption in the magnetosphere [11] and the solar corona [12] as well as in accretion disks [13]. Flows driven by Alfvén waves can also impact the star formation process [14]. Here, waves can spin up ionized protostellar clouds, thus triggering rotation-driven dynamos.

In fusion, our choice of the Alfvén frequency range is motivated by considerations of relatively good coupling efficiency, the (high) relative size of the frequency compared to the edge decorrelation rate $\Delta\omega$, simplicity, and the desire to facilitate insight into a recently proposed, related scheme for flow drive using oscillating fields driven by external resonant coils [15]. To minimize losses due to radial attenuation, we anticipate the location of the Alfvén resonance r_s [where $\omega = k_\parallel(r_s)v_A$, $v_A^2 = B_0^2/4\pi\rho_0$, ρ_0 is the ion mass density, and $\mathbf{B}_0(r_s)\cdot\nabla = B_0 ik_\parallel(r_s)$] to be in the plasma edge. The external rf field rapidly undergoes mode conversion to the kinetic shear Alfvén wave (KSAW) [16] for $r \le r_s$ in a hot edge ($v_A < v_{Te}$, v_{Te} is the electron thermal velocity) and for $r \ge r_s$ in a cold edge [17] ($v_A > v_{Te}$). In the former case, we envision placing r_s very close to the edge to minimize evanescence prior to the resonance, while in the latter case we would attempt to place r_s slightly inside the edge to set up a reflected wave. (We choose the frequency of the wave to avoid driving a standing wave at the edge.) Bulk plasma flows are then driven by the KSAW, for which the radial wave vector k^R is given by (under the assumption of electron temperature much greater than ion temperature)

$$k^R = k_r^R + ik_i^R \,, \tag{1a}$$

$$k_r^R \rho_s = \left[\Delta\left(\frac{\omega^2}{k_\parallel^2 v_A^2} - 1\right)\right]^{1/2} = \left[\Delta\left(\frac{n}{n_{\mathrm{crt}_{\omega,\mathbf{k}}}} - 1\right)\right]^{1/2} , \tag{1b}$$

$$k_i^R \rho_s = \left(\frac{\delta_{e\mathbf{k},\omega}}{2}\right)|k_r^R \rho_s| \,. \tag{1c}$$

Here, for $|r-r_s| > L_n$ (L_n is the density scale length), $\Delta = 1$ for a hot edge and $\Delta = -m_e\omega^2/\omega_{ci}^2 k_\parallel^2 \rho_s^2$ for a cold edge (where $\rho_s^2 = c_s^2/\omega_{ci}^2$, $c_s^2 = T_e/m_i$, ω_{ci} is the ion cyclotron frequency, T_e is the electron temperature, m_e is the electron mass, and m_i is the ion mass). Note that since $\Delta < 0$ for the cold edge, then the radial wave number is real (radially propagating wave) for $n < n_{\mathrm{crt}}$ while for $n > n_{\mathrm{crt}}$ the wave is evanescent. For $|r-r_s| < L_n$, $\Delta \rightarrow (\rho_s/L_n)^{2/3}\Delta$, on account of the Airy-function dependence of the KSAW wave near resonance. Thus, the wave propagates radially for $r \le r_s$ (hot edge) and for $r \ge r_s$ (cold edge), with wave damping determined by the nonadiabatic electron dissipation $\delta_{e\mathbf{k},\omega}$. For a collisionless hot edge plasma ($k_\parallel v_{Te} \gg \nu_e$)

$$\delta_{e\mathbf{k},\omega} = \sqrt{\pi}\left(\frac{\omega}{|k_\parallel| v_{Te}}\right)\exp\left(\frac{-\omega^2}{k_\parallel^2 v_{Te}^2}\right), \tag{2a}$$

and for a collisional hot edge plasma ($\nu_e > k_\parallel v_{Te}$)

$$\delta_{e\mathbf{k},\omega} = \frac{\omega k_\parallel^2 v_{Te}^2/\nu_e}{\omega^2 + (k_\parallel^2 v_{Te}^2/\nu_e)^2} \,. \tag{2b}$$

For a cold edge plasma,

$$\delta_{e\mathbf{k},\omega} = \frac{\omega}{\nu_e} + 2\sqrt{\pi}\left(\frac{\omega}{|k_\parallel| v_{Te}}\right)^3 \exp\left(\frac{-\omega^2}{k_\parallel^2 v_{Te}^2}\right).$$

The evolution of the average flow $\langle V_\theta(r)\rangle$ driven by the KSAW is determined by

$$\frac{\partial\langle V_\theta\rangle}{\partial t} + \mu\langle V_\theta\rangle = \frac{d\langle P\rangle_T}{dr} \,, \tag{3a}$$

$$\langle P\rangle_T = -\langle \tilde{V}_r \tilde{V}_\theta\rangle + \langle \tilde{B}_r \tilde{B}_\theta\rangle/4\pi\rho_0 \,. \tag{3b}$$

In this cylindrical model, $\langle V_\theta(r)\rangle = -c\langle E_r(r)\rangle/B_0$. Here μ is the neoclassical damping factor [18] [for the Pfirsch-Schlüter regime, relevant for tokamak edges, $\mu \simeq (a/R)^2 \omega_{T_i}^2/\nu_{ii}$, where R/a is the aspect ratio, ω_{T_i} is the ion transit frequency, and ν_{ii} is the 90° ion-ion collision time] and $\tilde{\mathbf{V}}, \tilde{\mathbf{B}}$ are KSAW fluctuations. $\langle P\rangle_T$ is the quasilinear Reynold's stress, determined by the difference of the fluid ($\langle \tilde{V}_r \tilde{V}_\theta\rangle$) and magnetic ($\langle \tilde{B}_r \tilde{B}_\theta\rangle$) stresses. Alternatively, $\langle P\rangle_T$ may be thought of as a ponderomotive pressure, driven by the self-beats of the KSAW. Equation (3) indicates that radially propagating waves, with radially varying wave energy density flux, are required for flow drive [19]. Equation (1) clearly implies that the KSAW satisfies both requirements. Furthermore, Eq. (3) implies that an imbalance between the fluid and magnetic stresses is required as well. Since, for the KSAW, $\hat{\psi}_{\mathbf{k}} = (1 + \Delta k_r^2 \rho_s^2)\hat{\phi}_{\mathbf{k}}$, where $\hat{\psi}_{\mathbf{k}} = \omega \hat{A}_{\parallel\mathbf{k}}/ck_\parallel$, ion inertia, via the polarization drift, provides such an imbalance. Thus

$$\langle P\rangle_T = k_r^R k_\theta \rho_s^2 c_s^2 \left|\frac{e\hat{\phi}(r)}{T_e}\right|^2 k_r^2 \rho_s^2 \Delta \,, \tag{4}$$

and, for steady flows (near the resonance),

$$\langle V_\theta(r)\rangle = \delta_{e\mathbf{k},\omega}\Delta^2 \left(\frac{k_\theta c_s^2}{\mu}\right)\frac{\omega^2|\tilde{\varepsilon}_r(r_s)|^2}{k_\theta^2 \rho_s^2 c_s^2}\left(\frac{\omega^2}{k_\parallel^2 v_A^2} - 1\right)^2 \exp(-2k_i^R|r-r_s|) \,. \tag{5}$$

Here Δ^2 accounts for the structure of the Airy function and the nature of the plasma response (cold or hot, etc.). We note here that details of whether the wave propagates to $r=0$ or reflects, or of the Airy function, do not qualitatively alter the simple physics of the flow drive mechanism. The result of Eq. (5) follows from the balance of the gradient of the ponderomotive pressure $d\langle P\rangle_T/dr$ with the neoclassical damping. Also,

$$\tilde{V}_r = -\rho_s c_s \nabla_\theta (e\hat{\phi}/T_e) = \partial\tilde{\varepsilon}_r(r_s)/\partial t$$

has been used to eliminate $e\hat{\phi}/T_e$ in favor of the KSAW displacement at the resonant surface, $\tilde{\varepsilon}_r(r_s)$. Finally, Eq. (5) is valid only for $r < r_s$ (hot edge) and $r > r_s$ (cold edge).

The structure of the flow field predicted by Eq. (5) has several interesting features. First, it should be noted that flow is driven by the electromagnetic torque induced by the quasilinear Reynold's stress and not by particle resonance with directed waves. Thus, the flow drive mechanism is macroscopic in nature, and does not involve modification of the microscopic structure of the ion distribution function. Second, from Eq. (1) $\langle V_\theta(r)\rangle \sim (k^{R2}\rho_s^2)^2\Delta^2$, consistent with the fact that ion inertia determines the imbalance between fluid and magnetic stresses. Third, since $k_i^R = (\delta_{e\mathbf{k},\omega}/2)k_r^R$, electron dissipation determines the width of the velocity shear layer, and is clearly necessary for a stationary mean flow. Note that the shear layer *width* will tend to *increase* with *decreasing* $\delta_{e\mathbf{k},\omega}$ and with increasing T_e in a hot edge plasma. Fourth, $\langle V_\theta(r)\rangle \sim k_\theta$, so that varying the poloidal orientation of the launched wave controls the direction of rotation. Finally, it is clear from Eq. (5) that the density profile, as well as the wave damping, contributes to determining the profile of $\langle V_\theta(r)\rangle$.

It is interesting to compare the results described above to those obtained from a consideration of the viscoresistive Alfvén wave (where $\omega \ll \nu_e$ and $k_\parallel v_{Te} \ll \nu_e$) in an infinite medium. Indeed, the problem of flow drive by these waves has much in common with viscoresistive current drive or the wave-driven dynamo, a topic of considerable present interest [10,20]. For these waves

$$k_{\parallel}=\frac{\omega}{v_A}\left[1+\frac{i(\nu+\eta)k_{\perp}^2}{2\omega}\right], \tag{6a}$$

$$\tilde{V}_{\mathbf{k}}=\frac{-k_{\parallel}v_A^2}{(\omega+i\nu k_{\perp}^2)}\left[\frac{\tilde{B}_{\mathbf{k}}}{B_0}\right], \tag{6b}$$

where ν is the kinematic viscosity. Equation (6a) expresses the fact that resistive and viscous dissipation damp Alfvén waves as they propagate along magnetic field lines. Using Eqs. (6a) and (6b), a short calculation yields

$$\langle P\rangle_T=-\left[\frac{(\eta-\nu)k_{\perp}^2}{2\omega}\right]^2 v_A^2\left\langle\frac{\tilde{B}_r}{B_0}\frac{\tilde{B}_\theta}{B_0}\right\rangle. \tag{7}$$

Note that, here, the difference between resistivity and viscosity, rather than ion inertia, determines the imbalance between fluid and magnetic stress. Indeed, the result of Eq. (7) is quite reminiscent of previous calculations of Alfvén-wave-driven dynamo activity [21] and current drive [10]. In the former case, the net "α effect," determined by the difference of fluid and magnetic helicities, vanished for equipartitioned, equialigned, and equidissipated 3D magnetohydrodynamic (MHD) turbulence. In the latter case, a circularly polarized Alfvén wave was shown to drive a quasilinear toroidal electric field of the form $\langle E_z\rangle=[(\eta-\nu)k_{\perp}^2/\omega]|\tilde{B}_r|^2$. Also, in this case, resistive and viscous dissipation, rather than electron Landau damping, provides the necessary "phase shift." It is also amusing to observe that while current drive is maximal for circular polarization, viscoresistive Alfvén-wave flow drive is maximal for plane polarization. Finally, it is readily apparent from Eq. (7) that dissipation and radial localization are necessary for viscoresistive Alfvén-wave flow drive.

The goal of flow drive is to suppress turbulence, thereby improving confinement. Thus, electric-field shear must be large enough that the suppression criterion [5]

$$k_\theta'\Delta x_{\mathbf{k}'}\frac{d\langle V_\theta(r)\rangle}{d}=k_\theta'\Delta x_{\mathbf{k}'}\frac{c}{B_0}\frac{d}{dr}\langle E_r\rangle>c\Delta\omega_{\mathbf{k}'} \tag{8}$$

is satisfied. Here $\Delta\omega_{\mathbf{k}'}$, k_θ', and $\Delta x_{\mathbf{k}'}$ refer to the poloidal wave number, radial correlation length, and decorrelation rate of the ambient edge turbulence. c is a parameter which accounts for the fact that the criterion $k_\theta'\Delta x_{\mathbf{k}'}\langle V_\theta\rangle'\geq\Delta\omega_{\mathbf{k}'}$ overestimates the velocity shear needed for turbulence suppression [22], i.e., $c\leq 1$. Relating $|\tilde{\varepsilon}_r(r_s)|^2$ to absorbed power via [23]

$$P_{\text{abs}}=\frac{\pi}{4}aR\frac{|\tilde{\varepsilon}_r(r_s)|^2}{k_\theta^2}\frac{\omega d\epsilon_r}{dx}, \tag{9a}$$

$$\epsilon_r=4\pi\rho_0\omega^2-k_{\parallel}^2B_0^2, \tag{9b}$$

it follows that the criterion for suppression is

$$S(r)\geq 1, \tag{10a}$$

where, under the validity constraints of Eq. (5),

$$S(r)=\left[\frac{k_\theta'\Delta x_{\mathbf{k}'}}{c\Delta\omega_{\mathbf{k}'}}\right]\left[\frac{2\delta_{e\mathbf{k},\omega}^2}{\pi^2}\right]\left[\frac{\omega_{ci}}{\omega}\frac{k_\theta c_S}{\mu}\frac{L_n}{\rho_s}\right]\left[\frac{P_{\text{abs}}}{\rho_0 aR\rho_s c_S^2}\right]\Delta^{5/2}\left[\frac{\omega^2}{k_{\parallel}^2v_A^2}-1\right]^{5/2}\exp(-2k_i^R|r-r_s|). \tag{10b}$$

In deriving Eq. (10b), $d\epsilon_r/dx\cong 4\pi\rho_0\omega^2/L_n$ was used, since $1/L_n\gg|k_{\parallel}^{-1}dk_{\parallel}/dr|$ for low-m, edge-resonant KSAW's. Note that the suppression criterion $S(r)\geq 1$ determines the amount of absorbed power (P_{abs}) necessary for an enhanced confinement zone of given width. For example, in the case of drift-wave turbulence ($\Delta\omega_{\mathbf{k}}\sim\omega_{*e}$, $k_\theta'\Delta x_{\mathbf{k}'}\sim 1$, $k_\theta'\rho_s\sim 0.2$) in the cold edge of TEXT ($B_0=2$ T, $n_0=5\times 10^{12}$ cm^{-3}, $T_e=30$ eV), $P_{\text{abs}}\sim 300$ kW is required for a 8-cm-wide zone of good confinement. Similar calculations in the marginally cold (almost hot) edge of DIII-D in the L mode ($B_0=2.1$ T, $n_0=10^{13}$ cm^3, $T_e=150$ eV) indicate that $P_{\text{abs}}\sim 300$ kW will produce a 5-cm-wide turbulence-suppression zone. It is interesting to note that the power requirement is determined primarily by the desired *width* of the suppression zone, and *not* by the *peak value* of the velocity shear. This is a consequence of the fact that

$$d\langle V_\theta(r)\rangle/dr\sim\delta_{e\mathbf{k},\omega}^2\exp(-2k_i^R|r-r_s|),$$

where $k_i^R=(\delta_{e\mathbf{k},\omega}/2)|k_r^R|$ and

$$k_r^R=(\Delta^{1/2}\rho_s)^{-1}(\omega^2/k_{\parallel}^2v_A^2-1)^{1/2}.$$

Hence, in an edge plasma with weak electron dissipation, the rf-driven shear layer tends to be strong but narrow, while in an edge plasma with strong electron dissipation, the shear layer is somewhat weaker, but broader at comparable P_{abs}. Thus, flow drive and turbulence suppression will be easier in future hotter (than DIII-D) edge plasmas. Also, as confinement improves temperatures will rise and the density gradient will steepen, thus broadening the width of the shear layer further [see Eq. (10b)]. Anomalous momentum transport processes active at the edge of the layer can be expected to broaden it as well. Indeed, a quantitative understanding of the temporal evolution of the shear layer structure during flow drive requires transport analysis, including self-consistent evolution of the angular momentum flux, as well as modeling of the effects of electric-field shear on transport. However, the parameter scalings of P_{abs} required for $S(r)\geq 1$ may be determined from Eq. (10b).

Several questions naturally arise concerning the practicality of this proposal. First, the external rf may directly drive anomalous transport processes or, via heating, aggravate ambient turbulence, thereby defeating shear suppression. Regarding the former, it should be noted that thermal electrons resonate with the KSAW when $k_{\parallel}=\omega/v_e < k_{\parallel}(r_s)$. Since $q(r)$ increases monotonically with radius in tokamaks, the resonant radius for thermal transport lies *outside* the Alfvén resonance radius r_s in a hot edge. Thus, while low-energy particles resonant with the wave when $k_{\parallel}(r)=\omega/V_{\parallel} > \omega/V_{Te}$ may be lost, thermal particles will be well confined. On the other hand, the cold-edge-induced resonance transport is small in the parameter v_e/ω (resonant transport is exponentially small). It should also be noted that the values of $P_{abs} \sim 200$–300 kW are almost always smaller than the total input heating power (Ohmic and auxiliary). Hence, the power transported to the edge region in steady state will also exceed the flow-drive power, rendering its associated heating effects irrelevant. Finally, since $\omega >$ diamagnetic frequency, the transport is *inward* and helpful to confinement.

A second concern pertains to the effect of the large ambient density fluctuations on the KSAW resonance. However, it should be noted that, if successful, flow generation will naturally tend to heal such resonance smearing by suppressing the turbulence which causes it. Further, ambient edge turbulence has much smaller decorrelation rates than the wave frequency, rendering it a static scatterer of the KSAW. Moreover, the KSAW wavelength is large and the propagation path is short so that phase fluctuations are small. As a result, the KSAW resonance should be relatively sharp for high P_{abs}.

It is also worthwhile to note that while rf flow drive is a new concept, experiments with edge-resonant Alfvén heating on the TCA tokamak have already noted a change in the observed propagation direction of low-frequency fluctuations [24]. This change is consistent with the appearance of an rf-driven radial electric field and flow. In this regard, an abrupt change in the measured plasma potential and phase velocity direction over a few cm was observed during heating.

In conclusion, shear suppression of edge turbulence utilizing rf-driven flows is an interesting and potentially viable technique for enhancing the performance and understanding of tokamak confinement, and thus merits further study. In addition to treatment of the cold-edge standing wave [17] and compressional effects, inclusion of realistic geometric and profile effects, and the development of a gyrokinetic theory of momentum transport, future work will focus on consideration of flow drive using ion-cyclotron-range-of-frequency and ion-Bernstein waves. Indeed, in addition to being more practical and economical, utilization of high-frequency waves for flow drive avoids the competition between fluid and magnetic stress intrinsic to Alfvén waves.

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