A review of experimental drift turbulence studies

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Abstract

Experimental drift turbulence and zonal flow studies in magnetically confined plasma experiments are reviewed. The origins of drift waves, transition to drift turbulence and drift turbulence–zonal flow interactions in open field line and toroidal closed flux surface experiments are discussed and the free energy sources, dissipation mechanisms and nonlinear dynamics of drift turbulence in the core, edge and scrape-off layer plasma regions are examined. Evidence that turbulence across these regions is linked and that turbulence-driven zonal flows exist is presented, and evidence that these flows help regulate the turbulent scale lengths, amplitude and fluxes is summarized. Seemingly contradictory reports exist regarding the scale of turbulent transport events; gyro-Bohm behavior of turbulence correlation lengths as well as evidence for long-range transport phenomena both exist. Changes in turbulence during and after transport barrier formation are summarized and compared. The inferred turbulent particle and heat fluxes due to turbulent transport are usually consistent with global confinement, and edge plasma momentum transport appears to be linked to plasma flows at the last-closed flux surface and in the open field line region. However, inconsistencies between observed transport and turbulence have sometimes been reported and are pointed out here. Special attention is given to open issues, and suggestions for future experimental studies are given.

(Some figures in this article are in colour only in the electronic version)

1. Background

Since the advent of magnetic fusion research, measurements of volume-averaged particle and energy confinement have been carried out using a variety of approaches to measure the transport. Particle and power balance studies showed that the effective transport coefficients are much larger than what would be obtained from purely collisional neoclassical transport processes (figure 1); the discrepancy is attributed in large part to the presence of turbulent transport processes in the plasma. This review paper provides a summary of the current...
1.1. Early evidence of drift turbulence in confined plasmas

Early evidence for the presence of drift turbulence in confined plasmas came from collective Thomson scattering of microwaves and far infrared radiation in the ATC and PDX tokamaks as reported by Mazzucato and independently by Surko and Slusher [1–4]. Later work showed that the observed broad spectra (see figure 2) were indeed due to a broad turbulent spectrum and were not due to a spatially varying Doppler shift that might occur if the plasma rotation
rate varied with minor radius [5], which could produce an artificially broadened spectrum due to the poor spatial resolution of the scattering diagnostics used in these studies. A summary of these early measurements of turbulence in tokamak plasmas is available in the literature [6]. In these early studies, the measured density fluctuation amplitude was shown to be roughly consistent with a mixing length estimate \( \tilde{n}_{\text{rms}}/n_0 \sim 1/k_\perp L_n \), where the \( \tilde{n}_{\text{rms}} \) denotes the root-mean-squared fluctuation amplitude of the density, \( n_0 \) denotes the mean density, \( k_\perp \) denotes the perpendicular wavenumber of the fluctuations and \( L_n \) denotes the density gradient scale length. This limit arises when the turbulence momentarily flattens the total instantaneous gradient, thereby momentarily stabilizing the linear instability. Estimated transport diffusivities from the turbulence were close to the values needed for global confinement, bolstering the thinking that the turbulence was somehow related to the global confinement.

Shortly after these results were reported, studies of the edge and scrape-off layer (SOL) region in several small tokamaks using multi-point Langmuir probe arrays also showed that large-amplitude (\( \tilde{n}_{\text{rms}}/n_0 \sim 0.1–1 \)) density fluctuations existed in the edge plasma region, that these fluctuations could cause a turbulent cross-field particle flux that was consistent with global confinement in the plasma and that the spectra of these fluctuations were also rather broad as expected for a turbulent plasma [7–13].

1.2. Summary of drift turbulence theory

We summarize the essential elements of the drift wave instability making reference to an existing theory review [14]. The focus here is primarily upon electrostatic disturbances; we provide a summary of expected electromagnetic effects afterwards. Elementary considerations show that in the presence of ion and electron pressure gradients, a magnetized plasma develops corresponding ion and electron diamagnetic currents that provide for the plasma equilibrium. If a small disturbance (e.g. a ‘wave’) in the ion and/or electron pressure gradient should occur, then the corresponding diamagnetic currents will develop a perturbed response. This perturbed current is perpendicular to the magnetic field and will be carried by the ions whose guiding centers move with the ion polarization drift velocity associated with the perpendicular wave electric field. For the spatio-temporal scales of interest in this review the plasma is quasineutral and thus the electrostatic charge density of the wave is small, and the divergence of the total perturbed current must then vanish. This constraint then forces a corresponding perturbed parallel current to develop. Due to their small inertia, electron motion parallel to the magnetic field effectively generates this current. The resulting perturbations propagate predominantly in the ion and electron diamagnetic drift directions, and are known as ion and electron drift waves due to their close relationship to the diamagnetic particle drifts.

In the absence of dissipation of the parallel electron motion (a condition which is usually denoted as the adiabatic limit in the jargon of drift wave theory), the resulting drift wave density fluctuations will be in phase with the wave plasma potential fluctuations, i.e. \( \tilde{n}/n \approx e\tilde{\phi}/kT_e \) (a relation sometimes referred to as the Boltzmann relation for the electron response). If the electrons can lose momentum to the background plasma as they move parallel to the magnetic field, the resulting dissipation creates a time delay between the plasma potential perturbation and the resulting electron density perturbation that then follows. As a result the potential perturbations are phase shifted relative to the density fluctuations, with \( \tilde{n}/n = (e\tilde{\phi}/kT_e)(1 - \delta) \) where \( \delta \neq 0 \). The dissipation of parallel electron motion can occur via several different processes including electron–ion Coulomb collisions, wave–particle interactions in a collisionless plasma or via collisions of passing electrons with either trapped electrons or trapped ions in a torus. The dominant mechanism depends upon the particular conditions of the experiment being studied.
The top panel of figure 3 [14] illustrates the resulting dynamics in the so-called adiabatic limit without parallel electron dissipation. In this case, a density perturbation gives rise to a potential perturbation that is in phase with the density fluctuation. Because of the low-frequency nature of the drift instability, the resulting electric field gives rise to an $E \times B$ drift of the ions and electrons. The resulting particle motion does not give rise to a time-averaged cross-field flux and convects the disturbance in the diamagnetic drift direction which satisfies a dispersion relation $\text{Re}(\omega(k)) = 0$ with a vanishing growth rate $\text{Im}(\omega(k)) = \gamma_k = 0$. When finite dissipation of the parallel electron momentum is included, the potential perturbations are phase shifted relative to the density fluctuations, with $\tilde{n}/\bar{n} \approx (e\tilde{\phi}/kT_e)(1+i\delta)$ where $\delta \neq 0$ resulting in the situation shown in the lower panel of figure 3. In this case, the disturbance is linearly unstable, resulting in a growing amplitude with a growth rate $\text{Im}(\omega(k)) = \gamma_k > 0$ for a limited range of wavenumbers. In this case the time-averaged product of the fluctuating density and $E \times B$ velocities does not vanish, and as a result the fluctuations can cause a net transport of plasma down the mean gradients and the perturbation then grows in amplitude. The figure illustrates the development of a turbulent particle flux in an isothermal plasma; a similar process where the density fluctuation is replaced by a temperature fluctuation will develop in a plasma with a temperature gradient, leading to the transport of heat down the background gradient.

As the drift wave perturbations reach a finite amplitude, they can interact nonlinearly with perturbations at other wavenumbers due to the presence of the convective derivative in the fluid conservation equations. This leads to the development of disturbances at other wavenumbers some of which are linearly stable and are thus damped. In the limit of very weak parallel electron dissipation, this process is often described in the Fourier domain via a set of mode
coupling equations for the power spectrum evolution of the form

\[
\frac{\partial \langle \phi_k^2 \rangle}{\partial t} = \gamma_k \langle \phi_k^2 \rangle + \text{Re} \sum_{k_1, k_2} \Lambda_{k_1, k_2} \langle \phi_{k_1}^* \phi_{k_1} \phi_{k_2} \rangle,
\]

where \(\langle \phi_k^2 \rangle\) denotes the ensemble-averaged fluctuation power spectrum for wavenumber \(k\), \(\Lambda_{k_1, k_2}\) denotes a coupling coefficient that is a function of wavenumbers \(k_1\) and \(k_2\) and the third order quantity \(\langle \phi_{k_1}^* \phi_{k_1} \phi_{k_2} \rangle\) denotes the complex bispectrum which determines the phase coherency and effective transfer of energy between Fourier modes. The matching condition \(k = k_1 + k_2\) can be seen to arise from the requirement that the wave momentum be conserved in this interaction. These new scales in turn interact and produce fluctuations on yet other scales. Eventually this progression saturates into a condition where the release of free energy is balanced by the nonlinear transfer of that energy to other scales where it is damped. The result is a saturated turbulent spectrum of density, temperature and velocity fluctuations. In the simple model above, this would then give a stationary power spectrum given as

\[
\langle \phi_k^2 \rangle = -\frac{1}{\gamma_k} \text{Re} \sum_{k_1, k_2} \Lambda_{k_1, k_2} \langle \phi_{k_1}^* \phi_{k_1} \phi_{k_2} \rangle.
\]

In configuration space, there is a distribution of turbulent eddies which have a distribution of finite lifetimes and spatial scales, and which have density, temperature and potential fluctuations associated with them. The term in the summation is sometimes referred to as the nonlinear transfer term. Attempts have been made to measure this term in experiment. When parallel dissipation is strong, then a coupled set of equations with a similar form arise and describe the nonlinear transfer of pressure fluctuations and turbulent kinetic energy. Theory predicts that fields such as the density will transfer energy toward higher wavenumber while kinetic energy will transfer toward smaller wavenumber, forming large-scaled structures in the plasma [15].

The fluctuating density, temperatures and plasma velocities that develop from the drift turbulence have cross-correlations which lead to time-averaged fluxes of these quantities, as can be seen by simple considerations [16]. In the electrostatic limit, when magnetic field fluctuation effects can be neglected, the primary guiding center drift is due to fluctuating \(E \times B\) drifts, which when correlated with the other fluctuating quantities give rise to a time-averaged electrostatic turbulent flux of particles \(\bar{T}\), momentum \(\bar{\mu}\) and heat \(\bar{Q}\) given by

\[
\bar{T} = -\frac{\langle n \nabla \phi \times \hat{B} \rangle}{B^2} + \langle n \bar{v}_|| \rangle \hat{b}_0,
\]

\[
\bar{\mu}_R = \left\langle \left( -\frac{\nabla \phi \times \hat{B}}{B^2} + \bar{v}_i \hat{b}_0 \right) \left( -\frac{\nabla \phi \times \hat{B}}{B^2} + \bar{v}_i \hat{b}_0 \right) \right\rangle
\]

and

\[
\bar{Q} = \frac{5}{2} \bar{n} \left[ \frac{1}{T} \left( -\frac{\langle \bar{T} \nabla \phi \rangle \times \hat{B}}{B^2} + \langle \bar{T} \bar{v}_|| \rangle \hat{b}_0 \right) + \frac{1}{n} \left( -\frac{\langle \bar{n} \nabla \phi \rangle \times \hat{B}}{B^2} + \langle \bar{n} \bar{v}_|| \rangle \hat{b}_0 \right) \right].
\]

Here the fluctuating density, potential, parallel velocity and temperature are given by \(\bar{n}, \bar{\phi}, \bar{v}_||\) and \(\bar{T}\), respectively, and \(\langle \rangle\) denotes a time average or ensemble average. The resulting particle, momentum and heat fluxes are then functions of position. Due to its significance in fluid turbulence, the momentum flux \(\bar{\mu}\) is also referred to as the Reynolds stress. The
elements of this turbulent stress tensor are composed of the ensemble-averaged products of the fluctuating velocities, e.g. $\langle \tilde{v}_r \tilde{v}_\theta \rangle$, $\langle \tilde{v}_r \tilde{v}_\parallel \rangle$, and so forth. Here $r$ denotes the radial direction that is antiparallel to the mean pressure gradient, $\theta$ denotes the azimuthal or poloidal direction and the parallel symbols refer to the magnetic field direction.

When the plasma pressure is high enough, the inductive electric field becomes significant and the fluctuating parallel current induces a perturbed magnetic field resulting in substantial changes in the drift wave characteristics [17]. When $\beta < m_e/M_i$ the Alfvén wave emerges as a second eigenmode at a much higher frequency than the drift wave. When $\beta > m_e/M_i$ the drift wave and Alfvén wave dispersion relations merge into a single mode referred to as the drift-Alfvén wave that has characteristics of both types of waves. When the drift-Alfvén wave becomes turbulent via a mode coupling process, the resulting correlations between fluctuating fields can lead to an electromagnetic component of the particle, momentum and heat fluxes. Theory has suggested that it may be possible for such fluctuating currents to destroy magnetic flux surfaces and form volumes filled with stochastic field lines, leading to large rates of electron heat transport due to parallel motion along magnetic fields [18]. This represents a possible additional loss mechanism that does not involve the drift turbulence $E \times B$ drift physics discussed above. The magnetic fluctuations can also induce a stress on the plasma resulting in a transport of momentum. The turbulent momentum flux and turbulent electron heat flux caused by magnetic field fluctuations are given as

$$\mu_{\text{EM}} = \frac{\langle \tilde{B} \tilde{B} \rangle}{\mu_0 n M_i}$$

and

$$\tilde{Q}_{r\parallel}^{\text{EM}} = \frac{\langle \tilde{q}_{r\parallel} \tilde{B}_r \rangle}{\tilde{B}}$$

respectively [18, 19]. These fluxes in general would exist in addition to the electrostatic fluxes given earlier.

There are measurements of some of these turbulent fluxes in the edge and SOL region of confinement experiments, as well as in laboratory plasmas. There are also a few reports of turbulent transport fluxes deeper in the plasma, usually made with probes in smaller, lower temperature confinement devices, or with heavy ion beam probe (HIBP) systems in hotter, higher powered confinement experiments [20]. Since the cross-correlation function forms a Fourier transform pair with the cross-power spectrum, these quantities can also be written in terms of the fluctuation amplitude, cross-phase and cross-coherence computed from the correlated fluctuating quantities as has been shown by Powers for the particle flux [21]. These points are discussed in more detail below in the section on data analysis techniques.

When they are measured or inferred, the turbulent particle and heat fluxes are usually compared against the plasma confinement in one of two ways. In the first approach a careful particle and power balance is carried out. The particle and heat sources and sinks are either measured or inferred from models, and then these terms are used in the particle and heat conservation equations of the species in question to infer the flux-surface-averaged fluxes. These so-called transport fluxes are then compared against the turbulent fluxes to determine the significance of the turbulent transport to overall confinement. The earlier review of Wootton et al provides an excellent summary of this approach and of the resulting implications in ohmically heated and beam heated tokamaks [22].

In the second approach, the total flux across a specific magnetic flux surface of interest (usually the last-closed flux surface (LCFS)) is calculated from measurements, usually with an assumption about the poloidal and toroidal symmetry of the flux. When coupled with the
volume integrated source and sink, this then yields an estimate for the confinement time, e.g. for the electron energy balance one obtains

\[ \tau_E \equiv \frac{\int_V P_e(x) \, d^3x}{\int_A Q_e(x) \cdot ds} \]

A similar expression can be written for the particle confinement time in terms of the particle flux and particle source and, in principle the ion energy confinement as well (although in practice this latter quantity has never been directly assessed despite the development of measurements of ion temperature fluctuations [23]). This so-called turbulent confinement time is then compared against the globally measured energy or particle confinement time.

A brief word on nomenclature: the plasma in magnetic fusion experiments is usually considered to consist of a core plasma region, the LCFS which forms the boundary separating open and closed field line regions, the SOL which is the open field line region and then an edge plasma region that interfaces the core region with the SOL. The designation between the core and edge is somewhat arbitrary but this distinction is used throughout the literature and thus we will make use of this nomenclature, keeping in mind that the regions are connected to and interact with each other. It is also important to note that, when considering the SOL region, the effects of the sheath which surrounds the divertor or limiter surfaces must be considered. As has been discussed in the literature [24–27], the sheath introduces a new resistive term to the parallel electron motion and the associated parallel current. In the SOL region the sheath dissipation of the parallel electron motion often dominates the volumetric collisional dissipation. In this case the sheath physics may become a significant contributor to the turbulence in the SOL region, influencing both the amplitude of the turbulence and the cross-phase between density and velocity or potential fluctuations.

1.2.1. Drift turbulence characteristics. In the direction perpendicular to the magnetic field the drift turbulence is characterized by the spatial scale \( \rho_S \equiv \frac{C_S}{\Omega_i} \) where \( C_S \) denotes the ion sound speed and \( \Omega_i \) the ion gyro-frequency. The drift wave parallel wavenumber satisfies \( k_\parallel \ll k_\perp \rho_S \); the range of \( k_\perp \rho_S \) depends upon the linear free energy source and dissipation mechanism as well as upon the nonlinear energy transfer. In the plasma fluid frame the electron (ion) pressure gradient driven modes are expected to propagate in the electron (ion) diamagnetic direction. When drift turbulence fluctuation measurements are made, these signatures or characteristics are usually examined and an attempt is usually made to identify the underlying free energy source and dissipation mechanisms. With the advent of numerical turbulence simulations, recent efforts have also begun to compare the model turbulence characteristics directly with experiment in order to more deeply test the theoretical understanding of the turbulence origins, saturation mechanisms and associated transport. These comparisons can also include multi-field effects and turbulence/zonal flow (ZF) interactions in addition to tests of the linear stability of the mode in question.

The collisional plasmas found in laboratory experiments with low (\( \sim \)few eV) electron temperatures usually exhibit resistive drift wave turbulence driven by the electron pressure gradient and where electron–ion collisions provide the dissipation. In a few cases, such experiments have also investigated the so-called universal drift instability, where the fluctuating parallel electron motion is inhibited by inverse-Landau damping of the drift wave. In smaller tokamak experiments where ohmic heating predominantly heats the electrons, the turbulence has been attributed to this universal instability or to the dissipative trapped electron mode. This latter instability is driven by the electron density gradient free energy source; parallel electron dissipation is then provided by Coulomb collisions between the trapped and passing electrons. Evidence for the existence of these drift instabilities in small confinement experiments has
been summarized in Wootton et al. [22] and in the references contained therein. At the point of that earlier writing it was not possible to clearly attribute the observed transport to any one of these classes of drift instabilities.

In larger confinement devices with strong auxiliary heating and/or with high enough densities such that significant ion heat transport occurs, the ion temperature gradient (ITG) can become large enough so that the inequality $\eta_i \equiv \partial_r \ln T_i / \partial_r \ln n > \eta_{\text{crit}}$ is satisfied, resulting in the ITG drift wave instability. There are several different linear branches of this mode having to do with the detailed geometry (e.g. slab, toroidal) and local values of collisionality, magnetic shear and so forth. The resulting ITG turbulence is expected to have $k_\perp \rho_S \sim 0.1–0.5$ and to propagate in the plasma frame with a phase velocity close to the ion diamagnetic drift velocity $V_{\text{ph}} \approx V_{\text{dia}}$. The precise value of $\eta_{\text{crit}}$ depends upon whether one is examining the slab or toroidal branches of the instability as well as upon the impurity gradients, but generally has a value in the range $\eta_{\text{crit}} \sim 1$. This stability threshold is often recast in terms of the ITG normalized to the major radius, i.e. to the dimensionless ratio $R/L_{T_i} > R/L_{T_i \text{crit}} \sim 3–5$. The precise value of the linear $R/L_{T_i \text{crit}}$ threshold depends upon the assumptions in the model and upon detailed plasma conditions.

Another instability that can arise in plasmas with strong electron heating is the collisionless trapped electron mode (TEM) which is driven by the electron pressure gradient and destabilized by wave–particle interactions between the toroidal precession of the trapped electrons and the parallel phase velocity of the TEM drift wave. The resulting turbulence can range from ITG scales down to smaller scales with $k_\perp \rho_S \sim 1$, and in the plasma frame is expected to propagate in the electron diamagnetic drift direction. Depending upon the plasma conditions, the linear mode can be unstable for a wide range of electron pressure gradient, or can exhibit a critical gradient behavior much like that of the ITG mode. As the TEM mode wavenumber increases such that $k_\perp \rho_S > 1$ the TEM mode transitions into the electron temperature gradient (ETG) drift instability. The ETG instability occurs when $\eta_e \equiv \partial_r \ln T_e / \partial_r \ln n > \eta_{\text{crit}} \sim 1$ and is expected to lead to fine-grained turbulence with $k_\perp \rho_S \sim 1–10$ that in the plasma frame propagates in the electron diamagnetic drift direction. The stability threshold can also be cast in terms of the dimensionless ratio $R/L_{T_e} > R/L_{T_e \text{crit}} \sim 3–5$.

For a more complete discussion of these instabilities and the individual characteristics of each, as well as for a comparison with confinement properties of tokamak devices the reader is referred to the paper by Doyle et al. [28] and references therein. Several older theory reviews also summarize the free energy sources, dissipation mechanisms and key linear stability properties of drift wave instabilities thought to be relevant in fusion confinement experiments [14, 29, 30]. The main point to note here is a recognition that the type of drift instability within a single experimental device can vary depending upon the gradients and local conditions of a particular experiment, and that the various classes of drift instabilities are usually characterized by their perpendicular wavenumber ranges, propagation directions and free energy sources.

1.2.2. Development of sheared ZFs from drift turbulence. In 1979, Hasegawa et al. studied the nonlinear dynamics of drift wave and Rossby wave turbulence using numerical simulations of collisional drift turbulence [31]. They found that the drift turbulence kinetic energy with perpendicular wave number in the range $0.1 < k_\perp \rho_S < 1$ condensed into an ordered anisotropic flow characterized by a scale comparable to the turbulence scale length in the direction parallel to the background density gradient scale length (i.e. the radial wavenumber of this flow $k_\perp \rho_S \sim 0.1–0.3$) while the scale length in the direction transverse to both the magnetic field and the background density gradient was long (i.e. in the azimuthal (poloidal) wavenumber of this flow $k_\theta \rho_S \sim 0$). Eventually closed iso-potential contours, which correspond to the
streamlines of the ZF $E \times B$ drifts, developed, indicating the formation of a region in which turbulent mixing across the barrier was reduced (see figure 4). No such contour structure developed in the density fluctuations. Due to the similarity to geophysical flows, they denoted this shear flow as a ZF. The emergence of this structure was similar to the generation of large-scale convective cells by drift waves which was discussed by earlier authors [32–34]; however, the key difference lay in the lack of cross-field transport from the ZF, while the convective cells were thought to give rise to rapid transport across the magnetic field due to their large spatial scales. The significance of this ZF on turbulence saturation and transport was not generally recognized until sometime later [35, 36].

The emergence of this sheared ZF arises due to the radial divergence of the cross-field poloidal or azimuthal momentum flux [37], which due to its importance and significance is known in the fluid turbulence literature as the turbulent Reynolds stress [37]. Within the fluid description, the Reynolds stress arises from decomposing the fluid fields into slowly varying and fluctuating components and then making the usual assumption that the time average of a single fluctuating component vanishes while the time average of higher order fluctuating products does not [19, 38]. When the ZF velocity $\bar{V}_{ZF}$ is damped at a rate $\nu_d$ then in a simple model (where $x$ denotes radius and $y$ denotes azimuthal or poloidal direction) the ZF velocity evolves according to the equation [37, 39]

$$\frac{\partial}{\partial t} \bar{V}_{ZF} = -\frac{\partial}{\partial x} \langle \bar{v}_x \bar{v}_y \rangle - \nu_d \bar{V}_{ZF}.$$

Since the time-averaged turbulent Reynolds stress $\langle \bar{v}_x \bar{v}_y \rangle$ is a flux of momentum, it acts to rearrange the momentum but does not act to inject or dissipate the momentum into the plasma. Thus a divergence of the momentum flux leads to either a concentration or diffusion of the momentum depending upon the sign of the divergence. The reader can also view the divergence of the stress as giving rise to a net force on the fluid element; for the simple model in the equation above, the force is in the $y$ direction and can either reinforce or damp $V_{ZF}$ for a negative (positive) stress divergence. Momentum sinks (due to either viscous damping effects or flow drag effects arising from trapped ion–passing ion or ion–neutral collisions) and sources then act to balance out the flux divergence and can give rise to both fixed point (FP) and dynamical behavior for the turbulence and ZF [40]. This shear flow drive from the Reynolds stress has also been viewed as the merging of anisotropic tilted vortices or eddies [41].

One can also view ZF generation as a nonlinear energy transfer process in the Fourier domain in which the turbulent kinetic energy is carried to larger spatial scales via the beating
of smaller scaled turbulent velocity fluctuations associated with the Reynolds stress which
then produces a fluctuation at large spatial scale [42,43]. Energy is conserved in this process,
leading to the ZF extracting energy from turbulence and thereby reducing fluctuation amplitude.
The ZF does not tap into free energy source because the associated flow velocity is normal
to the background gradients; thus it is linearly stable and does not cause transport. Due to
the importance of the subject an extensive theoretical literature has developed and has been
summarized in a recent review [39].

Because the sheared ZF has a finite correlation time, in the basic theoretical model the
increase in the turbulence radial wavenumber due to the effect of the ZF is expected to be
diffusive in wavenumber space with the change in radial wavenumber given as ⟨k^2⟩ ∼ D_k t
where the turbulent diffusion rate D_k ∼ k^2 (∩ZF)τ_c [39]. The net effect is a reduction in
turbulent radial correlation length, which is inversely proportional to the radial wavenumber
of the fluctuations and a decrease in the turbulence amplitude relative to what would be found in
the absence of the ZF. These two factors control the rate of turbulent transport and therefore the
ZF plays a critical role in determining the rate of transport from the associated drift turbulence.
In a cylindrical plasma the ZF has a zero-mean frequency component with a finite correlation
time proportional to the damping rate of the shear flow. In a torus, a second branch of the
ZF appears due to the geodesic curvature of the magnetic field, which gives rise to density
compressions as the plasma rotates poloidally with an azimuthally symmetric radially sheared
poloidal velocity. These compressions can then relax via a sound wave which propagates
along the magnetic field, giving rise to a finite frequency for this sheared flow with a frequency
C_s/R where R denotes the major radius [44]. This finite frequency sheared flow is known as
the geodesic acoustic mode (GAM). For recent in-depth reviews of the theoretical aspects to
ZF/turbulence interactions, the reader is referred to several recent publications and references
therein [45,46]. We also point out that a review of the status of experimental studies of ZFs
has recently become available [47]. These references also provide useful intuitive schematics
of the turbulence–shear flow interaction process.

1.2.3. Interaction of mean sheared flows with drift turbulence. In 1982, a regime of improved
confinement was reported in the ASDEX device [48] in which a steep edge density gradient
was observed to spontaneously form if sufficient plasma heating was applied and the device
wall conditions were sufficiently clean. This high confinement regime, denoted the H-mode
regime in order to distinguish it from the so-called low-mode confinement (L-mode) regime,
has subsequently been observed in nearly all medium and large-scale tokamak devices, as well
as in a number of stellarator devices. The H-mode regime showed that the plasma could support
the existence of a much steeper density gradient for the same value cross-field flux indicating
that the effective cross-field diffusivity had dropped in value. Subsequent measurements in
the DIII-D tokamak showed that the turbulent fluctuations just inside the plasma edge were
suddenly reduced in amplitude [49] and that at the moment of the transition between the two
states (denoted as the L–H transition), a spatially varying strong radial electric field associated
with a plasma $E \times B$ drift spontaneously developed [50]. Further evidence for the importance
of radial electric fields and plasma flows on the edge plasma turbulence came from observations
that the H-mode state could be induced by the generation of a strong $E \times B$ edge plasma flow
using a small biased electrode inserted into the plasma [51] and subsequent studies that showed
a large reduction in the turbulent transport during the application of the biasing [52–56].

Motivated by these observations, a basic theory of turbulence decorrelation by the action of
a sheared background mean $E \times B$ flow was developed [57,58]. This mean flow is distinguished
from the ZF by the fact that the mean flow can be sustained by non-turbulent processes such
as external momentum input or by the return current created by trapped ion orbit loss physics. This theory predicts that the turbulent fluctuation amplitude and radial correlation length should be reduced when the flow shear decorrelation rate $\omega_{sh} \approx k_\theta (\partial V_\theta / \partial r) / \Delta_r$ exceeded the natural decorrelation rate $\Delta \omega \approx 1 / \tau_{corr}$ of the turbulence [57, 58]. Here $k_\theta$ is the azimuthal wavenumber of the turbulence, $\partial V_\theta / \partial r = V'_\theta$ is the radial shear of the azimuthal $E \times B$ velocity, $\Delta_r$ is the radial correlation length of the turbulence, $\tau_{corr}$ is the decorrelation time of the turbulence and the turbulence-related quantities are to be evaluated with their unsheared values. A detailed review of the theory has been given elsewhere [59], and a comparison of primarily macroscopic experimental measurements with this model is also available [60]. Subsequent experiments showed that multi-field effects were important [53, 55, 56, 61, 62], and in particular that the shear flow has an important effect on the cross-phase (or equivalently on the zero-time delay cross-correlation) between the density and potential fluctuations, and therefore upon the time-averaged particle flux carried by the turbulence. Later modifications to the basic theory of flow shear decorrelation included two-field effects which gave a cross-phase delay that depends upon the flow shear decorrelation rate [63–65].

In these models of shear flow decorrelation effects the origin of the shear flow was left unspecified. In subsequent theoretical work, several possible trigger mechanisms were proposed including the turbulent Reynolds stress [37], Stringer spin-up [66] and neoclassical ion orbit loss [58, 67] which could act to induce a shear flow (for a review of these theoretical mechanisms of shear flow generation see [68]). This shear flow then inhibits transport which, given continued heat input into the core plasma, would then reinforce the development of the strong ion pressure gradient. The increase in ion pressure gradient can then sustain the sheared electric field, which then leads to the further development of a mean shear flow along the lines discussed by Burrell [60, 66]. This model has also been invoked to explain the formation of an extended region of reduced transport in the edge region [69, 70] known as the very high (VH) confinement mode as well as the formation of internal transport barriers (ITBs; see e.g. [71–73]). The experimental studies of the L–H transition and the evolution of turbulence in the shear flow region have recently been summarized by Wagner [74].

**1.2.4. Implications for experiment.** The implications of these theoretical expectations of the interactions between gradient-driven turbulence, mean sheared flows and turbulent-driven flows for controlled magnetized fusion are significant [75, 76]. Due purely to linear effects, the turbulent transport diffusivity can have a relatively small value if the free energy gradient is below a threshold value, and can then increase rapidly if the critical pressure gradient from linear theory is exceeded. Recent experiments have demonstrated this basic behavior (see figure 5) [77, 78]. The linear threshold for rapid transport (i.e. the ‘critical gradient’) depends, for example, upon collisionality, shape, magnetic shear and thus these quantities play a role in determining transport rates. The particular type of drift instability can vary depending upon the plasma conditions (e.g. the density, electron and ion temperatures, collisionality, magnetic shear, flow shear and plasma shape) and can lead to widely varying transport rates for particles, momentum, electron heat and ion heat. In discharges where heat is carried primarily by the electrons the electron pressure driven drift waves (e.g. either the universal, collisional TEM or collisionless TEM modes) are usually expected (but not always seen as discussed below) [22, 79]. The increase in confinement with density in ohmically heated tokamaks has been attributed to a reduction in this type of turbulence. If significant heat transport occurs through the ions due, e.g. to high plasma density and/or strong ion heating, then the ITG instability can be triggered. The onset of ITG turbulence has been associated with a degradation of confinement at high plasma densities in ohmically heated plasmas as well as in auxiliary heated discharges (the so-called L-mode regime) [80]. Finally, if the ETG is driven
to large values by strong electron heating the ETG instability may occur and has been predicted to lead to significant thermal transport [81]. However, the existence and significance of ETG turbulence are at this writing a contentious issue [82]. The threshold gradient needed for the onset of the large diffusivity can be increased above the linear value by the coupling of finite amplitude drift turbulence to the damped ZFs [36] (a phenomenon now referred to as the Dimits shift in the literature). The resulting kinetic core plasma profiles can be relatively insensitive to the details of heating profiles (i.e. they are ‘stiff’), making the core plasma conditions highly dependent upon the boundary conditions at the edge region.

Due to the turbulence–ZF coupling, the damping of the ZF can influence the turbulence saturation level and thus the transport [39]. Theory and simulation suggest that a competition between formation of longer range transport events such as avalanches and streamers and the shearing effect of a ZF or mean shear flow can occur [83–88]. This can result in a transition from the so-called gyro-Bohm to Bohm like transport where the characteristic transport step size changes from a local gradient scaling to a system size scaling. If a reduction in effective diffusivity then occurs due to the effects of sheared flows in the presence of a fixed cross-field flux, then the mean pressure gradient can increase in a localized region. An increased localized ion pressure gradient can then enhance the sheared radial electric field. Such sheared electric fields can then react back on the turbulence via the shear decorrelation mechanism and lead to a further decrease in transport, thereby reinforcing the reduction in transport and the associated transport barrier. The L–H transition has been attributed to such a phenomenon [60] as have ITBs [89].

Thus there is a rich set of potential dynamics of turbulence and transport in magnetized confined plasmas that depend on a number of key theoretical expectations for the behavior of drift turbulence. The approach taken in this review is to highlight published experimental studies that support, refute or challenge our understanding of these underlying key turbulence physics mechanisms, and then try to determine the extent to which these turbulence observations are consistent with the macroscopic transport behavior of the confined plasma.

1.3. Scope, key issues and physics questions of interest

Space constraints force us to limit the scope of this review. We have chosen to focus this paper on the major experimental findings that have emerged since several earlier reviews of
drift turbulence in confined plasmas [22, 90, 91]. Readers seeking a summary of such results for the period before 1990 are encouraged to consult these earlier works to learn of the earlier development of this subject. This review also excludes important studies of instabilities and turbulence in inhomogeneous plasmas that are of relevance to other applications such as space physics, solar and stellar physics and astrophysics. The focus of this review paper is the entire plasma region encountered in magnetic fusion experiments, and thus it is somewhat broader in scope than reviews that focused on the edge plasma turbulence [91, 92]. Other reviews have focused upon studies of particle, momentum and heat transport in magnetic fusion [80, 93, 94] using either power balance and/or transient transport analyses approaches; relatively little attention is given to these approaches in this paper although there are clear links to the subject of drift turbulence. The discussion of relevant theory is only cursory and the interested reader is referred to several recent reviews [14, 39, 59] that are available.

This paper addresses several key issues related to the observations of turbulence in confined magnetized plasmas and comparison with basic theoretical understanding. These are of course constrained by the available data. The questions include the following.

- How do coherent drift waves evolve into drift turbulence?
- What evidence exists that links turbulence with the actual plasma confinement in closed flux surface geometries?
- What are the dominant free energy sources and dissipation mechanisms that drive turbulence in the core, edge and SOL regions, and how do these vary with operating conditions? Is the usual picture of a linear instability saturated by nonlinear transfer to other linearly stable scales suitable and, if not, when does it break down?
- What is the evidence for the existence of ZFs and GAMs in confined plasmas?
- What role do ZFs and GAMs play in setting the turbulence amplitudes and associated transport rates?
- Do experiments shed any light on the mechanisms by which ZFs and GAMs might saturate?
- Are the turbulent transport scales determined by local gradients, or are they determined by the system size? Is there any indication in the turbulence data for a transition from one regime to the other?
- Does the turbulence play a role in triggering the formation of transport barriers at the edge and in the core regions of the plasma?
- Is the response of the turbulence to a shear flow consistent with theory?
- Is there evidence that the turbulence can develop a long-range character via triggering avalanche or streamer dynamics? Does turbulence spread or propagate from unstable regions into locations that are expected to be linearly stable?

The paper is organized as follows. First, we briefly summarize the key diagnostic techniques that are used in the study of turbulence in confined magnetized plasmas. Second, we examine a number of key results from laboratory plasma experiments performed (usually) with open field lines, since such studies usually allow more detailed measurements than in larger confinement devices. Third, we look at the evidence linking particle, momentum and heat confinement with turbulence and we also summarize the origins and behavior of turbulence in ohmically heated and auxiliary heated L-mode plasmas with weak flow shear. We then also summarize the evidence for the existence of ZFs and GAMs in confined plasmas by looking at the spatio-temporal structure of the flows and looking at their interaction with the turbulence. Next, we turn our attention to the changes in turbulence that occur during transport barrier formation, and we take a critical look at the role that turbulence might play in triggering such confinement transitions. We then examine what is known about the behavior of turbulence in well-established regions of strong flow shear, and we summarize existing experimental tests of
the theoretical predictions of the effect of a mean sheared flow upon turbulence and transport. We also briefly summarize the evidence that longer range transport events such as avalanches or streamers exist and play a significant role in the overall transport. We then return to the questions posed above, and try to point out what answers can reliably be drawn based on the available experimental results; we also point out areas where the data are unclear, contradictory or non-existent. The paper then closes with a set of suggestions for further work.

2. Fluctuation diagnostics

Space limitations do not permit a full discussion of the techniques used to measure turbulence quantities of interest and the data analysis approaches used to extract the physics information of interest. Here, we provide a brief summary of the techniques of interest and provide an introduction to the literature that discusses the diagnostic approaches and analysis techniques in more detail.

2.1. Langmuir probe-based techniques

Langmuir probe arrays have been used to measure density, potential and electron temperature fluctuations along with correlations between these quantities that then yield the fluxes of particles, momentum and heat (see e.g. [7–13, 19, 38, 52, 53, 55, 56, 61, 62, 90, 95–130]. The particle and convective heat fluxes have been measured by simultaneously measuring the density and electron temperature fluctuations and then cross-correlating these with the inferred $E \times B$ drift velocity found from two-point measurements of floating potential fluctuations. Similarly the momentum flux or Reynolds stress has been measured by cross-correlating fluctuations in the different components of the $E \times B$ velocity. Such measurements have been made in the edge and SOL regions of large tokamaks as well as across the entire plasma of smaller lower temperature experiments. Due to the heat flux from the plasma, probes are usually limited to plasmas with densities below $10^{19} \text{m}^{-3}$ and electron temperatures below 50 eV or so. Probe data are most commonly analyzed using Fourier transform based approaches which yield frequency and/or wavenumber spectra at different positions, times and conditions; more recently wavelet and other non-stationary schemes have also emerged to provide physics insight into non-stationary and intermittent conditions. Because of the recent theoretical interest in the nonlinear interaction of turbulence with larger scaled phenomena, nonlinear higher order techniques have been adopted and used to look, e.g., at the interaction of turbulence and shear flows, L–H transition physics and so on [42, 43, 131–137].

2.2. Collective Thomson scattering

Collective Thomson scattering occurs when an electromagnetic wave with $k \lambda_D \ll 1$ (here $k$ is the wave number of the incident wave and $\lambda_D$ is the Debye length) is scattered by the variations in the plasma refractive index caused by density fluctuations and early in the history of plasma physics research was proposed for use in studying plasma instabilities [138]. Such a wave incident upon a periodic plasma density perturbation of wavelength $\lambda$ will be scattered by an angle $\phi_B = 2 \sin^{-1}(\pi/k\lambda)$ due to the variation in the refractive index of the plasma, as long as the wavenumber and density fluctuation scale satisfy the Bragg scattering condition $k_S = k_i + k$ with $|k_S| = |k_i|$. Here $k_S$, $k_i$ and $k$ denote the scattered, incident and turbulent fluctuation wavenumbers, respectively. By measuring the scattered radiation at several angles, the wavenumber spectrum can be obtained [139]. The spatial resolution of such measurements is generally poor (i.e. the radiation is collected over a spatial scale, $L$, along the beam path
that can be significantly larger than the wavelength of the density fluctuations and for typical drift turbulence wavenumbers can approach the minor radius of the plasma. If the fluctuations exhibit a significant spatial variation, as found, for example, in the edge region of the plasma, the resulting measurement will provide only an average measurement of the turbulence over the scattering volume, and important spatial inhomogeneities may not be observed. We note that recent techniques to reduce these limitations have been proposed and used in tokamak experiments [140, 141]. A detailed discussion of collective scattering is given by Slusher and Surko [142].

2.3. Phase contrast imaging

Phase contrast imaging (PCI) launches a coherent wave of electromagnetic radiation that propagates through the plasma normal to both the magnetic field and the local mean plasma density gradient [143]. The transmitted wave is then composed of the incident wave (also known as the direct wave) and a wave that has a small phase disturbance (known as the diffracted wave) created by the passage of the radiation through the turbulent plasma density fluctuations. If the direct wave is phase shifted by \(90^\circ\) and then interfered with the diffracted wave, then the resulting intensity field will have a spatial variation that is proportional to the line-averaged phase shift suffered by the diffracted beam. The measured wavenumber range is determined by the width of the beam and the imaging optics. Within the context of this review paper, the technique has been used to study the response of edge plasma turbulence to the L–H transition in the DIII-D tokamak and core plasma fluctuations in ALCATOR C-MOD. The technique does suffer from the line averaging that occurs in the usual approach. Due to the strong fluctuation amplitude at the plasma edge, PCI results obtained through central plasma chords can potentially have significant imprints on them from this region. We note that approaches to reduce such effects have been discussed [144]. A detailed discussion of the technique can be found in the literature [143, 145].

2.4. Reflectometry

In reflectometry, a plane electromagnetic wave with wavevector normal to the magnetic surface propagates into the plasma until it reaches the cutoff layer, which is determined by the wave polarization, plasma density profile and local magnetic field. The wave undergoes Bragg scattering at the cutoff layer and is reflected back toward the launcher yielding, with suitable detection and data analysis schemes, the frequency and wavenumber spectra of the plasma density fluctuations. Good (~centimeter scale) radial localization is provided by the wave amplification that occurs at the cutoff layer due to the reduction in wave propagation speed near this region. The poloidal and toroidal localization of the measurement is determined by the details of the beam and detector optics. There are several variants to this basic technique that have recently been introduced. In the correlation reflectometry technique, two waves with slightly different frequencies are launched at the plasma, and interact and scatter at two radially displaced locations. If this separation distance lies within a turbulent correlation length, then information about the radial scale length of the turbulence can be obtained from the cross-correlation of the reflected radiation [146]. The technique of Doppler reflectometry has been successfully developed and deployed to measure the poloidal propagation velocity of density and has been introduced in [147, 148]. This technique can achieve high spatial and temporal resolution, allowing for measurements of the turbulence poloidal velocity, which can be compared with the background \(E \times B\) drift velocity; high time resolution measurements have successfully measured the GAM component of the ZF spectrum in ASDEX-UG [149]
and DIII-D. The interested reader is referred to the literature for detailed discussions of the application of reflectometry to the study of turbulence in confined plasmas [150, 151]. The upper hybrid resonance has also been proposed for correlation reflectometry to measure fine-scaled (electron gyroradius scaled) turbulence [152, 153].

Microwave imaging reflectometry (MIR) [154] expands reflectometry to 2D imaging of the wave cutoff layer. An MIR system images a reflected microwave beam onto multiple discrete detectors with an optical arrangement. Multiple lenses are used to tailor the reflecting surface to match the magnetic flux surfaces as closely as possible. Using multiple frequencies, it should be possible to expand the method to full 3D imaging of the density fluctuation structures. An implementation of the MIR diagnostic concept has been performed at the TEXTOR tokamak [154].

2.5. Beam emission spectroscopy

Beam emission spectroscopy (BES) measures localized, long-wavelength \( (k_\perp < 3 \text{ cm}^{-1}) \) density fluctuations by observing collisional fluorescence emitted by an injected neutral beam [155]. As neutral particles interact and collide with plasma electrons and ions, a fraction is excited and undergoes spontaneous radiative decay. Typically, the intensity of the Balmer H\( \alpha \) (or D\( \alpha \)) emission line \((n = 3 \rightarrow 2, \lambda_0=656.1 \text{ nm})\) is measured because of its relative brightness and the high efficiency and throughput of visible optics and detectors. The beam emission is Doppler shifted by the neutral beam velocity (several nanometers for \( \sim 80 \text{ keV} \) deuterium beams) allowing for spectral isolation of the beam emission from the intense edge recycling H\( \alpha \) (or D\( \alpha \)) emission. The normalized fluctuation intensity, \( \tilde{I}/I_0 \), of the emission could be related to the local normalized density fluctuation, \( \tilde{n}/n \), through a collisional radiative model [156]. For most implementations, a high-throughput optical system views the neutral beam with the line of sight nearly tangent to the local magnetic field lines so as to achieve optimal spatial resolution perpendicular to the field. Spatial resolutions of about 1 cm could typically be achieved and, with multiple channels deployed in the radial–poloidal plane, the 2D characteristics of the turbulence could be studied. BES systems have been deployed on both tokamaks [157, 158] and stellarators, allowing the determination of various fluctuation properties such as wavenumber and frequency spectra, amplitudes, radial and poloidal correlation lengths, decorrelation times, equilibrium and time-resolved poloidal turbulence velocity and 2D images of the turbulence [159]. The application of velocimetry techniques (which are discussed later in this review) [160, 161] to poloidally resolved measurements allowed identification and isolation of ZF features for both the low-frequency ZF [162] and GAM [163].

2.6. Gas puff imaging

Gas puff imaging (GPI) systems have been developed and deployed to image edge plasma fluctuations [164–166]. This technique allows for the visualization of light intensity fluctuations that are caused by the interaction of a turbulent plasma edge region with a quasi-two dimensional plume of neutral gas that is intentionally injected using a suitably designed gas delivery manifold. A camera capable of recording visible light images at a rate high enough to resolve edge turbulence fluctuations (e.g. typically framing rates of 100 kHz or higher) is then arranged with a view that is predominantly tangential to the magnetic field, and views the injected gas plume at roughly normal incidence as described in the literature [166, 167]. Working gases typically include the hydrogen isotopes and noble gases. Due to the thermal energies of the working gas and the ionization mean free path found in typical confinement discharges, the resulting light emission is peaked near the LCFS or separatrix region, providing
data over a region a few centimeters wide in the radial direction, while providing data over a region spanning tens of centimeters in the poloidal direction.

2.7. Electron cyclotron emission imaging

Electron cyclotron emission imaging (ECEI) uses a quasi-optical configuration to image ECE emissions onto an array of detectors. The radial locations are separated by examining different frequencies, as with conventional ECE, while poloidally distributed measurements are provided by the optical system. The imaging methods allow for radial and poloidal resolutions of about 1 cm, potentially allowing for long-wavelength ($k_{\perp} \rho_i < 1$) measurements of electron temperature fluctuations. The basic ECE method has been optimized to provide 2D imaging measurements in the radial–poloidal plane. An ECEI system has been deployed at the TEXTOR tokamak [168] to provide high time resolution images of electron temperature fluctuations associated with an $m = 1$ sawtooth oscillation over a 5 × 16 cm region. This system has been deployed in conjunction with the MIR system described in the previous section. By aligning the measurement volumes, the potential for spatially overlapping measurements of density and electron temperature fluctuations and the corresponding phase relationship between them may be accessible in future implementations, offering a very powerful technique for core 2D turbulence measurements.

2.8. Heavy ion beam probe

The HIBP has measured the electrostatic potential and fluctuations of density and potential simultaneously in the high temperature core of toroidal plasmas [20, 169–181] and is the only diagnostic capable of direct evaluation of the turbulent flux in this region. The HIBP measurement [169] detects the change in kinetic energy and current of a high-energy singly ionized heavy ion beam that becomes doubly ionized during passage through the plasma. In most applications a singly charged beam is injected into the plasma, and the doubly ionized beam born through the collisions with plasma electrons is detected. The former and the latter beams are usually termed primary and secondary beams, respectively. The energy difference between the primary and secondary beams corresponds to the electrostatic potential at the birth point of the detected beam. Therefore, the potential and its fluctuations could be evaluated in the measurement of the detected beam energy. The fluctuations in the detected secondary beam intensity are roughly proportional to the electron density fluctuations at the birth point of the secondary beam if the effect of the beam attenuation along the beam orbits could be neglected. As the plasma density is increased, however, the contribution of the beam attenuation, called path integral effects, becomes influential and disturbs the local density fluctuation measurement [182–185]. The position of the birth point of the secondary beam could be controlled by altering the beam energy and/or by sweeping the primary beam. The observable range of the HIBP measurement covers a 2D region of the plasma determined by the beam parameters and by device geometry and magnetic field. In addition, an HIBP system is usually equipped with multi-point detection of the secondary beams whose birth points are located in approximately ~centimeter range. The geometry and size of the sample volume is determined by the beam diameter usually of less than 1 cm and the entrance slit width of the secondary beam energy analyzer. Furthermore, multi-point detection of potentials at neighboring two points allows the direct measurement of electric field fluctuation by making difference between the two potential values, an essential capability for the study of ZFs and their interaction with turbulence. The simultaneous detection of both density and potential fluctuations also allows the evaluation of the fluctuation-driven local particle flux. The HIBP
could also detect magnetic field fluctuations in tokamaks [186] and more recently in non-axisymmetric devices [187]. With the correct measurement geometry, HIBP can be used to measure nonlinear energy transfer and turbulence saturation in the core plasma region (although this has not yet been done to our knowledge). HIBP measurements have been applied on a number of toroidal devices including ISX-B [20, 188], TEXT [181, 189], ATF [190], TEXT-U [179], JIPPT-IIU [191], T-10 [192], JFT-2M [173] and CHS [193].

2.9. Correlation ECE

In addition to ECEI diagnostics, discussed previously, correlation ECE (CECE) was deployed to measure turbulent electron temperature fluctuations in tokamak and stellarator plasmas [194, 195]. This diagnostic operates by deploying two closely spaced radial channels that achieve good poloidal and radial resolution, allowing for long-wavelength $T_e$ fluctuation measurements. The two channels are within a turbulence correlation length and are simultaneously acquired and a cross-correlation analysis applied. This technique eliminates uncorrelated noise in the measurements, allowing for isolation of the relevant electron temperature fluctuation signal. Measurements from this diagnostic are being quantitatively compared with gyro-kinetic simulations of turbulence to help validate the codes [196].

2.10. HF-CHERS

Ion temperature and toroidal velocity fluctuation measurements were obtained with the high-frequency charge exchange recombination spectroscopy (HF-CHERS) diagnostic system on TFTR [23, 197, 198]. This specialized diagnostic system exploits the charge exchange emission often used for measuring impurity (often carbon) ion temperature and toroidal velocity. The major differences with conventional charge exchange systems is the high light throughput needed to achieve vastly higher time resolution ($\sim 1 \mu s$) and good spatial resolution to isolate long-wavelength ion thermal fluctuations with an acceptable signal-to-noise ratio. The fluctuating ion temperature is found by applying a moments-based technique to the time-varying broadened emission line.

3. Data analysis techniques

These diagnostics produce one or more time-varying signals that are related usually to the fluctuating density, potential or electric field, temperature and, occasionally, magnetic field. The resulting set of time-series data are then analyzed using a wide variety of tools that have been developed to try and extract elements of the physics, and to facilitate comparison with theory.

3.1. Linear signal processing: correlation functions and spectra

The spectral analysis approach is one of the oldest techniques for turbulence study and is based on the Fourier transform. The method has been often used to indicate the wave-component structure of turbulence, and to investigate causal relationships between two variables by decomposing the turbulence into frequency and wavenumber components. The mutual relationship between two physical quantities of turbulence could be investigated using the cross-power spectrum $P_{crs}(\omega) = (1/N_{ens}) \sum_{i=1}^{N_{ens}} \tilde{f}_i(\omega)\tilde{g}_i^*(\omega)$ where $N_{ens}$ is the number of ensembles, $f_i(\omega)$ and $g_i(\omega)$ are the Fourier transforms of the time series $f(t)$ and $g(t)$ and subscript $i$ is the ensemble index. The spectra are defined as ensemble averages and therefore
the technique characterizes the stationary features of turbulence, but has essentially no ability to resolve dynamic and time-dependent characteristics of turbulence. The cross-power allows the coherence and cross-phase, given as

\[ \gamma_{12}(\omega) = \frac{|P_{crs}(\omega)|}{\sqrt{P_1 P_2}}, \quad \text{and} \quad \tan \phi(\omega) = \frac{\text{Im}(P_{crs}(\omega))}{\text{Re}(P_{crs}(\omega))}, \]

respectively, to be computed from the spectra. Here \( P_1 \) and \( P_2 \) denote the autopower spectra of the time series \( f(t) \) and \( g(t) \). The coherence indicates the strength of the linear relation between two physical quantities and is bounded between \([0,1]\).

The correlation function \( C_{crs}(\tau) = (1/N_{ens}) \sum_{n=1}^{N_{ens}} (1/T) \int_{-T/2}^{T/2} \tilde{f}(t) \tilde{g}(t+\tau) \) is related to the spectrum via the Wiener–Khinchine theorem, \( P_{crs}(\omega) = \int_{-\infty}^{\infty} C_{crs}(\tau) \exp(i\omega\tau) \, d\tau \), where \( \tau \) is termed the time lag. Usually, the correlation functions are normalized by \( |\tilde{f}(t)||\tilde{g}(t)| \) and are then bounded by \([-1,1]\). The normalized correlation function computed from turbulent fluctuations is usually a decreasing function of \( \tau \), and a characteristic decorrelation time can be defined (usually from the delay needed to give a 1/e-folding decrease in \( C \)). By replacing time by position and frequency by wavenumber, these quantities could be readily extended to the wavenumber domain. Using these relations, the fluxes discussed earlier could be Fourier decomposed, e.g. as

\[ \Gamma_r = \frac{1}{B} \sum_{\omega} k_{\theta}(\omega) \sqrt{P_{\tilde{n}n}(\omega) P_{\tilde{n}\phi}(\omega)} y_{n\phi}(\omega) \sin(\phi_{n\phi}(\omega)) \Delta \omega, \]

where \( P_{crs}(\tilde{n}, \tilde{\phi})(\omega) \) is the cross-power of density and potential fluctuations, \( \Delta \omega \) is the resolution of the Fourier transform and \( k_{\theta}(\omega) \) is the wavenumber which could be obtained from, for example, the two-point technique [199].

### 3.2. Nonlinear signal processing: energy transfer and multi-scale interactions

Turbulence results from the nonlinear interaction between Fourier components, which leads to a nonlinear energy transfer that moves the energy to other spatio-temporal ranges. Bispectral analysis is one of the techniques to investigate the wave–wave couplings, or the interactions between the fundamental waves. The bicoherence based on the Fourier spectral analysis [2–4] is the normalized bispectrum and is defined as

\[ b^2 = \frac{\left| \sum_{n=1}^{N_{ens}} \tilde{f}(\omega_2) \tilde{f}(\omega_1) \tilde{f}(\omega_3) \right|^2}{\left( \sum_{n=1}^{N_{ens}} |\tilde{f}(\omega_2)|^2 \sum_{n=1}^{N_{ens}} |\tilde{f}(\omega_1)\tilde{f}(\omega_3)|^2 \right)^{1/2}}, \]

where \( \tilde{f}(\omega) \) is the Fourier components, and the condition \( \omega_1 = \omega_2 + \omega_3 \) should be satisfied for non-zero bicoherence value. In case that \( \omega_1 = \omega_2 + \omega_3 \) is satisfied, the numerator takes a form of \( \sum |A_i| \exp(i\delta_i) \), where \( \delta = \delta_2 - \delta_1 - \delta_2 \) denotes the phase difference between the phases of the three waves. The ensemble-averaged numerator vanishes if the individual waves have a random phase distribution \( \delta_i \) across the ensembles. On the other hand, if the wave phases are coherent (i.e. the phase takes on a repeatable value) across many ensembles, the numerator gives a finite value. Therefore, the bicoherence is an indicator of the coupling strength between three elemental waves to satisfy the matching condition.

This approach was first used to show the presence of three-wave couplings in broad-band spectra measured with Langmuir probes in the plasma edge region in the ATF torsatron [137]. More recently, this approach has been applied in order to elucidate the processes of turbulence generated shear flows [43]. Bicoherence analysis could reveal hidden phase couplings between fundamental wave components and distinguish between nonlinear coupled phase coherent components and independent phase incoherent components.
The nonlinear dynamics of weakly dissipative drift turbulence is thought to be described by the Hasegawa–Mima equation [200, 201]

$$\frac{\partial \phi(k, t)}{\partial t} = (\gamma_k + i \omega_k) \phi(k, t) + \frac{1}{2} \sum_{k_1, k_2} \Lambda^Q_k(k_1, k_2) \phi(k_1, t) \phi(k_2, t),$$

where $\phi(k, t), \gamma_k, \omega_k$ and $\Lambda^Q_k(k_1, k_2)$ represent the Fourier coefficient with the wavenumber $k$ of the fluctuating potential field, the growth rate, the dispersion relation and the wave–wave coupling coefficient, respectively. This model has been extended to cases where the dissipation is stronger, resulting in coupled nonlinear equations for the advection of vorticity and density [200, 201] in an isothermal plasma. The equation can be used to describe the evolution of the power spectrum $P_k = \langle \phi^2_k \rangle$,

$$\frac{\partial P_k}{\partial t} \simeq 2 \gamma_k P_k + \sum_{k_1, k_2} T_k(k_1, k_2),$$

in terms of the linear growth rate and nonlinear energy transfer $T_k(k_1, k_2) = \Lambda_{12}(k_1, k_2) (\phi^*_k \phi_{k_1} \phi_{k_2})$ which represents the power transfer function from $k_1$ and $k_2$ to the wave $k$ and is found from the bispectrum $(\phi^*_k \phi_{k_1} \phi_{k_2})$ computed from the measured fluctuations.

In the actual analysis, the coupling coefficients are obtained by a finite difference equation between $X_k = \phi(k, t)$ and $Y_k = \phi(k, t + \tau)$, which is derived from the above equation (see details in [133–135]). These techniques must be modified when the turbulence is described by more than one field; however, no formal published application of such a multi-field turbulence nonlinear dynamical model currently exists in the literature.

The amplitude correlation function (ACF) method is a combined method of the numerical filtering technique and correlation function. In this method, the envelope or amplitude evolutions of the filtered fluctuation signals in two frequency regimes are evaluated first, then the correlation between them is evaluated, which is expressed as

$$C_{ACF}(\tau) = \langle A(t; \omega_1, \omega_2, \tilde{f}) A(t + \tau; \omega_3, \omega_4, \tilde{f}) \rangle,$$

where $A(t; \omega_i, \omega_j, \tilde{f})$ represents the amplitude of the band-pass filtered signal, $\tilde{f}$, from $\omega_i$ to $\omega_j$ as a function of time. Since the amplitude should be the measure for the energy contained in the frequency band, therefore, the existence of a significant correlation indicates the presence of energy transfer between them. In addition, the time lag or delay giving a maximum of the correlation should reflect the direction of the energy transfer.

### 3.3. Wavelet based approaches

Fourier analysis cannot resolve time-dependent feature of turbulence, although turbulence is characterized by intermittent or short-lived events. This issue could be avoided by substituting the Gabor or van Milligen wavelet transform, defined as

$$f_G(a, t) = \int \Phi(a, t - t') f(t') \, dt$$

for the Fourier transform in the above expressions for spectra and energy transfer, where $\Phi(a, t) = 1/\sqrt{a} \exp[i2\pi t/a - (t/a)^2/2]$ [202, 203]. Here, if the parameter $a$ is chosen as $f = 1/a$, then the parameter $f$ corresponds to the frequency in the traditional Fourier analysis. The width of the temporal window $\Delta t = a$ and depends on the frequency.
other hand, the Fourier transform of the wavelet base shows that the frequency resolution is 
\( \omega = 2\pi/a \pm \pi/a \).

The van Milligen wavelet [202–204] is a suitable tool to resolve the time-dependent or intermittent phenomena widely observed in turbulence plasmas, but it has not received wide application. In particular, the wavelet biocoherence analyses have been carried out for the Langmuir probe data from the ATF torsatron and W7-AS stellarator, when the method was proposed [202, 203], and for the comparison between L- and H-mode in the DIII-D tokamak [205]. In addition it has recently been applied to the cylindrical plasma device, VINETA, to investigate the two-dimensional intermittent couplings in the domain of wavenumber and frequency [206] and to show the existence of nonlinear couplings between turbulence and stationary ZFs in CHS [207].

3.4. Conditional sampling and conditional averaging

In turbulent plasmas, many kinds of outstanding intermittent or transient phenomena have been observed to occur, such as transition to improved confinement, sawtooth crashes, blob generation and so on. Although these phenomena should have a statistical nature, the most probable evolution of such phenomena could be extracted from measurements by making an average over many repeated examples of the phenomena of interest. Usually one defines e.g. a threshold that must be exceeded for an event to be considered to have been observed. All such independent events are then extracted from data, and averaged together to find an average spatio-temporal view of the phenomena of interest. The usefulness of the method of the conditional average was demonstrated and discussed in a conventional double-plasma device [208] in an experiment that examined inverse-Landau damping of ion acoustic waves (IAWs), and has since been applied in a number of turbulence experiments.

3.5. Image based approaches

The recent development of imaging based diagnostics, using either some type of plasma radiation emission measurement or using large numbers of spatially distributed point measurements using, e.g., Langmuir probe arrays, has led to the introduction of several new image processing based techniques.

3.5.1. Birth/death statistics, structure motion

Using a large array of probes in the TORPEX simple torus device, Muller et al introduced the use of schemes borrowed from image processing to automatically determine the birth, evolution and death of turbulent 'structures' which were defined usually in terms of an event or disturbance with an amplitude that exceeded a critical value specified by the experimentalist [209–211]. Such events could be identified at their birth, and low order statistics such as the structure centroid location, major and minor elliptical axes and orientation and so forth could be automatically tracked during the subsequent evolution. Eventually the structure would no longer exceed the specified criteria, and would then be considered to have died. Using automated algorithms, these measures would be identified for a large population of such structures, and then this information used to determine some element of physical understanding of the origins and dynamics of such structures.

3.5.2. Velocimetry

Velocimetry techniques use imaging diagnostics to track the motion of scalar quantities such as density or temperature fluctuations, and then infer the flow field form of the motion of the scalar field. The conceptual underpinning of velocimetry techniques is quite simple. If one assumes that the dynamics of the scalar field are dominated by the
advection of that field, then by comparing the time history of the field at two spatial locations one should be able to watch the field structure ‘sweep past,’ with the time history of the field at the second point being equal to the first, but offset by a time lag dependent on the spatial separation of the channels and the magnitude of the velocity. More rigorously, if the dynamics of the conserved scalar field $n(x, t)$ are dominated by advection of a steady, spatially uniform velocity $U$, one can then write

$$\frac{\partial n}{\partial t} + U \frac{\partial n}{\partial x} \approx 0 \Rightarrow n(x, t) = n(x - U t).$$

If then one computes the cross-covariance between $n(x, t)$ and $n(x + \Delta x, t)$, one finds

$$R_{\Delta x}(\tau) = \int n(x, t)n(x + \Delta x, t + \tau) \, dt \propto \int P_n(k) \cos(k(\Delta x - U \tau)) \, dk$$

(where $P_n(k)$ is the spatial power spectrum of the scalar fluctuations). It is easy to show for experimentally relevant spectra that $R_{\Delta x}(\tau)$ has a ‘global’ (in $\tau$) maximum at $\tau U = \Delta x/U$. Therefore, by computing $R_{\Delta x}(\tau)$ and finding the lag $\tau_U$ where it peaks, one could infer the velocity $U$. Alternatives to finding the location in the peak of the cross-covariance function include examining the cross-phase between the two scalar field signals, possibly using wavelets [160, 212, 213]. Other algorithms have been proposed to infer the flow field [214] as well by identifying structures in the scalar field and looking at the displacement of those structures in two adjacent time frames.

There have been only a few studies in which the inferred flow field is compared with other measurements to determine whether the inferred flow field corresponds to the turbulent $E \times B$ velocity. Yu et al measured the azimuthal flow velocity using the TDE algorithm and compared it with the inferred plasma fluid flow and concluded that the TDE algorithm measured the Doppler shifted diamagnetic velocity of drift fluctuations [215]. Munsat and Zweben used a velocimetry algorithm that combined aspects of pattern matching techniques with optical flow techniques constrained by a conservation law imposed on the scalar intensity field and assuming an incompressible velocity field [216]. Their results allowed them to infer two-dimensional turbulent flow fields from GPI data in NSTX. However, there were no turbulent measurements with which to compare and thus it is not clear whether the inferred flow field is a good representation of the actual turbulent flow. At this writing, it seems that this approach could infer large-scale coherent flows with reasonable fidelity [217] but the quality of the inferred two-dimensional stochastic flow fields associated with turbulence remains to be determined in future work. Clearly if such flow fields could be inferred accurately, significantly new measurements in the core regions of fusion experiments would be made possible, and detailed turbulence physics studies could be enabled by inferring e.g. the vorticity and enstrophy from the velocity field data.

3.5.3. Searching for long-range correlations using structure functions and R/S analysis

Long-range correlations in space and time could be associated with, for example, long-lived spatially localized coherent structures, or with the generation of avalanche and streamer events which could propagate over a distance that is significantly larger than the typical turbulence correlation length. The structure function analysis technique has been proposed to search for such correlations in existing turbulence data [218, 219]. In this approach a time series, $Y(t)$, obtained from a suitable turbulence diagnostic is used to compute the function

$$S_q(\tau) = \langle |Y(t + \tau) - Y(t)|^q \rangle$$

where $\tau$ denotes a time delay. If the physics that determines $Y(t)$ is scale-invariant and self-similar on some range of time delays bounded by $\tau_1 < \tau < \tau_2$, then the $q$th structure function will vary as

$$S_q(\tau) = C_q \tau^{\zeta(q)}$$

where $C_q$ is a slowly varying function of $\tau$, and the exponent $\zeta(q)$ determines how quickly the structure function varies with
Values $0.5 < H < 1.0$ indicate that the physical processes that determine $Y(t)$ exhibit long-range positive correlations. A value of $H = 0.5$ indicates no such long-range correlation, and values $0 < H < 0.5$ indicate long-range anti-correlations. The calculation of $H(q)$ involves a number of important technical issues; we refer the interested reader to the literature (see [218] and references therein). We also note that the so-called rescaled range (R/S) technique has been applied to permit the rapid computation of the Hurst parameter [220], and that the validity of the R/S approach has been questioned recently [221, 222]; however, a recent in-depth comparison of this approach with the more computationally intensive structure function approach summarized here concluded that the two techniques did indeed give the same result for $H$ provided that long enough time series were used, and confirmed that the interpretation of $H$ summarized above was indeed correct [219].

4. Laboratory plasma device results

Although the parameters achieved in laboratory plasma devices are often quite different from those in fusion confinement experiments, the universal nature of the drift instability, coupled with a number of common elements of nonlinear dynamics and turbulence, make it possible to gain significant insight into the basic physics of drift turbulence and transport in such experiments.

Here we highlight several key results from studies of drift waves and drift turbulence in laboratory devices which verify the linear physics of various drift instability branches, show how coherent drift waves transition into drift turbulence, how this turbulence could lead to the development of organized structures and processes within the plasma and which show the existence of turbulent-driven azimuthally symmetric sheared ZFs in simple plasma configurations.

4.1. Coherent wave studies

Early work in linear Q-machine (for quiescent plasma) devices showed the presence of collisional and collisionless coherent drift waves when the wave became linearly unstable due to a change in a controlling plasma parameter (usually either the magnetic field, which decreased ion collisional viscosity, or parallel current, which created a new free energy source) was increased. When this free parameter exceeded a critical value, a spontaneously generated electrostatic wave would be observed in the plasma; the dispersion relation of this wave would then be compared with linear stability theory of the appropriate class of drift wave. These early results usually had the electron pressure gradient as the primary source of free energy driving the drift wave, and provided confirmation of the essential elements of the linear drift wave instability, and showed the important role that parallel electron dissipation played in causing linear instability. The interested reader is referred to these original papers [223–228] or to textbooks for a discussion of the various classes of electron pressure gradient driven linear drift instability. We also note that similar studies of the onset of coherent collisional drift waves have been provided in a toroidal heliac device [229], and gave good agreement with a finite beta collisional drift wave theory developed for a high aspect ratio circular cross-section torus [230].

In fusion confinement devices, it is thought that ITG-driven drift instabilities may be an important driver for transport in the core region of many experiments [28], and thus there is interest in studying the basic physics elements of this class of drift wave instability. Work in the Columbia Linear Machine (CLM) using a plasma source capable of producing a radial gradient
in the parallel ion temperature showed the clear development of an ITG-driven coherent \( m = 2 \) wave in the plasma when the critical value \( \eta_{i\perp} = d \ln T_{i\perp}/d \ln n > \eta_{i\parallel} = d \ln T_{i\parallel}/d \ln n \) was exceeded [231]. The dispersion relation of the measured oscillation agreed with the linear theory that was developed for this experiment with anisotropic ion temperature, confirming the nature of the observed oscillation. Subsequent work in the same experiment demonstrated heat transport caused by this ITG instability [232], looked at the effect of flow shear stabilization on the ITG mode [233], and looked at the transition from the slab branch to the toroidal branch of the ITG mode [234]. We also note that work in the LAPD device showed the merger of collisional drift waves and Alfvén waves at finite beta [235], in agreement with expectations from linear theory.

4.2. Parametric decay of drift waves

Following the early exploration of the linear drift instability, workers moved to the study of the beginning of the nonlinear evolution of coherent drift waves—the first step on the road to turbulence. The key signature of the development of a parametric decay process involving three waves was shown in two experiments [236, 237]. Results from one are shown in figure 6. The plasma was operated at a condition where a single coherent drift wave was present at finite amplitude, and then an exciter electrode was energized to drive a second, pump mode. The phase of this pump mode could be varied relative to the phase of the drift wave, and the pump mode excitation voltage could also be varied. When the pump mode phase was such that it would nonlinearly beat with the drift wave in a constructive manner, then a parametrically excited half-harmonic decay wave would be observed if the pump mode exciter amplitude exceeded a critical value. If the pump mode phase was incorrect, then no decay wave was observed. The results were also compared successfully with an analytic three-wave parametric decay model, confirming the essential physics of this nonlinear process.

The nonlinear interaction of spontaneously excited coherent drift waves was also studied using a bicoherence analysis to identify phase coherence between triplets of waves [238]. The results showed that in the presence of finite amplitude coherent drift waves, nonlinearly driven oscillations satisfying the three-wave interaction frequency summation rule began to appear and were phase coherent with the parent waves. This latter observation was taken to indicate that these waves were in fact nonlinearly driven perturbations and were not phase independent, linearly unstable eigenmodes of the system.

4.3. Transition to drift turbulence

Detailed studies of the transition from coherent drift wave activity to broad-band drift turbulence have been carried out in several linear plasma devices [239–240]. In the first experiment coherent drift waves were excited by a parallel current driven between a biased mesh located between the plasma source and the downstream solenoidal magnetic field. Langmuir probes were then used to measure the fluctuating density. As the parallel current was increased, first a single coherent drift wave was observed (figure 7) and then, as the parallel current (which represents the free energy source in this experiment) was raised further, multiple frequencies were observed. Eventually these narrow band frequencies merged into broadened frequency peaks.

Further insight to the physics of the transition was obtained by reconstructing a projection of the phase space explored by the drift waves (figure 7), and calculating the embedding dimension, \( D_2 \), and the Lyapunov exponents, \( \lambda_i \), of the fluctuating system. These approaches,
which find their origin in the study of nonlinear complex systems, complement the more traditional Fourier-transform based approaches that have traditionally been used.

The transition to turbulence was found to be characterized by two Hopf bifurcations, the first of which result in a transition from a FP to a limit cycle (LC) attractor in a three-dimensional projection of the phase space. The second Hopf bifurcation results in a two-torus (T2) attractor; further increases in parallel current lead first to a mode-locked (ML) state which eventually transitions to a state of chaos. With a strong enough parallel current this chaotic phase space portrait was found to break up and fill a bounded phase space homogeneously, indicating that a
Figure 7. Evolution of the reconstructed phase space portrait during transition to turbulence; reprinted with permission from [240]. Copyright 1997 Institute of Physics.

state of turbulence had been reached. This transition was also described by the corresponding evolution of the embedding dimension, \( D_2 \), which describes the minimum number of degrees of freedom of the oscillations, and by the Lyapunov exponents of the dynamics, which describe how quickly two neighboring points in phase space diverge. The parameter \( D_2 \) was found to increase to unity in the LC regime, to 2 in the T2 regime, and then collapsed back to 1 in the ML regime. \( D_2 \) then diverges as the state of chaos and turbulence is reached, indicating that many degrees of freedom exist in that state. The Lyapunov exponents show the existence of exponentially diverging phase space trajectories in the state of chaos and turbulence. These results were consistent with the Ruelle–Takens route to turbulence that has been reported in many fluid mechanics experiments.

The development of broad drift wave spectra was also studied using Fourier transform based approach in the works of Ricardi et al [241] and Burin et al [240] which is naturally suited to viewing the problem from the perspective of wave–wave interactions. In both experiments, the transition from coherent drift waves to drift turbulence was studied as the magnetic field was increased; the first experiment was carried out in a simple magnetized torus, while the second was carried out in a cylindrical magnetized plasma. Here we show selected results from the linear device [240] to illustrate the key findings of these experiments and to facilitate comparison with those of Klinger et al.

The gradual development of the spectra (figure 8) showed that linearly unstable drift wave eigenmodes were initially excited; these modes then began to interact via three-wave coupling process as determined by bicoherence analysis. Above a critical magnetic field of \(~700 \text{ G}\) a ML spectrum developed; further increases in the magnetic field then led to the loss of the ML state and the emergence of much broader spectra. The development of the wavenumber resolved energy spectrum showed that, as a state of turbulence is approached, the magnitude of the linearly stable, low wavenumber portion of the energy spectrum increased by over a factor of 100, while the intermediate range of wavenumber, which corresponded to linearly unstable
The nonlinear dynamics of energy transfer in drift wave turbulence were also studied in the VINETA device [206]. In these experiments, an azimuthal array of probes was used to measure the time variation of the density fluctuation azimuthal mode number spectrum. Higher modes, only increased in amplitude by factors of 2–3. This result was taken as an indirect measure of the existence of a nonlinear transfer of energy to the low azimuthal wavenumber, large scale length portion of the spectrum.

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wavenumber fluctuations were observed to precede the development of intermediate mode number fluctuations; at the same time, the degree of phase coherence, as judged by the wavelet basis bicoherence, was seen to increase. The results were interpreted as being consistent with the transfer of turbulence energy from smaller scaled features into larger scaled features.

4.4. Observations of fluctuation-driven ZFs in laboratory plasma devices

One essential element of the nonlinear dynamics of drift turbulence is the emergence of sheared ZFs, which are linearly stable $E \times B$ drifts within the magnetic surface driven by the nonlinear transfer of energy from the drift waves into the ZF. Recent work in the CLM and CSDX laboratory devices [129, 242] has shown that such flows exist in confined plasmas, and that they are nonlinearly driven by the turbulence.

In the CSDX experiments, probe measurements confirmed the existence of a radially sheared azimuthal flow. The $m = 0$ nature of this flow has been confirmed with fast imaging techniques. This sheared flow was then shown to be consistent with a turbulent-driven ZF using probe techniques [97]. The essential results are summarized in figure 9. A multi-tipped Langmuir probe array was used to measure the turbulent Reynolds stress in the experiment, and then this momentum flux was used in the turbulent momentum conservation equation along with an estimated collisional viscosity profile inferred from line-averaged ion temperature measurements. The results showed that the observed sheared flow was self-consistent with the turbulent Reynolds stress. A two-field numerical simulation of the drift turbulence was also performed, and showed the emergence of the shear flow during the development of the drift turbulence; the simulation result was also consistent with the experimental observations.

These results were dependent on the estimated ion–neutral and ion–ion collisional dissipation terms. Recent work on a similar device in which the ion temperature and neutral temperature profiles were measured [243] found that the ion viscosity is more peaked on axis, and that the neutral gas temperature was significantly lower than reported in the work on CSDX, resulting in a lower degree of neutral gas depletion and thereby increasing the ion–neutral flow
drag in the experiment. These effects would have important consequences for the interpretation of the ZF results, and it would be valuable to have simultaneous measurements of the turbulent Reynolds stress, shear flow and flow dissipation profiles; at present no such measurements exist.

In the CLM experiments [242, 244], the existence of a low-frequency ZF was inferred using a clever azimuthally symmetric ‘ring’ shaped probe along with conventional point-like probes. The ZF modulates the lab-frame frequency of the ITG mode, and the resulting frequency modulation was inferred using short time Fourier transform techniques. In addition, the azimuthal and axial wavenumbers of the ZF were also measured, and confirmed that the ZF was azimuthally symmetric and was uniform along the magnetic field. The amplitude of the ZF was also deconvolved from the spectral data, and showed a significant flow shear rate at the periphery of the plasma column.

The linearly stable ZF appears to act as an energy sink to the drift turbulence in these experiments. Interestingly, work in the CLM device also points to the possibility that drift waves could couple nonlinearly to ion acoustic waves (IAWs) [245, 246]. Since the damping rate of the IAWs is dominated by ion Landau damping and ion temperature is low in the CLM device, an increase (decrease) in ion mass could then decrease (increase) the damping rate of the IAW. Since these waves are then acting as an energy sink, and the amplitude of the IAW modulates the rate of energy transfer from the drift turbulence, the saturated drift turbulence amplitude will depend upon the ion mass [247] in a manner analogous to the predator–prey model of drift turbulence–ZF interactions [248] where the IAW replaces the ZF. Sokolow and co-workers have argued that this may help explain the origin of the isotope scaling of confinement in toroidal devices. It is not known whether this mechanism is indeed operative in confinement devices.

4.5. Generation of intermittent transport events

The discovery of the generation of isolated plasma density structures in the open field line region of tokamak devices using GPI, which subsequently propagate toward the wall, led to several studies in linear devices [249–252]; these experiments were motivated in part as a test of theoretical predictions of blob dynamics that were specific to toroidal devices [253–256]. Results from the LAPD linear device are presented here to represent the essential findings of these studies.

Figure 10 shows the raw ion saturation current fluctuations (proportional to the plasma density) at three radial locations in the LAPD device. The left-hand plot shows data obtained 1.3 cm inside the edge of the plasma (defined by the cathode radius), the central plot shows results 0.7 cm outside the edge and the right-hand side shows results obtained 2.2 cm outside the edge. The density fluctuations show a significant negative (positive) going tendency at the location inside (outside) the cathode radius. The PDF of these data in panels (b) and (c) of the figure confirms these trends. Conditionally averaged measurements show monopole-like density blobs propagating outward, and are associated with potential perturbations that give rise to a convective cross-field $E \times B$ drift of the density structures as seen in figure 11. Density holes are also observed to propagate inward with a more complex spatial structure.

Similar dynamics have been reported to occur for larger scaled coherent structures which occur in simple magnetized toroidal devices which have a purely toroidal magnetic field (see [241, 257, 258]). Results from the TORPEX device, which could vary the magnetic geometry from a simple torus with a purely toroidal magnetic field to one with both toroidal and vertical field, provide additional insight into the relationship between drift turbulence, blob generation and background $E \times B$ shear flow. In a series of papers from this device
interchange instabilities are observed to transition into a state of turbulence driven by a combination of drift and interchange instability when a sufficiently strong vertical magnetic field is added to the plasma, and blobs were generated when the pressure gradient exceeds a critical value. The resulting radially elongated structures then protrude out into a background $E \times B$ shear flow, which breaks the structures off and forms isolated coherent density perturbations which are then able to evolve and move independently from their point of origin.

We also note that evidence for the existence of an azimuthally localized radially elongated streamer in drift turbulence has been reported in a linear cylindrical plasma device [263]. Such structures, which could degrade confinement by providing a rapid means of plasma transport across the magnetic field, were identified in this work as being generated by the nonlinear phase locking, or bunching together, of drift waves to form the elongated structure. These observations await confirmation from other experiments.

5. Confinement experiment results

We first summarize the evidence linking turbulence and transport in confinement experiments. We then turn our attention to attempts to identify the ‘origin’ (i.e. the linear instability) of the turbulence in the core and edge/SOL plasma regions. From there, we then show the evidence for the existence of ZFs, and examine how they affect the turbulence. We also examine
Figure 11. Temporal behavior of the conditionally averaged density distribution of (a) positive density perturbations (‘blobs’) and (b) negative going density perturbations (holes), (c) spatial distribution of conditionally averaged positive density excursion (blobs) which clearly show a monopolar structure. (e) The potential fluctuations associated with this density excursion show a clear dipolar structure, which leads to an outward convection of the density perturbation as shown in (a). (d) The negative going density structures (holes) show a more complex structure and have a potential distribution (f) that leads to inward propagation of these hole structures. Thus in this linear experiment, outward going density blobs and inward going density holes are generated at the edge of the plasma, and lead to significant transport in that region of the device; reproduced with permission from [252]. Copyright 2006 American Institute of Physics.

5.1. Evidence linking turbulence and transport in toroidally confined plasmas

A number of studies have searched for evidence linking turbulence to confinement. These studies usually measure the parametric variation of turbulence with some system parameter (e.g. density, heating power, plasma current, etc) and then compare this variation with the variation of inferred particle and/or thermal diffusivity. Usually it is found that an increase in diffusivity is linked to changes in the turbulence that are consistent with an increase in turbulent transport.
Occasionally the turbulence measurements permit the particle and/or (electron) heat flux to be inferred. In such cases, the flux is integrated over the flux surface and the resulting total particle or heat loss is compared against expectations from global confinement measurements. We summarize results from both approaches here. The reader is also encouraged to consult the review by Wootton et al for a more extensive discussion of earlier studies linking turbulence and global confinement [22].

5.1.1. Parametric studies of particle and heat fluxes. Many experiments have documented parametric variations of particle and/or heat confinement and turbulence (see [22] and references therein for summaries of earlier results). These earlier results provide somewhat contradictory results. For example, early work on the TFR tokamak showed evidence linking energy confinement time changes and changes in turbulence [264]. Measurements in ISX-B [20, 265] (figure 12) showed that neutral beam heating gave a large increase in the turbulent cross-field particle flux consistent with changes in global confinement. In beam heated ISX-B discharges, the convective thermal heat flux carried by the particle flux, given as $T_e \langle \hat{n}_e \hat{v}_e \rangle$, was roughly (within a factor of ~2) consistent with global energy confinement time. In ohmic discharges this flux was too small to account for the thermal confinement, suggesting that perhaps some other process was operative in that regime.

Detailed studies of transport and inferred turbulent heat flux in TEXT were less conclusive, with the inferred turbulence flux being substantially smaller than that expected from transport analysis [79]. However, the possible presence of counter-propagating drift modes in the TEXT data could have caused a spurious reduction in the measured turbulent wavenumber.
which would then erroneously reduce the inferred turbulent flux. In other work on TEXT, measurements using a HIBP system [20, 265] were combined with probe measurements in ohmically heated TEXT discharges. Results showed that the turbulent particle flux in that device was large enough to be significant and was close to that needed to explain global particle balance [22]. In work on ohmically heated ASDEX discharges, the particle flux was measured [25] and compared with estimates of the global particle confinement time made in the SOL region [266]. The results showed that the inferred particle flux integrated over the LCFS would yield a particle confinement time that was within the systematic and random errors of the turbulence measurements and of the inferred particle confinement time. Results in ohmically heated Tore Supra discharges examined the variation of the electron thermal diffusivity, inferred at a position 2/3 of the distance from the plasma center to the plasma edge, as the plasma density is increased, and compared this variation with that of the relative density fluctuation amplitude. The electron heat diffusivity and density fluctuations were both found to decrease with increasing line-averaged density (figure 13), qualitatively consistent an increase in energy confinement with increasing density [80].

We also note that the electron heat flux due to magnetic fluctuations, $Q_{EM} = \langle \tilde{q}_e \tilde{B}_r \rangle / \tilde{B}$, has been measured [267, 268] using a novel diagnostic in both an RFP and a small ohmically
heated tokamak, CCT. In the RFP, the magnetic fluctuations were attributed to MHD tearing modes and thus fall outside the scope of this review. The magnetic heat flux was, however, large enough to impact the energy confinement in the device. In the CCT tokamak result, the magnetic fluctuation-driven heat flux was not large enough to account for the electron heat flux in the ohmic discharge except near the X-points of magnetic islands located at low order $q = m/n = (3/2, 2, 3)$ resonant surfaces (here $q$ is the safety factor and $m, n$ are the poloidal and toroidal Fourier mode numbers. It was speculated that MHD activity localized to the rational surface caused localized stochastic field lines in the tokamak. Again, these observations would seem to suggest that mechanisms other than drift instabilities are responsible for these results. To our knowledge these are the only such measurements correlating heat flux fluctuations with magnetic field fluctuations in a tokamak device under any operating conditions.

Auxiliary heated discharges gave rise to the so-called L-mode confinement regime in which the energy confinement time is reduced from what was experienced during ohmic heating (see e.g. [269] for an early report of this phenomena; a more recent summary of L-mode transport could be found in [28]). The results from ISX-B, summarized in the preceding paragraphs, illustrate this phenomenon via turbulence measurements. More recently, spatially localized, multi-point measurements of density fluctuations were obtained in TFTR L-mode discharges using BES [270] and were used to study the role of turbulence in strongly heated L-mode discharges. These multi-point measurements allow the effects of any potential counter-propagating modes, which made interpretation of the TEXT results ambiguous [79], to be accounted for in the estimation of turbulent transport. The TFTR measurements demonstrated a radial correlation length of about 2 cm, while the poloidal correlation exhibited an oscillatory decaying behavior, showing a wave-like character in the poloidal direction. The corresponding wavenumber spectra demonstrated a peak in $S(k_\theta)$ near $k_\theta = 1$ cm$^{-1}$, $(k_\perp \rho_s \approx 0.1$–0.3) while the $S(k_r)$ spectrum peaked at $k_r = 0$. The radial profile of the fluctuation amplitude increased gradually from $\tilde{n}/n \sim 1\%$ in the core region to nearly 20% in the edge region at $0.9 < r/a < 1.0$, where multiple counter-propagating modes are sometimes observed [271]. This type of spatial distribution of relative density fluctuation amplitude is nearly universally observed in confinement devices [90] and was reported in early reviews of the subject [6]. It should be pointed out that the origins of this apparently universal observation are, at best,
poorly understood at this writing. These measurements were then used to infer the turbulent diffusivity, which was then compared against diffusivities inferred from a particle and power balance analysis [270] (figure 14). The results were taken to suggest that turbulence plays a determining role in the global confinement of L-mode plasmas.

The same approach was used in the TEXTOR tokamak which was capable of inducing biased H-mode confinement regimes, resulting in increases in the particle and energy confinement. Measurements of the edge turbulence then allowed comparisons of the changes in the turbulent particle flux with changes in global confinement. The results showed that the turbulent particle flux appeared to contribute a significant, but not dominant, fraction of the total particle flux leaving the confined plasma [55]. In work on the CCT tokamak device, the walls were typically coated with a Ti getter, eliminating most of the neutral gas recycling at the wall, and therefore making measurements of the particle confinement time relatively easy by simply terminating the plasma fueling during the discharge and measuring the resulting decay of line-averaged density. Measurements of the poloidally distributed turbulent particle flux were then made with a poloidally distributed array of probes [53] which, when integrated over the LCFS, gave values of particle confinement that were within factors of 2 of the experimentally observed confinement during unbiased ohmic discharges. During biased H-mode operations, the turbulent flux essentially vanished, and the global particle confinement time became very large, qualitatively consistent with the changes in turbulent flux. However, no quantitative comparisons were ever given from that device. Similarly, work in DIII-D provides measurements of the reduction in turbulent particle and heat flux in low-powered H-mode discharges. However, no comparison of these reductions with changes in global confinement was provided. Recently, work on RFP devices has also shown that the inferred turbulent particle flux, when integrated over the boundary of the edge plasma, is sufficient to yield a particle confinement time that is quite comparable to the globally observed particle confinement time inferred from particle balance techniques [272].

Initial estimates of turbulent electron heat flux were made in the ohmically heated CalTech tokamak [7]. In this work, the density, floating potential and electron temperature fluctuation amplitudes were measured. However, no measurements of the cross-phase between these field quantities were available, and thus the results provide only an upper limit on the heat transport caused by electrostatic $E \times B$ heat convection and conduction as well as by magnetic fluctuations. The results suggested that, at the edge of the plasma, the convective heat flux
was consistent with global estimates for plasma confinement. The first measurements of the convective and conducted turbulent heat flux including the relevant cross-phase information were performed in the TEXT tokamak [95] and are shown in figure 15. The work in TEXT took into account the radiated power from the plasma, which represents a significant loss term, thus permitting a direct comparison between the magnitude of the turbulent heat flux, the actual transported heat flux obtained from power balance and the global energy confinement time. The results showed that at the LCFS of these plasmas, the turbulent heat flux was sufficient to explain the loss of heat from the core plasma. However, deeper in the plasma the convective turbulent heat flux was significantly lower than the flux estimated to be exhausted from the plasma core as could be seen in the figure. The implication is then that some other transport mechanism could be responsible for the loss of heat from the central plasma region. Alternatively, the turbulent heat flux measurement technique could be breaking down in this region due, e.g., to probe tip heating effects. Measurements in the DIII-D edge plasma also showed that the turbulent heat flux was of the correct order of magnitude to explain global energy confinement [273]. Recently the turbulent electron heat flux has been measured in H-mode; results are discussed below. To our knowledge, there are no published studies that directly measure ion turbulent heat flux in strongly heated discharges thought to be subject to ITG instabilities.

Thus in a variety of toroidal confinement experiments, the turbulent particle flux appears to be within a factor of 2 or so of the value needed to explain the observed global particle confinement. To our knowledge there are no published results in the literature comparing the parametric scaling of how the turbulent particle flux and corresponding global particle confinement time scale with parameters such as line-averaged density, plasma current and so on. The state of measurements of turbulent ion and electron heat flux is even weaker. Given the importance of thermal transport for burning plasma experiments, additional measurements of ion and electron temperature fluctuations and associated thermal transport would be extremely useful.

5.1.2. Turbulent momentum transport. A number of studies have compared the core plasma momentum confinement and diffusivity against ion thermal diffusivity; these studies have usually used a transport analysis approach that balances calculated momentum input against
observed momentum profiles and then infers an effective momentum diffusivity profile. Reference [274] contains a review paper of this subject. It is usually found that the plasma momentum transport coefficient behaves in a manner that is quite similar to the ion thermal transport coefficients, suggesting a link between the two transport channels. Since ion thermal transport is linked to turbulence, by implication it then appears that plasma momentum transport is also related to turbulence. However, the links are only indirect and to our knowledge there are no published studies of core turbulence and core momentum transport.

In the edge plasma region, evidence linking edge turbulence and the development of large-scale toroidal and/or poloidal plasma flows has recently been obtained in experiments on tokamaks, stellarators and in RFP devices. One such set of measurements have been reported by the CIEMAT group working on JET [275, 276]; figure 16 shows the key results. Working in the edge SOL plasma, these workers reported that strong parallel plasma flows were more common when a large outward burst of particle flux occurred at the LCFS. This would then suggest that the correlation \( \langle \tilde{v}_r \tilde{v}_\parallel \rangle \) between the fluctuating radial and parallel velocity components, which is a component of the Reynolds stress, would be largest during such times. The radial particle transport bursts in turn were observed to be correlated with steep instantaneous plasma density gradients. Together, the results were taken to suggest that the edge plasma turbulence exhibits aspects of a critical gradient instability, where a large outward flux of plasma occurs if the critical density gradient is exceeded. If such events are localized to the low field region (which appears to be the case [53]), they will then result in the subsequent formation of a rapid parallel flow of plasma that then tries to equilibrate the pressure perturbation that has developed on the flux surface, yielding a finite Reynolds stress \( \langle \tilde{v}_r \tilde{v}_\parallel \rangle \) that will then transport parallel momentum radially.

Sanchez and co-workers also measured the net nonlinear transfer of kinetic energy between the turbulence and the large-scale sheared poloidal flows in L-mode JET discharges [277]. In limiter experiments, the results showed that at the LCFS and just outside into the SOL, the turbulent kinetic energy was transferred into the sheared poloidal flow, while deeper inside the plasma, the larger scaled flow transfers energy into the turbulence. In diverted discharges the results showed that the transfer was from the large-scaled shear flow into the turbulence, consistent with the usual picture of turbulent eddy viscosity which damps larger scaled shear flows.

Measurements of the radial–parallel Reynolds stress \( \langle \tilde{v}_r \tilde{v}_\parallel \rangle \) were also reported in TJ-II, with the rate of energy input per unit plasma mass into the large-scale parallel flow of the edge plasma. As the line-averaged density was increased, the turbulent Reynolds stress term increased. This in turn caused a corresponding increase in the rate of energy input into the parallel flow in the region just inside the separatrix. At high line-averaged densities, this rate was estimated to be strong enough to influence the parallel flow in TJ-II. Taken together, these results from JET and TJ-II suggest that the turbulent momentum transport at the plasma boundary influence, and perhaps dominate, the large-scale parallel flows in the edge and SOL regions.

A strong link between edge turbulence and sheared plasma flows has also been reported in RFP devices [38, 118, 119, 278]. In these experiments, the electrostatic component of the Reynolds stress was the dominant momentum transport mechanism in the edge region of the RFP, and the divergence of this Reynolds stress was consistent with the measured radial profile of sheared toroidal plasma flow. The transfer of kinetic energy from the smaller scaled turbulent fluctuations into the larger scaled sheared flow, analogous to the energy exchange measured in TJ-II, was also measured [19]. The results showed that the turbulent kinetic energy is transferred into the larger scaled sheared toroidal \( E \times B \) flow in the region located just inside the separatrix in agreement with the TJ-II results.
We conclude that results from tokamaks, stellarators and RFPs show a significant nonlinear exchange of momentum and kinetic energy between the edge turbulence and the large-scale shear flows that are universally observed at the edge of confinement devices. A number of results suggest that at or near the LCFS the turbulence could actually drive the shear flows in ohmic or L-mode discharges. The initial results from JET indicate that this picture is sensitive to the type of discharge (i.e. limiter versus divertor discharges), and the results from TJ-II indicate that this result is also dependent on the edge density. There do not appear to be any such studies from the core plasma region. In light of recent work in ALCATOR C-MOD linking edge and SOL flows with core plasma rotation and global thermal energy confinement [279–281], additional studies of the role of turbulence momentum transport in determining plasma rotation and flow should be strongly encouraged.

5.2. Origins of core plasma turbulence

The above results show evidence suggesting a link between the turbulence and the resulting rate of particle, momentum and heat transport in the plasma across a range of conditions including ohmically heated, auxiliary heated L-mode and H-mode discharges. However, these results say nothing about the underlying drive mechanisms that lead to the turbulence. As seen in the summary of drift instability theory, in principle this task requires identification of the free energy source(s), parallel electron dissipation mechanism(s), knowledge of the nonlinear interactions that result in the generation of turbulence on the observed spatial and temporal scales, as well as knowledge of the cross-phase between the different fluctuating field quantities and the resulting cross-field flux of particles, momentum and heat. These mechanisms could vary significantly depending upon the experimental conditions. Typically in low density ohmically heated discharges that are dominated by electron heat transport, comparisons with either the universal or dissipative trapped electron drift instabilities or to non-drift wave instabilities are made. We refer the interested reader to earlier works for discussions of these plasma conditions and studies [22, 90]. In higher density ohmic discharges or in auxiliary heated discharges, comparisons with either ITG or collisionless TEM instability theory have usually been made. The possible role of ETG drift instability in electron heat transport is currently under active investigation.

In practice a complete identification of the origin of the turbulence as described above is not possible due to diagnostic limitations. Instead, identification of the type of instability is usually made from a linear stability analysis of the experimental profiles which takes into account the stabilizing effects of equilibrium flow shear, magnetic shear and flux surface compression via the Shafranov shift. Such an analysis is then often followed by an attempt to predict some characteristics of the turbulence itself, such as the wavenumber spectrum, the dispersion relation and so forth using either a fluid, gyro-fluid or gyro-kinetic numerical simulation performed in either flux tube or global geometry. In a few cases, turbulence measurements are available and are then compared against these modeling results in order to then identify the origin of the turbulence. This procedure assumes a primary free energy source and parallel electron dissipation mechanism exists, and that the plasma does not significantly deviate from the time-stationary equilibrium assumed in the stability analysis and quasilinear or turbulent modeling. If these assumptions break down in the experiment, then the analysis results may be questionable. With these caveats in mind, let us summarize key observations of the origins of turbulence in more experiments.

Achievement of fusion conditions requires significant auxiliary heating and thus transport in these conditions is of interest. The ITG instability has been implicated to drive turbulence and transport in the core region of several tokamak experiments operating in the auxiliary
heated L-mode confinement regime [23, 282–284]. We summarize results from L-mode experiments on DIII-D that were taken to demonstrate the existence of fluctuations driven by ITG instability [285] in plasmas where the complicating effects of flow shear, magnetic shear and Shafronov shift stabilization were minimized. In these experiments, the density was systematically varied in a series of ohmically heated discharges, and fluctuation measurements were obtained with a far infrared scattering system and correlation reflectometry. The density spanned from a low value that exhibits linear ohmic confinement (LOC) to a high value that exhibits saturated ohmic confinement (SOC), as originally identified in the neo-Alcator confinement scaling studies formed on DIII-D [286] which were motivated out of earlier observations on other tokamaks [80]. The density profiles for the low and high densities showed a broad profile at high density with a large density gradient near the outer region of the plasma, and a more monotonically decreasing density with radius for the lower density condition. At low density, the global energy confinement is proportional to line-averaged density; this scaling then saturates and perhaps begins to decay at higher line average density (figure 17(a)). At the conditions where confinement level saturates the fluctuation level at $k_0 = 2 \text{ cm}^{-1}$, observed with an FIR scattering diagnostic which integrates fluctuation levels over the core region of the plasma, exhibits a marked increase as the saturated confinement regime is reached (figure 17(b)). This result was taken as suggesting a correlation between the increased fluctuation amplitude and the saturation of the confinement improvement.
Spectral measurements of the density fluctuations are shown in figure 18. Of particular note was the increase in low-frequency ($f < 200$ kHz) fluctuations that were observed when the high density SOC regime is reached. A baseline level of fluctuations over a wider frequency range was observed in both the LOC and SOC regimes; this feature was apparently unchanged as density is increased. The central frequency of the low-frequency feature was positive, which was consistent with fluctuations propagating in the ion diamagnetic direction, as expected for ITG turbulence. Corresponding measurements near $r/a = 0.6$ with a poloidal correlation reflectometer showed the fluctuation direction changing from the electron diamagnetic direction at low density to the ion direction at the higher density, consistent with the scattering results. Importantly in these experiments the $E \times B$ Doppler shift was measured, allowing the authors to rule out a lab-frame propagation direction change simply due to $E \times B$ drift effects. Linear stability analysis showed that at the higher density conditions, the growth rate is dominated by ITG instability, since the threshold for instability was exceeded by the mean gradients. Nonlinear simulations likewise showed significant linear stability variation as the plasma density was changed, with the ITG mode being stable in the low density discharges, and becoming unstable at high density; ITG-driven turbulent energy flux likewise increases markedly in the higher density SOC regime in the simulation.

Further evidence for the existence of ITG-driven turbulence comes from localized measurements of turbulent ion temperature fluctuations in TFTR using the HF-CHERS diagnostic [197] which were used to help identify turbulent driving mechanisms [23] in these experiments. Models predicted that normalized electron and ion temperature fluctuations, $\tilde{T}_e/T_e$ and $\tilde{T}_i/T_i$, would behave in a specific manner relative to normalized density fluctuations, depending on the dominant instability driving the turbulence (e.g. trapped electron modes versus ITG-driven modes). Thus ion temperature fluctuations were measured using HF-CHERS and compared with density fluctuations. It was found that the normalized ion temperature fluctuations exceeded the normalized density fluctuation amplitude by a factor of $\sim 2$ as expected for ITG-driven turbulence. This result was found across the outer region of the plasma $0.6 < r/a < 1.0$, consistent with the expectations of an ITG stability analysis of this
plasma equilibrium, suggesting that the ITG instability was responsible for the turbulence in these beam heated discharges.

There is some limited evidence that under some conditions the turbulence in the core region is consistent with the TEM turbulent drive. Doppler reflectometry measurements in ASDEX-UG [282] were used to search for such evidence by increasing the plasma density in ohmically heated discharges. No independent measurements of $E \times B$ Doppler shifts were available; thus theory was used to estimate this Doppler shift, and then the plasma-frame propagation was estimated from lab-frame propagation velocity measurements and computed plasma flows. The resulting inferred change in the plasma-frame propagation direction of the turbulence at elevated collisionality was interpreted as indicating a transition from TEM instability to ITG instability—an interpretation that would seem to be consistent qualitatively with the results from Rettig et al discussed above that showed evidence for onset of the ITG instability at higher densities (and thus presumably higher collisionality). This would also suggest that the turbulence reported by Rettig et al in the low density LOC regime in DIII-D could be due to TEM fluctuations. Unfortunately no examination of that possibility was provided in the work of Rettig et al. In the ASDEX work, given that the key piece of evidence for TEM—namely the plasma-frame propagation direction—is not measured but rather is inferred, it would seem that additional studies are needed to confirm the hypothesis that lower density ohmically heated tokamaks have TEM turbulence.

Additional evidence for the existence of TEM turbulence is provided by PCI measurements of the density fluctuation wavenumber spectrum in the ALCATOR C-MOD experiment [287]. The measurements were obtained in conditions where the TEM was computed to be the dominant linear instability in the core plasma. The line-integrated density fluctuation spectrum was found to be in good agreement with the spectra computed using gyro-kinetic simulations as shown in figure 19. However we note that these measurements are line integrated through both the core, edge and SOL plasma regions. The simulations only examine the turbulence in the core plasma conditions and do not contain any relevant information from these other locations. Thus if edge and/or SOL fluctuations give any significant contribution to the measured spectrum, the resulting agreement could then be fortuitous.
When strong electron heating occurs, the ETG could be strong enough to trigger the ETG instability, which is distinct from the ITG and TEM class of drift turbulence due to the much smaller scaled fluctuations that are expected. This regime is of particular importance, since in future burning plasma experiments the thermal transport will occur via both electron and ion channels. Recent work indicates that small-scaled ($k_⊥\rho_i > 1$) fluctuations consistent with ETG-driven turbulence exist in some core plasmas [288]. In this work, electron cyclotron heating (ECH) was used to heat the electrons, and the response of the large-scaled ($k_⊥\rho_i \lesssim 1$), intermediate scale ($k_⊥\rho_i \sim \text{few}$) and small-scaled turbulence was monitored. The results showed a marked increase in the amplitude of $k_⊥\rho_i > 1$ fluctuations that was concurrent with an increase in the electron thermal diffusivity. These increases both occurred in the outer regions of the core plasma, consistent with linear gyro-kinetic ETG stability calculations; no such change in the larger and intermediate scale turbulence was observed, suggesting that the high $k$ turbulence changes could not be attributed to changes in this part of the turbulence spectrum. We note that very recent work in NSTX using high harmonic fast wave heating of the electrons has also provided evidence for the existence of ETG-scaled turbulent fluctuations [289] when the plasma is calculated to be linearly unstable to the ETG drift instability. The contribution of these fluctuations to the electron thermal transport is still to be determined.

In summary, results from tokamak devices based upon measurement of turbulence spectra, fluctuation propagation direction and linear stability analysis along with comparisons with turbulence simulation suggests that a combination of ITG and TEM instabilities could generate core turbulence and transport. There is initial evidence of higher wavenumber ETG-driven fluctuations as well, but the significance of the contribution of these fluctuations to transport is not yet clear. Evidence exists to support the hypothesis that the onset of ITG turbulence corresponds to a reduction in confinement. It would be very useful to continue such studies where independent plasma rotation measurement allows the Doppler shift effects to be accounted for. In addition, simultaneous multi-field measurements would allow the cross-phase to be measured and used for identification of turbulence origins as has been recently done in DIII-D [196]. However, the reader should note that the ITG and TEM wavenumber spectra may sometimes overlap [290] and making $E \times B$ flow measurements with sufficient resolution and accuracy to resolve diamagnetic flows is challenging, making a unique attribution of the fluctuations to a particular linear instability difficult in many cases.

5.3. Origins of edge and SOL plasma turbulence

Experiment shows that the edge and SOL turbulence consists of bursts of very large fluctuation amplitudes with strongly non-Gaussian statistics (figure 20 shows a GPI image of edge/SOL turbulence [164–167]) which, as shown by HIBP measurements in ISX-B, make a drastic departure from the Boltzmann equilibrium for the electrons in the SOL (figure 21) [20]. Langmuir probes have been used on tokamak devices to study the PDFs and associated moments of these bursts and their average behavior using conditional averaging techniques [105, 127, 291]. The results show that positive bursts of density propagate radially outward into the SOL while negative bursts of density (i.e. ‘holes’) propagate inward toward the core region, yielding a density fluctuation skewness profile that changed sign at the LCFS (figure 22) [105]. Work in the ALCATOR C-MOD tokamak [291] using a combination of probes and visible light fluctuation diagnostics showed that these density structures propagate as a result of the $E \times B$ drift associated with the dipolar potential structure of the events (see figure 23). These tokamak observations are strikingly similar to those found in the linear experiments summarized earlier, suggesting a common set of underlying physics mechanisms. These results also suggest that the large-amplitude density perturbations being generated at the LCFS cause a significant effect
Figure 20. Left panel: fluctuating He I light emissions at the edge region of the NSTX tokamak. Right panel: radial profile of relative emissivity fluctuation amplitude; reproduced with permission from [167].

Figure 21. Radial profile of the ratio of normalized plasma potential fluctuations to normalized density fluctuations in the ISX-B experiment. The ratio makes a rapid deviation from a value of $\sim 1$ in the edge region, indicating non-Boltzmann parallel electron dynamics in this region. Upper panel: ohmic discharge; lower panel: L-mode discharge; reproduced with permission from [20]. Copyright 1987 American Physical Society.

at least several centimeters inside the LCFS, and thereby couple the SOL turbulence to the core region.

The origin of the turbulence in the SOL is thought to be fundamentally distinct from that of the core region due to the fact that the field lines terminate on resistive sheaths tied to the
Figure 22. Left panel: conditionally averaged density and poloidal electric field measurements showing that positive going density events propagate radially outward, while negative going density holes propagate inward. Right panel: the skewness of the density turbulence changes sign at the LCFS, showing that the outward (inward) going density bursts (hole) are being generated at the LCFS; reproduced with permission from [105]. Copyright 2003 American Institute of Physics.

grounded conducting divertor or limiter material surfaces [25, 26, 292], thereby introducing a new dissipation mechanism for the electrons that does not depend on collisions within the SOL plasma volume [26, 27]. Comparisons with SOL turbulence data in ASDEX [293] showed that this model accurately reproduced the cross-phase between potential and ion saturation current fluctuations, that the most unstable range of poloidal wavenumbers corresponds to the observed correlation lengths in the experiment and that the poloidal phase velocity of the fluctuations agreed with this model. This basic physics model has been incorporated into turbulent fluid simulations of the SOL region [125–127] and provided good agreement with the radial variation of the normalized fluctuating density amplitude, the PDF of the density fluctuations, of the cross-field particle flux and of the effective diffusivity [126] in TCV, suggesting that this fundamental picture of turbulence at the LCFS and in the SOL is likely correct. Furthermore, this mechanism would cause the turbulence to be generated on the low-field side of the tokamak as has been reported in several experiments [53, 120, 122]. If the asymmetric cross-field transport from this mechanism is fast enough relative to parallel flow timescales then it could lead to the development of pressure gradients within a magnetic flux surface which would then drive parallel flows and fluctuations within the flux surface [54] in a manner reminiscent of the JET and TJ-II momentum transport results discussed above.

It is important to note that even in ohmic discharges a mean $E \times B$ velocity shear maintained by the plasma equilibrium exists at the LCFS [8]. Work in the TORPEX device has shown that such flows could act to pinch off, or shear, the developing instability at the
LCFS [210, 260] and thereby help to regulate the magnitude of the edge fluctuations. It seems possible that if the magnitude of the shear flow were to increase at the LCFS due to heating or momentum input into the plasma, then this increased flow shear would modify the cross-field transport. Thus the conditions necessary for the development of the L–H transition seem to be naturally created at the LCFS. Another important outcome of this turbulence generation mechanism is the development of an outward convective velocity to the structures which then leads to a non-diffusive convective transport of plasma toward the first wall [105, 166, 167, 249–251, 253, 294–297]; this could then lead to strong plasma–wall interactions with possible deleterious effects on the wall and on the plasma. These observations and interpretations have emerged in the period since the earlier reviews of edge turbulence in confined plasmas [22, 90, 91, 265, 298] and thus represent a significant change in our understanding of edge and SOL turbulence.

5.4. Transition between edge and core turbulence

There are observations that show that the turbulence from the edge/SOL region has an impact on the outer regions of the core plasma flux surfaces, as evidenced, for example, by the large fluctuation amplitude not only in the SOL but also in the region immediately inside the LCFS, the rapid deviation from the Boltzmann relation for density and potential fluctuations in the same region and the penetration of the holes that are generated at the LCFS several centimeters into the closed flux surface region. At this writing there is no clear experimental understanding of how the core plasma fluctuations transition in a continuous manner into the fluctuations that occur near the LCFS and on open field lines, nor of how these regions might be dynamically coupled to each other. Given the fact that it is precisely this region where the L–H transition occurs, where the H-mode pedestal forms and where edge localized modes (ELMs) are formed, studies of how the turbulence evolves and is coupled across the edge/LCFS region would seem to be critically important in understanding the origin of the L–H transition.

In TFTR Durst et al reported on how the turbulence characteristics change across the core–edge plasma interface [271]. Fluctuations were observed to propagate in both the ion and electron diamagnetic drift directions in the edge region of L-mode discharges in this circular limiter tokamak plasma; these two distinct types of fluctuations exhibited a significant variation with radius. At $r/a = 0.74$ the fluctuations were observed to propagate in the ion drift direction, and have a poloidal wavenumber spectrum that peaked at $k_\theta = 1 \text{ cm}^{-1}$. Closer
to but still inside the LCFS a low-frequency set of fluctuations propagated in the ion drift direction, while in the same region at higher frequencies the fluctuations propagate in the electron drift direction. In the SOL region, this edge ion mode region was observed to have the largest amplitude, while the edge electron mode amplitude decreased to small values. At a location about 5 cm inside the LCFS, the amplitude of the edge ion mode was observed to fall to a small value inside this region; deeper inside the fluctuation amplitude of the ion mode recovered. The radial correlation length was also reported to show an abrupt drop in the interface region 5 cm inside the LCFS. Thus it would seem that the fluctuations in the core, edge and SOL regions, while exhibiting different types of fluctuations, made a continuous and smooth transition across these regions, and influence and couple to each other.

Results from DIII-D also show the existence of such counter-propagating turbulent fluctuations in the region just inside the LCFS [299]. Two poloidally separated BES density fluctuation channels show that fluctuations propagate in both electron and ion directions near the LCFS. As the observation location is moved deeper into the plasma, the relative strengths of the two counter-propagating fluctuations changes (figure 24). These experiments also showed that the ion grad-$B \times B$ drift direction had a strong effect on these counter-propagating turbulent fluctuations. In particular, it was reported that a time-averaged mean shear layer existed in discharges with a lower single null (LSN) which then disappeared in upper single null discharges. A reduction in the radial correlation length of the edge turbulence was also observed to occur at the location of the shear layer in LSN discharges, in qualitative agreement with theoretical expectations from flow shear decorrelation effects. As discussed above, this localized reduction in radial correlation length is similar to that reported from TFTR above by Durst et al.

These results may have a relationship to the L–H transition. Recent C-Mod results show that the L–H transition power threshold (which corresponds to the heat flux transported through the plasma edge) is a sensitive function of the divertor geometry [300] (i.e. LSN/USN/double null) with the L–H transition threshold being lower when the ion grad-$B$ drift direction is toward the single null divertor X-point. This is precisely the condition where the L-mode turbulence was found to have a flow reversal and shear layer [299]. In addition, as pointed out elsewhere in this review, this region of the plasma has GAM-type oscillatory shear flows present and perhaps these dynamics play a role in the counter-propagating modes in the edge/core transition region.

These somewhat fragmentary observations would seem to suggest that there may indeed be a link between these various effects. Given the potential impact on the L–H transition power threshold and global confinement, it would seem that the subject of the coupling between core and edge plasma turbulence, and the impact of plasma shaping and divertor geometry upon that coupling, deserves significant additional study.

5.5. Variation of turbulent scale length with $\rho^*$

The heat flux across a magnetic surface could be written in terms of a thermal diffusivity $\chi = \rho_2^2 \omega^\ast (G(p_\parallel)/\langle k_{\perp}\rho_S \rangle)$ and the gradient of a temperature [301]; in this expression the factor $\rho_2^2 \omega^\ast$ arises from the natural step size and step time of drift fluctuations, while the dimensionless factors $G(p_\parallel)$ and $\langle k_{\perp}\rho_S \rangle$ denote the variation of the diffusivity with the dimensionless plasma parameters and the mean scale size of the turbulent fluctuations, respectively. Current experiments can match all of the dimensionless parameters that will be found in a burning plasma experiment except for the dimensionless scale length $\rho^* \equiv \rho_S/a$. A key issue then arises: how does $\langle k_{\perp}\rho_S \rangle$ scale with the system size, $a$, which is usually measured in terms of the dimensionless scale length $\rho^*$. If $\langle k_{\perp}\rho_S \rangle$ is independent of $\rho^*$ then turbulence scale size can only depend on local dimensionless quantities and results in the so-called gyro-Bohm scaling
and yields a favorable scaling of confinement time with both magnetic field and system size. If the turbulence scale size increases as the system size increases (i.e. if $\langle k_\perp \rho_S \rangle$ is proportional to $\rho^*$) then the resulting thermal diffusivity gives Bohm scaling which has much less favorable scaling with magnetic field and system size. Thus, knowledge of how the thermal diffusivity scales with $\rho^*$ is important in that it provides insight into the quality of confinement that will be obtained in such next-step experiments.

There are two basic experimental approaches to addressing this key question. First, one could infer the electron and ion thermal diffusivity from experiments in which $\rho^*$ is varied while the other dimensionless parameters are held constant and then infer how the thermal diffusivity and thus turbulent scale length vary with $\rho^*$. The success of this approach is determined in large part by the degree to which these other parameters could be held constant while varying $\rho^*$.

Many different authors have reported on such studies in L-mode and H-mode discharges using both single fluid and two fluid power balance analyses [302–306] with significantly varying conclusions as to the $\rho^*$ scaling of turbulence correlation lengths. As this approach does not examine the turbulence characteristics first-hand, it falls outside the scope of this review and we refer the interested reader to the literature.
The second approach to addressing the $\rho^*$ scaling is to directly measure the turbulence scale length variation with $\rho^*$. This approach has been used on DIII-D with fluctuation measurements provided by the BES diagnostic [307], in TJ-K using a probe array [308], and on Tore Supra using laser scattering to measure the turbulent wavenumber spectrum [309] (which is just the Fourier transform of the spatial correlation function). Let us summarize these important experiments.

In DIII-D, a pair of discharges were created with $\rho^*$ varying by a factor of 1.6 ($B_T$ was varied by a factor of 2), while all other transport-relevant dimensionless parameters are held nearly constant. Over this small variation in $\rho^*$ the radial correlation length normalized to the local ion gyroradius (figure 25) was roughly independent of $\rho^*$ consistent with a gyro-Bohm scaling, and had a value of approximately $L_{c,r} = 5\rho_I$. The decorrelation times were measured multiple poloidally displaced channels to find the true plasma-frame decorrelation time. The decorrelation times, which are approximately 10 $\mu$s, varied with $a/c_s$ as expected and are of the same magnitude.

In Tore Supra dimensionless $\rho^*$ scaling experiments, measurements of higher wavenumber fluctuations ($3 \text{ cm}^{-1} < k_{\perp} < 26 \text{ cm}^{-1}$) with CO$_2$ laser scattering show evidence that the local turbulence statistics are consistent with a gyro-Bohm scaling, i.e. that the turbulence scale size is independent of the system size [310]. In this work, the scattering measurements were used to examine the variation of wavenumber spectrum and decorrelation times scaled with $\rho^*$. The results (figure 26) showed that the shape of the fluctuation spectra normalized as $|n(k)|^2/n^2/\rho_I^2$ and the decorrelation time normalized to the transit time $a/c_s$ did not change.
with $\rho^*$. This result was interpreted as being consistent with a gyro-Bohm scaling for the turbulent correlation length. Interestingly, these experiments also showed that there was a transition in the spectral dependence of fluctuation power from a $k^{-3}$ dependence at lower wavenumber to a much stronger decay of $k^{-6}$ above approximately $k_\perp \rho_i = 1.5$, hinting that these higher wavenumber fluctuations may not play a significant role in transport in these discharges due to vastly reduced power available above $k_\perp \rho_i = 1.5$.

In the TJ-K device, a probe array was used to measure the density fluctuations as the plasma composition was changed from hydrogen to helium, neon and finally into argon, allowing a range of 10 in $\rho^*$ to be explored [308]. For the smaller values of $\rho^* < 0.1$, the results gave roughly a gyro-Bohm scaling for the turbulence scale length; in heavier ion plasmas with larger values of $\rho^*$, the scaling was closer to a Bohm-like $\rho^*$ scaling, suggesting if the device size is too small, then the net result is a Bohm-like transport scaling while, if the device scale is large enough, the expected gyro-Bohm scaling is recovered. The results also showed that the cross-phase between density and potential varied with ion mass for unknown reasons. No plasma equilibria were reported in this work and thus the possible role of changes in collisionality in these results is unclear at this time.

These results suggest that theoretical expectations for gyro-Bohm transport scaling of the turbulence scale lengths are likely correct. However, this interpretation might be complicated by the so-called critical gradient behavior wherein the diffusivity is small for small temperature gradients and then, when the critical temperature gradient is exceeded, the thermal diffusivity exhibits a rapid increase. Such behavior should be accompanied by significant changes in the turbulence scaling lengths, and perhaps in the parametric scaling of the turbulent correlation lengths as the critical gradient is exceeded. Perhaps the Bohm/gyro-Bohm scaling might then depend upon how far the critical gradient is exceeded. Computations also suggest that the system size could influence the turbulence size scaling as well [311] in a manner that is qualitatively similar to these TJ-K results. There are no data at present available to determine whether this is indeed the case, but it should be possible to perform such a test in existing devices. Given the significance of this issue for ITER performance projections, it would seem that additional direct study of the scaling of turbulent structure size across a wider range of experiment should be encouraged.
5.6. Existence of ZFs in confined plasmas

As discussed earlier, theory predicted that drift turbulence should nonlinearly generate large-scale ZFs within a magnetic flux surface, that these flows should act as a nonlinearly driven sink of turbulence kinetic energy and should act to partially decorrelate the turbulence via the shearing action of the ZF and thereby play a role in the saturation of turbulence and transport for a given mean pressure gradient. Experiment has verified some (but not all) aspects of this theoretical picture; we summarize representative results here. We also point out that since the first submission of this review paper, a separate review paper focused specifically upon experimental studies of ZFs has been accepted for publication elsewhere [47].

The first experimental evidence for such ZFs was reported in work on the H-1 stellarator device [312–314]. A low-frequency radially sheared, radially localized electric field was observed. This fluctuation had a long correlation length within the magnetic flux surface, and was nonlinearly coupled to higher frequency waves. This shear flow did not produce any transport itself, and had an amplitude that was modulated by corresponding modulates in the higher frequency waves. These results were obtained under conditions where the ions were nearly unmagnetized and as a result are in a different regime than the fluctuations found in larger scaled, higher magnetic field fusion experiments which tend to be dominated by $E \times B$ dynamics.

The first identification of stationary ZFs in conditions where the ions were magnetized was performed in a medium size toroidal device, CHS ($R = 1$ m, $a = 0.2$ m), using HIBPs [171]. The device was equipped with two HIBPs, located by 90° apart in toroidal direction, as is shown in Figure 27 to confirm a long-range correlation of ZFs. Each HIBP has three channels, thus, could measure the electric field directly by making a difference between potentials at two channel locations.

The measurement of electric field fluctuation was carried out in the plasmas produced with electron cyclotron resonance heating of $\sim$200 kW in low-beta plasma ($\sim$0.2%) to avoid
MHD activities which may interfere the ZF detection. In this discharge condition, the density was kept constant approximately at $n_e \sim 5 \times 10^{12}$ cm$^{-3}$, electron and ion temperatures $T_e \sim 0.5$ keV, $T_i \sim 0.1$ keV (i.e. in collisionless regime), ion Larmor radius $\rho_i \sim 0.1$ cm, time scale of micro-instabilities $\omega^*/2\pi \sim 50$ kHz with $k_{\perp}\rho_i \sim 0.3$ and energy confinement time $\tau_E \sim 2$ ms (or characteristic frequency of global confinement $\tau_E^{-1} \sim 0.1$ kHz), where $k_{\perp}$ is the wavenumber and $\omega^*$ is the drift frequency defined as $k_{\perp}T_e/eBLn$ with $L_n$ being a characteristic length of density gradient.

Figure 27 also shows a spectrum of electric field fluctuation at the position of $r = 12$ cm (or $\rho \sim 0.6$), which gives a maximum signal-to-noise ratio to the HIBP measurement in this condition. The spectrum shows the existence of fluctuation of the lower frequency of less than $\sim 1$ kHz, a peak at 17 kHz and a broad-band peak around the drift wave frequency of $\sim 50$ kHz. As shown in the figure, the coherence (the blue line) shows that the low frequency fluctuations with frequency $f < 1$ kHz have a long-range toroidal and poloidal correlation. Consequently, the fluctuations of this low frequency range were considered to be ZFs, while the peak at $\sim 17$ kHz was conjectured as a GAM. The evaluation of the absolute values of the electric field strength of the low-frequency fluctuations shows that the fluctuation amplitude is approximately $\sim 1$ V, and the corresponding amplitude of the stationary ZF electric field is 0.05–0.1 kV m$^{-1}$ (corresponding to an $E \times B$ drift speed of 0.06–0.11 km s$^{-1}$).

The electric field fluctuations at two toroidal locations were coherent but had a finite constant phase shift that was proportional to the radial separation of the two toroidal observation points. Therefore, the radial structure of the fluctuation could be deduced by making radial correlation function, which indicates the phase between electric fields at two observation points. Figure 27 shows the radial correlation coefficient (closed circles) as a function of the radial distance of observation positions $r_1 - r_2$. The result shows a quasi-sinusoidal ZF radial structure with a radial wavelength of $\sim 1.5$ cm. Accordingly, the fluctuations of the low frequency range show the zonal structure with a radial wavelength that lie between $\rho_S$ and $L_n$.

The BES diagnostic on DIII-D was used to infer the coherent advection of small-scaled density fluctuations by the larger scaled GAM ZF that exists in the edge region of L-mode DIII-D discharges [367]. These flows were radially sheared and had a long poloidal scale length (see figure 28) with a peak amplitude in L-mode at a normalized minor radius $r/a \sim 0.9$ for the conditions of the experiment; there are no published results from H-mode available. The frequency of these GAMs in DIII-D and ASDEX-UG showed good agreement with theory over a wide range of operating conditions [149]. Recently these velocimetry techniques were applied to BES data from the core region of L-mode DIII-D discharges using a lower noise system. The results showed the existence of a zero-mean frequency ZF in this region of the plasma [162]. In related work, a multi-point probe array was used to measure radial electric field fluctuations within a flux surface for large poloidal and toroidal separations in the HL-2A tokamak [315]. The results show that the GAMs also have a toroidally symmetric spatial structure, consistent with theoretical expectations. Thus in experiments across a number of confinement devices, the existence of shear flows that are consistent with the expected spatio-temporal structure of zero-mean frequency ZF in the central plasma region, and with the GAM in the edge plasma region has been demonstrated. However, these results are silent as to the origins of these flows, and upon their impact on the turbulence. This is the subject of nonlinear turbulence dynamics which we take up next.

5.7. Nonlinear dynamics of turbulence

The development of a saturated power spectrum from a linear drift wave instability necessarily involves the re-arrangement, or transfer, of fluctuation energy across a range of spatio-temporal
scales by nonlinear mechanisms as was summarized earlier. The first attempt to study nonlinear energy transfer in tokamak turbulence was carried out on the turbulence data in density fluctuations in the plasma edge region in the TEXT tokamak [134]. The model used for such measurements is naturally cast in the wavenumber domain. However, the necessary wavenumber measurements require multi-point detection, which was not possible in this experiment. Instead, in this experiment, the density fluctuations were measured at two positions with the Langmuir probes separated by a small poloidal distance at a location \(\sim 1\) cm behind the limiter. The spectral analysis was performed on \(\phi(f, x)\) instead of \(\phi(k, t)\), and thus Ritz and co-workers implemented this technique with probe data in the edge region of TEXT using a frozen flow type hypothesis where the wavenumber could be translated into a frequency dependence. The analysis deduced the essential parameters such as dispersion relation \(k_f\), the growth rate \(\gamma_f\) and the power transfer function \(T_f(f_1, f_2)\) to demonstrate the cascading process, and demonstrated that the fluctuations in the frequency range from 30 to 110 kHz were linearly unstable while the other frequency ranges are linearly damped. The energy in these unstable frequencies was then inferred to be transferred predominantly into the linearly stable fluctuations with frequencies in the range 10–40 kHz.

Ritz’s energy transfer analysis was modified to include the effects of statistical noise [316] and was then applied to density fluctuation data obtained with BES for the supershot discharges in the TFTR experiments. In these experiments, the observed frequency was dominated with the Doppler shift due to the background plasma flow, i.e. \(\omega = \omega_0 + k_\theta v_\theta \sim k_\theta v_\theta\). Therefore, in this situation, \(\phi(f, t)\) was regarded as being a proxy for \(\phi(k, t)\) since the frequency is directly proportional to the poloidal wavenumber [3]. The results of this spectral transfer analysis on the plasma core at the normalized radius of 0.7 in the so-called supershot plasmas in TFTR suggested that the peak in the fluctuation spectrum located at \(k_\theta \rho_s \sim 0.07\) corresponds to the
region of linear stability, implying that the fluctuations of this regime were nonlinearly driven by modes at higher wavenumber. The results showed a clear difference in the linear growth rate and nonlinear coupling strength for different beam heating configurations.

Recent work in the TJ-K stellarator device has examined the nonlinear energy transfer of both kinetic energy and density fluctuations via the use of a 2D spatial array of probes to measure the density and potential fluctuations across the entire poloidal cross-section of that device [317]. The results (see figure 29) show the existence of a dual energy transfer process; kinetic energy is transferred into the large spatial scales while the energy associated with the density fluctuations are transferred into the larger wavenumber range. The results were compared with numerical simulation results with reasonable agreement suggesting, for the first time in a plasma experiment, that essential elements of a dual energy cascade in two-dimensional turbulence expected from theory [15, 318] could actually be observed in experiment. It is interesting to note that these results are qualitatively consistent with expectations of the emergence of a large-scale ordered flow such as a ZF, as well as the shearing action of such a flow upon the density fluctuations.

The initial applications of nonlinear energy transfer in TEXT and TFTR were carried out prior to the discovery that ZFs could play an important role in the regulation of drift turbulence and thus no discussion of this physics is contained in these earlier papers. Furthermore, the single field model is of questionable validity, particularly for the conditions of the TEXT experiment, and thus the conclusions of these initial studies of nonlinear drift turbulence dynamics should not be taken to be conclusive. Nonetheless, this approach—of directly examining the underlying nonlinear dynamics and linear growth rate from experimental data—would appear to provide great promise in identifying the origins of the turbulence and providing a deep test of numerical turbulence simulations provided that the necessary experimental data could be obtained.
5.8. Nonlinear drift turbulence–ZF interactions

The nonlinear transfer of energy across different spatial and temporal scales is a key element of the saturation of drift waves into a state of turbulence, of the generation of ZFs from the turbulence and of the back reaction of the ZFs upon the turbulence. Here we summarize existing experimental attempts to study these important elements of the drift turbulence problem in toroidal confinement experiments.

5.8.1. Nonlinear drive of ZFs by turbulence. Recently, both the spectral transfer analysis discussed above and amplitude correlation technique were carried out for the floating potential fluctuation data in the H1-heliac [319]. Results in L-mode showed that energy flows from the unstable frequency regime (20–50 kHz) to the lower frequency regime (0–15 kHz) that is associated with a ZF. Application of an amplitude correlation technique confirmed that energy from higher frequency unstable modes was transferred into lower frequency stable modes. It would be extremely interesting to apply similar techniques to the exchange of energy between turbulence and ZFs in more highly magnetized plasmas found in fusion experiments or to repeat the Reynolds stress measurements that have been performed in laboratory plasma devices. To our knowledge, no such studies have yet been reported in the literature.

5.8.2. Nonlinear effect of ZFs on turbulence. The ZFs also have a nonlinear effect back upon the turbulence. HIBP measurement in JFT-2M revealed the existence of a GAM at the plasma edge, which manifests its presence as a sharp peak at ~15 kHz in the spectra of potential and density fluctuations [320]. Figure 30 shows the temporal evolutions of the band-pass filtered potential fluctuation amplitude at the GAM frequency, and that of the wavelet spectrum of the density fluctuations. The density fluctuations appear to be modulated by the GAM frequency; Fourier analysis confirms that the density fluctuations ranging from 100 to 120 kHz is modulated at the GAM frequency and that this modulation has a significant coherence with the GAM. The biphase computed between the turbulence and the GAM was shown to have a value around \( \pi \), consistent with theoretical expectation that the modulational instabilities should be the cause of ZF generation [39]. Furthermore, the coupling of the modulation to the GAM, as measured by the bicoherence, increases as the GAM amplitude becomes larger. This tendency is consistent with the theoretical prediction based on the Hasegawa–Mima equation, and the experimental proportional constant is found to agree with the theoretical prediction within the experimental error bars [131, 132, 321]. A bicoherence analysis has often been used to prove the coupling between GAM and background turbulence in several experiments such as H1-heliac [313], JFT-2M [131] and HL-2A [315].

In addition, HIBP measurements in the CHS device confirmed, using wavelet analysis, that the turbulent electric field power (which corresponds to the turbulence kinetic energy) in the drift frequency range \( 30 < f < 100 \) kHz was modulated by the low-frequency ZFs with \( 0.25 < f < 1.5 \) kHz [322]. The turbulence power was found to increase (decrease) when the ZF power decreased (increased), qualitatively consistent with a conservation of total power in the coupled turbulence–ZF system. Wavelet analysis also showed that the strength of the coupling between the ZF and turbulence is maximized when the ZF amplitude is peaked [207].

Recent work on DIII-D showed that the nonlinear interaction of the GAMs with the turbulence in the edge region could act to regulate the turbulence scale lengths [323]. In this work, the spatio-temporal data from the BES diagnostic were used to give the density fluctuation spectra and velocimetry techniques gave the GAM velocity. The nonlinear interaction between the GAM shear flow and the density fluctuations was then studied using the power transfer and ACT techniques. The results showed that a measurable finite cross-coherence exists between
these two signals, indicating that there is a statistical relationship between GAM amplitude and turbulent density fluctuation amplitudes. In addition, the phase of GAM velocity amplitude was found to lead the phase of the density fluctuation envelope, indicating that an increase in the GAM shearing rate then subsequently leads to a decrease in the energy contained in the density fluctuations. The nonlinear power transfer results also showed that the radial scale length of the turbulence was regulated by the GAM ZF.

In DIII-D, the measured radial correlation lengths of long-wavelength turbulence were compared with predictions made with analytical models and simulations [324]. The results showed that the correlation lengths lay between the \( \rho_S \) and \( \rho_{\theta S} \) (which corresponds to \( \rho_S \) evaluated with the poloidal magnetic field) scale lengths. The self-consistent ZFs could also be artificially quenched in the numerical simulations of these experiments, and the effect on radial correlation length could then be examined. The results (figure 31) showed that ZF dynamics had to be included to obtain reasonably good agreement with experiment. Without ZFs, the simulations predicted radial correlation lengths on the scale of the minor radius or gradient scale length, clearly at odds with the experimental observations. This work provides indirect evidence that ZFs indeed play an important role in the self-regulation of the scale length of the turbulence.

Thus these results provide evidence that the GAM branch of the ZF acts to regulate the amplitude and radial scale length of the turbulence; and thereby could play a role in regulating transport in the edge region of toroidally confined plasmas. The available results also provide indirect evidence that the turbulence indeed drives the GAM branch of the ZF. No such studies of zero-mean frequency ZF interactions with turbulence exist to our knowledge nor has a direct measure of ZF drive by turbulence in a strongly magnetized toroidal confinement device yet been reported. A summary of the current status of experimental studies of ZFs is available [325], as is a summary of the theoretical background [39, 46].

5.9. Turbulence during transition to improved confinement regimes

In the H-mode regime, first reported in ASDEX [83], a rapid spontaneous reduction in turbulent transport at the plasma edge is observed if sufficient plasma heating was applied and the device wall conditions were sufficiently clean (see [48] for a summary of early observations of the
H-mode regime). This regime has subsequently been observed in nearly all medium and large-scale tokamak devices, as well as in stellarator devices [74]. In subsequent work on DIII-D, it was found that a strong radial electric field developed in the edge region during the L–H transition, and that this radial electric field was associated with the development of a strong poloidal plasma flow in the region immediately inside the separatrix [50]. At nearly the same time, it was found that a biased emitting electron inserted into the region of the CCT tokamak could also trigger an H-mode confinement regime without any strong auxiliary heating [51]. These findings clearly showed that the development of a radial electric field and associated plasma flow had a strong effect on the transport in the edge region of the plasma. Similar results have been reported in other confinement experiments (see e.g. [326] for recent results from ASDEX-UG as an example). These observations have been explained using the flow shear decorrelation model described earlier in this review. A theoretical review of edge transport barriers is available [327].

Core transport barriers [89] have been observed to form in most types of confinement devices and are associated with the development of a strong $E \times B$ shear layer in the core plasma [71, 72]; this is sometimes accompanied by the formation of reversed magnetic shear region (for a review of these observations see e.g. [60]). A summary of the macroscopic profile evolution in several such experiments, and of the combined role of linear stabilization via magnetic shear, via the Shafranov shift of the flux surfaces and via nonlinear $E \times B$ shear decorrelation physics on the formation of these barriers can be found in the literature [328, 329]. Initially only reductions in the effective ion thermal diffusivity were reported; in later work reductions in the effective electron thermal diffusivity were also reported. We summarize representative turbulence measurements from several experiments here.

The plasma gradients undergo a substantial evolution after the L–H transition or ITB formation, and thus the underlying turbulence drive is significantly changed. This complicates the comparison of measurements taken before and after barrier formation. To avoid this complication, we examine the evolution of the turbulence across the L–H transition on timescales short compared with the profile evolution timescale, and then examine the turbulence
changes from L-mode state to the steady-state ELM-free H-mode state. We then use a similar approach to look at the turbulence evolution during and after ITB formation. We then summarize the existing tests of flow shear decorrelation mechanism thought to be responsible for the change in transport.

5.9.1. Turbulence amplitudes and particle transport during the L–H transition. Reflectometry measurements at the edge of DIII-D [49] showed that there was a rapid reduction in the density fluctuation signal in a region several centimeters wide, located just inside the LCFS. Measurements with a PCI diagnostic [145] (figure 32) on DIII-D confirmed the earlier reflectometry measurements, and showed that the radial correlation length underwent a significant decrease in the same region where the electric field developed; this decrease occurred simultaneously with the development of the strong electric field and the reduction in turbulence amplitude. Evidence for similar turbulence dynamics has been reported in a number of other tokamak devices and, recently, from the W7-AS stellarator device [330].

Probe measurements on the PBX-M tokamak provided the first measurement of the turbulent particle flux at the separatrix during a spontaneous L–H transition and showed that the particle flux was indeed quenched at the moment of the L–H transition [62]. Subsequent probe measurements in the DIII-D tokamak in discharges that operated very close to the L–H transition power threshold demonstrated that the poloidal plasma rotation velocity seems to play the critical role in triggering the transition [331]. This work would then suggest that the transition physics is determined by the physics that governs the poloidal rotation in the edge region.

![Figure 32. (a) RMS fluctuation amplitude, (b) radial correlation length and (c) Dα emission from the separatrix region during a DIII-D L–H transition; reprinted from [145] with permission from Elsevier.](image-url)
5.9.2. Turbulent Reynolds stress in the L–H transition. Theory suggests that the poloidal (and toroidal) plasma velocities are determined in part by the turbulent Reynolds stress, which acts to transport momentum in the plasma edge region [37]. Work on the HT-6M tokamak tested the hypothesis that the turbulent Reynolds stress acts to trigger a bifurcation in the poloidal plasma velocity, and thereby trigger the L–H transition [332]. In this experiment, a confinement transition was triggered by momentarily applying a large loop voltage to the plasma, presumably by rapidly pulsing the ohmic heating coil [333]. Several hundred microseconds after this transient is initiated a large increase in the poloidal plasma rotation and radial electric field was observed, while no significant changes in the toroidal rotation or ion pressure gradient were reported. The time and spatially resolved toroidal and poloidal main ion flow velocities were measured with Mach probes, while the ion pressure gradient was measured with a retarding field energy analyzer. A multi-tip probe array was then used to measure the temporal evolution of the turbulent Reynolds stress profile during an L–H transition. All measurements were reported with high time and space resolution, allowing an evaluation of the turbulent momentum conservation equation.

The results show the transient formation of a highly sheared Reynolds stress profile (see figure 33) just before and during the L–H transition. Analysis of this shear stress in the turbulent ion momentum equation showed that it led to the development of a strong poloidal plasma flow at the LCFS. The spatial and temporal evolution of the poloidal flow arising from the Reynolds stress was found to agree with the measured flow and the associated radial electric field during and immediately (<1 ms) after the L–H transition. The results were interpreted to support the hypothesis that the turbulent Reynolds stress could act to trigger the formation of a strongly sheared poloidal plasma flow and sheared radial electric field, which could then act to decorrelate the turbulence, resulting in the formation of an edge transport barrier formation.

These results would seem to provide strong support for the Reynolds stress trigger hypothesis; however, we note that another group has re-analyzed the HT-6A results [130] and concluded that the turbulent Reynolds stress was too small to induce strongly sheared flows in the HT-6A experiments. To date this disagreement has not been resolved, and to our knowledge there are no other published results of the evolution of the Reynolds stress in an
Figure 34. Left panels: relative density fluctuation time evolution in ERS discharge (left panel), and in RS discharge (right panel). In the RS discharge the fluctuation amplitude jumps between a state of higher and a state of lower density fluctuation amplitude; reproduced with permission from [334]. Copyright 1996 American Physical Society.

L–H transition. Thus it would appear that the crucial physics that triggers the formation of an H-mode is still not completely understood. This is an obvious area that should receive attention in future work.

5.9.3. Turbulence behavior during ITB formation When ITB formation occurred in TFTR discharges, Mazucatto et al [334] found that the density fluctuation amplitude was decreased in the region lying at and inside the reversed magnetic shear region, which is where the ITB is formed (see figure 34); similar discharges with reversed magnetic shear that did not show a sustained confinement improvement showed that the density fluctuations became bursty, but did not exhibit a sustained reduction in amplitude. Analysis of ITB discharges showed that the linear growth rate of the instabilities thought to be present in the plasma (both ITG and TEM) was reduced or even stabilized by the combined effect of reversed magnetic shear and $E \times B$ shear (which was inferred from an ion pressure balance analysis). Subsequent measurements and analysis showed that the formation of the ITB is preceded (and thus presumably caused) by the formation of a sheared $E \times B$ poloidal flow in the central plasma region [71, 72]. Thus this ITB formation exhibits many features that are reminiscent of the formation of the H-mode barrier at the plasma edge. The mechanism that triggers the formation of the sheared electric field and plasma flow was not addressed by these experiments.

In DIII-D, a central plasma transport barrier was combined with an edge plasma transport barrier to form the so-called quiescent double barrier (QDB) discharge [335, 336]. Measurements of the turbulence correlation length in the central plasma region that exhibits reduced thermal transport indicate that the turbulence scale length undergoes a significant reduction as compared with the L-mode case (figure 35). Turbulence simulations that took into account the mean shear flow from the experiment gave turbulent correlation lengths in reasonable agreement with the QDB experimental results. These changes were attributed to the shearing effect of a mean $E \times B$ flow.

Results from JT-60U using a similar correlation reflectometry technique give additional insight into the evolution of central plasma turbulence during the development of an ITB [337]. During the ramp up of the toroidal plasma current and neutral beam power in JT-60U, a rapid reduction in core plasma transport rates is observed to occur. A multi-frequency correlation reflectometry system was used to monitor the fluctuation amplitude and radial
turbulent correlation length during the formation and growth of this transport barrier. The radial correlation length exhibits a gradual (many tens of milliseconds) large (∼10 x) reduction as the core plasma develops into the ITB regime (see figure 36). This result is reminiscent of the radial correlation length in the L–H transition, but occurs over a much longer time period.
We note that similar studies and results have been published from JET [338, 339], ASDEX-UG [340] and T-10 [341] in both ion ITB and electron ITB conditions (which were defined by reductions in the respective thermal diffusion coefficients for the two heat transport channels). Furthermore, improved confinement regimes have also been demonstrated in stellarator devices and have been associated with the formation of sheared radial electric fields and plasma flows as well. For reviews of such stellarator results see e.g. [342, 343]. Taken together, these results show that the central plasma turbulence could undergo significant reductions in amplitude and radial scale length, and that these reductions are associated with reductions in core plasma thermal and particle transport. The radial electric field and associated $E \times B$ shearing rate appear to play an important role in both regulating the turbulence amplitude and scale length, and in triggering the transition between different confinement regimes. However, there are also important linear stabilization mechanisms involving magnetic shear and flux surface compression from the Shafronov shift that also appear to be important.

The origins of the electric field in the central region of these tokamak experiments were considered in work on JET [338]. The initial formation of the radial electric field was attributed to a large toroidal plasma rotation driven by intense core neutral beam momentum deposition. The resulting $E \times B$ shear was then calculated to be sufficient to begin to decorrelate the turbulence, leading to a reduction in the turbulent cross-field transport. In JET the resulting increased ion pressure gradient was calculated to be capable of sustaining the mean shear flow at a rate sufficient to continue the shear decorrelation suppression of the turbulence in a manner that is nearly identical to that posited to occur in the H-mode edge barrier [60, 73]. We note that in stellarator devices, similar rapid changes in core radial electric field have been reported and have been found to be consistent with neoclassical effects alone [344, 345]. It is not clear whether similar processes could explain core plasma radial electric field formation in tokamak.
plasmas. It is also intriguing to note that recent work in DIII-D suggests a relationship between ZF evolution, local minimum \( q \) and the trigger of an ITB [346, 347]. Additional work is clearly needed to understand the mechanism(s) that govern the formation of the core plasma radial electric field and associated sheared \( E \times B \) flow.

5.10. Turbulence, transport and shear decorrelation in established shear layers

Probe measurements have been made in the steady-state phase of spontaneously formed H-modes in the PBX-M and DIII-D tokamak experiments and provide simultaneous localized measurements of ion saturation current density and floating potential fluctuations in the shear layer region and in the SOL [61, 62]. The results show that, after the plasma pressure reaches the new equilibrium in ELM-free H-mode, the fluctuation response has a complex behavior and cannot simply be considered to have undergone a reduction in amplitude. Across most of the plasma edge and SOL region, the relative fluctuation amplitude is indeed decreased relative to the values seen in L-mode. However, in the region of maximum \( E \times B \) shear the floating potential fluctuation amplitude is actually increased. However, the inferred particle flux made from the correlation of density and potential fluctuations showed very large decreases across the whole edge and SOL region. This decrease occurs across the broad frequency range (figure 37(a)). As discussed above, the particle flux depends not only on the turbulence amplitude, but also on the degree of coherence between the two fluctuating fields, and on the cross-phase between the two fields. The results from DIII-D showed very significant changes in the cross-phase (figure 37(b)), which result in a large decrease in the time-averaged particle flux at the LCFS of the plasma. These results clearly showed that the original one-field model of shear flow decorrelation needed to be supplemented by a theory that included the cross-phase. Work in TEXTOR biased H-mode experiments also showed that the change in the cross-phase was important [55].

Probe measurements subsequently allowed the simultaneous measurement of density, electron temperature and potential fluctuations in the edge of DIII-D and TEXTOR [104, 106]. The results clearly showed a reduction in the cross-field heat flux carried by the edge turbulence during ELM-free H-modes—in fact the relative decrease is much larger than the corresponding increase in thermal energy confinement. The development of an ELM would then momentarily increase the cross-field heat flux back to values that were close or equal to the values found in L-mode [78].

Results from several experiments show a slightly different picture for the evolution of the multi-field turbulence statistics during H-mode. In particular, work in the PBX-M spontaneous H-mode [62, 108], and in both the CCT and CASTOR biased H-mode experiments [53, 122] shows that the density and potential fluctuations become decorrelated in the shear layer during H-mode. These latter results also provided turbulence measurements at multiple poloidal locations during biased H-mode experiments. The results showed that the response of the edge turbulent at the shear layer depended significantly on the poloidal position. There were significant changes in fluctuation amplitude reported on the low-field side of the tokamak. The large decrease in turbulent particle flux during the biased H-mode was attributable primarily to a large decrease in the zero-time delay cross-correlation between density and potential. These multi-field results are still not explained by a single field shear decorrelation model that was extended to include toroidal effects [348].

5.10.1. Tests of the shear decorrelation model. In the PBX-M tokamak, probes were used to compare the turbulence decorrelation rate in L-mode that is operating just below the L–H transition power threshold with the shear decorrelation rate in an H-mode that was just above
the L–H power threshold [62]. The relative change in other turbulence statistics across the transition was also studied. The results showed that a modest change in the ratio of shearing rate to turbulence decorrelation rate lead to a very large change in the turbulence statistics and in the associated cross-field particle flux. This result suggests that either the shearing rate was not the key parameter involved, or that the transition was very sensitive to the shearing rate, much like a phase transition is highly sensitive to the critical parameter at the boundary between two different phases. Work in W7-AS also concluded that the response of the turbulence to the $E \times B$ during the L–H transition exhibited signatures of a phase transition in that very small changes in the critical parameter ($E \times B$ shear) led to drastic changes in the turbulence [330].

The biased H-mode experiments in the CCT and TEXTOR tokamaks [51, 349–351] permitted a controlled study of the effect of a mean shear flow upon the edge turbulence. Detailed studies of this nature have been reported [56], and are summarized in figures 38 and 39. The cross-phase between ion saturation current and floating potential fluctuation showed a gradual change with velocity shear, resulting in a gradual reduction in the cross-field particle flux carried by the turbulence. The variation of the cross-phase and associated reduction in particle flux due to that cross-phase variation as reported in this paper gave reasonable agreement with a theoretical model for the variation of cross-phase with mean flow shear [352]. The resulting reduction in particle flux yields a substantial reduction in the effective cross-field diffusion coefficient, defined as the global particle flux averaged over the LCFS, divided by the plasma density gradient measured at the LCFS [353].
There are also studies of the response of turbulence to the gradual growth of the H-mode $E \times B$ shear layer during the development of the so-called VH-mode in DIII-D. Space restrictions do not allow a full discussion of these results; however, the published literature show that the gradual increased spatial extent of the edge transport barrier coincides with the increased extent of the $E \times B$ shear layer and that the shear decorrelation mechanism gradually begins to operate over a wider range of the edge plasma region [70]. It is not clear what governs the spread of this shear flow region, although it would seem that cross-field transport of momentum must play a role. Likewise, there are reports of reductions of turbulence and transport during the formation of a poloidally and toroidally symmetric radiating mantle caused by the controlled injection of neon or argon noble gases at the edge of tokamak devices [354]. The interested reader should refer to the literature for a discussion of these observations.

### 5.11. Evidence for long-range transport mechanisms

All of the previous results view the turbulence and transport as occurring primarily on a flux surface, with only small excursions away from that surface due to turbulent $E \times B$
drifts which quickly decorrelate. The net result is a transport process that occurs via a large number of turbulent diffusive steps from flux surfaces that are otherwise isolated from each other. However, there are a number of proposed mechanisms which could cause transport to occur over longer ranges via non-diffusive mechanisms. Avalanche-type processes have been hypothesized in which a marginally stable critical gradient is exceeded, leading to the formation of dual propagating fronts—a positive excursion of particles or heat which propagates down the background gradient, and a negative going excursion that propagates up the gradient much like is found in, e.g., sandpiles and avalanches [83]. Such processes could occur somewhat uniformly over a flux surface, or could be isolated to a narrow poloidal and toroidal location, leading to the formation of radially elongated streamer-like disturbances [81, 87, 355, 356] in the density and temperature which could propagate over long spatial scales. It is expected that such structures would be regulated by the existence of ZFs and mean shear layers in the plasma, which could act to quench the formation and propagation of such streamers via the usual shear decorrelation mechanism. Such structures have been hypothesized to lead to an enhancement of otherwise small levels of thermal transport from sub-ion gyroradius scaled ETG turbulence in numerical simulations [81]. Finally, it has been proposed that turbulence could propagate, or spread, from unstable regions into stable regions [45, 357, 358] via either toroidal mode coupling, nonlinear propagation of turbulence or by the self-trapping of drift turbulence in ZF structures which then propagate away from their point of origin carrying the trapped turbulent structures with them [359]. All of these mechanisms could lead to the loss of gyro-Bohm scaling for drift turbulence; indeed it is this property that has in part motivated the development of these theoretical constructs. The experimental study of such mechanisms is not well developed; however, there are a few publications that bear on the existence of such dynamics which are summarized here.

The examination of low-frequency electron temperature fluctuations obtained from ECE diagnostic on DIII-D [360] showed evidence for radially propagating structures. From the observed range of timescales for these perturbations, the temperature fluctuations should have thermalized within the flux surface and thus it was argued that this temperature disturbances represent large-scale fluctuations that occur on the entire flux surface. The timescale for density fluctuations to relax within the flux surface was estimated to be much longer; thus the reflectometry measurements of a Hurst parameter $H > 1/2$ over a significant fraction of the discharge volume [218] may indicate spatially localized disturbances which have not yet had time to equilibrate within the flux surface. These events were estimated to be capable of carrying a significant fraction, if not most, of the electron heat for the discharges studied in the papers [360]. Thus these processes may represent an important transport mechanism which could also lead to a breaking of gyro-Bohm transport scaling in the core plasma region. We point out, however, that to our knowledge no other tokamak group has reported reproducing these results.

We also note that work on JET showed that the core plasma turbulence amplitude is rapidly (~milliseconds) reduced in response to the L–H transition that occurs in the edge region, and that the evolution of the temperature gradients across the entire plasma was consistent with a model where the power balance thermal diffusivity undergoes a similarly rapid reduction over the outer half of the plasma column at the moment of the L–H transition, even though the velocity shear layer is localized to the region immediately inside the LCFS [362]. These results were interpreted to imply a clear coupling between edge turbulence and core turbulence, and motivated an initial theoretical study of this problem [358]; however, at this writing turbulent transport is usually considered to be a locally determined process, which is inconsistent with these JET observations. In this vein, we note with interest recent BES density fluctuation frequency spectra from DIII-D taken at $r/a = 0.6$ during an L–H transition [363] showed...
that the turbulence well inside the LCFS undergoes a significant decrease in amplitude within \( \leq 10 \text{ ms} \) of the L–H transition. We also note that no core simulation has successfully reproduced the rapid increase in turbulence amplitude that occurs as the minor radius approaches the LCFS. These observations would appear to suggest that large-amplitude turbulence at the LCFS and SOL is coupled to the deeper core plasma turbulence, and that the nearly instantaneous reduction of this edge turbulence at the moment of the L–H transition then effects a change in transport over a region that extends well inside the velocity shear layer. However, further experimental study supported by theory and simulation are needed to confirm this picture and determine the mechanism(s) that could be responsible for such coupling. If confirmed, this would represent a significant departure from the usual picture of transport being determined by local conditions alone, and would also imply that core plasma simulations that do not simultaneously incorporate an edge and SOL turbulence model would be deficient.

6. Discussion

Earlier in this review we framed a number of questions which could be addressed using published results. Here we recap these questions, and provide a discussion of the current state of understanding of these issues based on the available literature.

How do coherent drift waves evolve into drift turbulence? Although we simply referred readers to the literature, there are a number of published studies that show good agreement between linear drift wave theory and measurements of collisional and collisionless electron pressure driven drift waves, and ITG-driven drift waves; there are no such measurements for trapped electron modes or for ETG modes. An examination of the literature also showed convincing evidence of the development of the parametric decay of a large-amplitude unstable collisional drift wave into two nonlinearly driven drift wave eigenmodes. Two separate experiments in linear plasma devices using two different control parameters (parallel current in the first and a combination of radial pressure gradient and magnetic field in the second) also show the development of a broad turbulent drift wave spectrum from coherent waves. In both cases, the plasma first develops several harmonics which transition to a mode-locked behavior, leading to the collapse of the multi-peaked spectra into a spectrum dominated by a finite width large-amplitude quasi-coherent mode. A further increase in the critical control parameter then leads to a breakup of this mode and the formation of a broadened spectrum that characterized drift turbulence. In one of these papers, this sequence was shown to be consistent with the Ruelle–Takens route to turbulence that has been observed in fluid experiments. Although the techniques have now been developed, there have been no studies of the development of nonlinear energy transfer processes during such transition to turbulence experiments; such measurements would provide a valuable test of our understanding of the balance between linear stability, nonlinear energy transfer, nonlinear emergence of ordered structures and development of a stationary turbulence spectrum.

What evidence exists that links turbulence with the actual plasma confinement in closed flux surface geometries? In confinement devices with closed magnetic field topology, the effective thermal diffusivity deduced from a steady-state power balance is usually proportional to the turbulence amplitude (i.e. in spatial locations or operating conditions with relatively higher thermal transport the turbulence amplitudes are larger). Furthermore, probe measurement at the edge are capable of inferring the turbulent particle flux and heat flux and show that, for the limited set of conditions under which such measurements are available, the inferred
fluxes integrated over the LCFS would yield a confinement time that is within factors of 2–3 of the observed global energy confinement time. These results are consistent with turbulence causing the global confinement. However, such measurements have only been carried out under a limited set of conditions; it would clearly be desirable to extend such measurements over a wide range of conditions (e.g. approaching the density limit, high beta, L-mode and H-mode). It would also seem useful to increase the precision of the measurements to determine if there is a missing channel for transport that is not accounted for by the drift turbulence.

Transient transport experiments clearly show that thermal diffusivities exhibit a critical gradient behavior. This should have clear manifestations in the underlying turbulence behavior (e.g. turbulence amplitudes and/or radial correlation lengths and ZF characteristics) when the critical gradient is momentarily exceeded. It would be useful to combine time-resolved non-stationary studies of turbulence and ZF response during such transient transport experiments to determine the links between the microscopic turbulence behavior and macroscopic system response.

What are the dominant free energy sources and dissipation mechanisms that operate on ion and electrons in core, edge and SOL regions, and how do these vary with operating conditions? The origin of turbulence and transport in the ohmically heated collisional plasmas dominated by electron thermal transport would appear to still be poorly understood, much as was the case in an earlier review by Wootton et al [22]. There is evidence that in strongly heated tokamaks the ITG instability could be present and drive thermal transport at higher collisionality. In particular, the spatial scales of the turbulence and the propagation direction are consistent with this interpretation of the data. Linear stability analysis also showed that ITG modes are unstable in such plasmas. There is also evidence that TEM turbulence could occur at low values of collisionality and that the ITG instability then appears as collisionality is increased in tokamaks with strong auxiliary heating. In addition, there are some initial reports of small-scaled ETG turbulence co-existing with larger scaled ITG and TEM driven turbulence in tokamak experiments. The possible contribution of these ETG fluctuations to thermal transport is not known. No such identifications have been attempted in non-tokamak systems.

At the LCFS and in the SOL, the observations are consistent with turbulence driven by a combination of drift wave and interchange instability, triggered by the transition to open field lines that change the fundamental physics governing the cross-phase between density and potential fluctuations due to the introduction of sheath dissipation of parallel electron motion. The results also show that this LCFS/SOL turbulence appears to disturb the plasma up to several \( \rho_s \) inside the LCFS and thus perhaps couple to the turbulence found in the closed flux surface region. Results from JET L–H transitions provide additional evidence that the turbulence at the LCFS does affect core turbulence. However, there has been relatively little attention given to the coupling of turbulence across this boundary. Given the fact that the L–H transition is triggered precisely at this location, and that the overall plasma performance is sensitive to the transport rates in this region, it would be extremely useful to develop a much better and deeper understanding of how LCFS/SOL turbulence could influence and couple with turbulence on closed field lines, and to understand the origins and distribution of plasma flows in this region.

What is the evidence for the existence of ZFs and GAMs in confined plasmas? What role do ZFs and GAMs play in setting the turbulence amplitudes and associated transport rates? Do experiments shed any light on the mechanisms by which ZFs and GAMs might saturate? There is strong evidence from laboratory devices, from small confinement devices and from a number of tokamaks and stellarator devices that radially sheared ZFs and GAMs exist and that
these flows play a role in regulating the turbulence amplitude and cross-field transport on closed field lines. These flows have the expected poloidal and toroidal spatial structure, and have a shearing rate that is strong enough to influence, but not completely dominate, the turbulence dynamics. In laboratory devices and small confinement devices the direct drive of the ZFs by turbulence has been demonstrated, while in the larger confinement experiments measurements show that the turbulence and the shear flows are nonlinearly coupled as seen, e.g., in the finite bicoherence and bispectrum computed between these multi-scaled fluctuations. Measurements in these larger experiments show that the turbulence fluctuation amplitude rises and falls out of phase with the energy contained in the ZFs and GAMs. This observation, coupled with the nonlinear coupling between the turbulence and shear flows in confinement devices strongly suggest, but do not prove, that the shear flows are driven by the turbulence and that the shear flows in turn regulate the turbulence amplitudes and transport rates. Results from tokamaks show that the GAMs in the edge region could regulate the cross-field flux and radial scale length of the turbulence. There is no existing study of the saturation physics of the ZFs and GAMs, nor have any experimental studies been performed to identify and study the damping mechanisms that operate on these flows. In addition, there are no studies of the response of the turbulence and ZFs during transient transport experiments. Thus, although it is now clear that drift turbulence is crucially linked to ZF physics, additional work aimed at linking this interaction to macroscopic transport and confinement observations, and experiments aimed at identifying ZF damping mechanisms is needed.

Are the turbulent correlation lengths determined by local gradients, which will lead to the so-called gyro-Bohm transport scaling, or are they determined by the system size, which would lead to the so-called Bohm transport scaling, and is there any hint in the turbulence data for a transition from one regime to the other? This important question bears on the scaling of turbulent transport to ITER; however, there are only three published studies on how the turbulence scale length varies with the system size. The tokamak studies showed that, over a limited range of dimensionless system size, the turbulence correlation length depended upon the local plasma gradients in a picture that is consistent with expectations for gyro-Bohm scaling. However, these studies have only been performed over a very limited range (factor of ∼ few) of dimensionless system size. Experiments in a torsatron device showed that the turbulence correlation length scaling transitions from a gyro-Bohm scaling to a Bohm-like scaling when $\rho_s > 0.1$.

Transient transport experiments show strong evidence for critical gradient transport phenomena. However, there are no existing studies on how turbulence scale sizes vary as the critical gradient is exceeded, nor are there simultaneous perturbative transport/turbulence correlation length scaling studies. It seems likely that the turbulence intensity, correlation length and ZF magnitude must respond to how hard the system is driven beyond the critical value. An experimental study of this would be useful, and would also provide valuable insight into the system size scaling of transport physics—an issue of significant interest for ITER.

Does the turbulence play a role in triggering the formation of transport barriers at the edge and in the core regions of the plasma? This question bears on the important issue of the threshold for triggering a transition to an improved confinement regime. Our review of the literature shows only two papers which have tested the hypothesis that the trigger for the L–H transition is caused by the edge turbulence. These two papers provide contradictory conclusions, and thus the question of the role of turbulence in triggering the L–H transition must be considered open. There are no similar data available in the formation of ITBs; and
thus a similar conclusion holds for that region of the plasma as well. This is clearly an area requiring further experimental effort. It is also known empirically that the wall conditions of the experimental device play an important role in determining whether the H-mode regime could be reached, as well as the quality of the confinement in H-mode. There have been no studies of how wall conditioning affects edge turbulence and shear flows. Thus a better understanding of how the edge turbulence and shear flows are influenced by and respond to the wall conditions, and of how they will behave in walls which have reached equilibrium with the plasma are needed. Such work is of particular importance for steady-state high performance burning plasmas anticipated to occur in advanced tokamak scenarios.

**What detailed turbulence measurements are available in established mean shear flows, and is the response of the turbulence to the shear flow consistent with theory?**

Detailed probe measurements of turbulence in mean shear flows located just inside the LCFS show reasonable agreement between the experimentally measured variation of the relative density fluctuation amplitude and cross-field particle flux with shearing rate and theoretical predictions. The data are consistent with the existence of a critical value of shearing rate, roughly equivalent to the natural shear decorrelation rate present in the plasma in the absence of strong flow shear, required to quench most of the cross-field transport. The cross-phase between density and potential fluctuations also exhibits important changes with increased mean shearing rates, and plays an important role in reducing the inferred transport. These results show that the turbulence could respond gradually and continuously to flow shear rates below a critical value but then, when the shearing rates are increased further, a large change in the turbulence and associated transport occurs.

Such detailed multi-field measurements are not available in the core plasma region during ITB formation. However, there are measurements showing a reduction in turbulence amplitude and radial correlation length during the formation of ion channel ITBs. In addition, it is known that these barriers are associated with the formation of core plasma sheared electric fields that are of a magnitude to be consistent with the shear flow decorrelation hypothesis.

**Is there evidence that the turbulence could develop a long-range character by triggering avalanche or streamer dynamics?**

This question also bears on the scaling of turbulence with system size. We summarized work from one experiment that showed evidence for the long-range propagation of electron temperature perturbations over a significant fraction of the plasma minor radius. In that work it was argued, but not directly measured, that these perturbations represented disturbances that were uniform within the magnetic flux surface, and thus represented avalanche-like phenomena propagating across flux surfaces and it was argued that these events could have a significant impact on the power balance. Thus, although they may or may not be turbulent phenomena, such phenomena may be important for confinement. However, we also note that independent confirmation of this result is needed. In separate work, the Hurst parameter approach was used to show evidence for long-range correlations in the core plasma turbulence of L-mode discharges. However, this approach is indirect (i.e. it does not directly measure such structures but rather infers their existence statistically). There is also limited evidence from JET L–H transitions that a rapid reduction of turbulence at the plasma edge is quickly communicated to a large fraction of the plasma, resulting in a rapid reduction of the effective thermal transport over much of the plasma. These observations suggest that the edge plasma could somehow influence the deeper core plasma turbulence, although the mechanism for this coupling (if it indeed exists) is unknown at present.
These observations are difficult to reconcile with the usual view of turbulence being determined by local conditions only, and could lead to the breaking of gyro-Bohm thermal transport scaling. They clearly point out that studies of non-local transport mechanisms should be carried out to more clearly understand the existence and significance of long-range plasma transport processes in confined plasmas.

7. Open issues and suggestions for future work

The questions that were defined at the beginning of this review and that were summarized above are constrained by the available published turbulence experimental results. However, there are also a number of questions related to turbulence and transport for which there are no substantive experimental studies of turbulence available. Thus, in addition to pursuing the suggestions above, it would be useful to devise experimental studies of turbulence and transport in response to the following questions and topics.

- Is there evidence for turbulence arising from other (i.e. non-drift wave) instabilities in confined fusion plasmas?
- Why is wall conditioning important to the formation of the H-mode? Why does reduced recycling provide a significant increase in confinement across much of the plasma profile? Could good confinement be maintained when walls get saturated and recycling returns to equilibrium?
- Is the generation of intermittent edge and SOL transport linked to the density limit and, if so, what mechanism links the phenomena?
- Does turbulent momentum transport play a role in spontaneous toroidal plasma rotation? How are the edge plasma flows linked to plasma shape, divertor X-point location and ion grad-$B$ drift directions?
- What triggers the L–H transition, and how is this physics linked to the macroscopic parameters needed for the transition?
- Does turbulence recover during pedestal formation and if so, does this play a role in the saturation of the pedestal pressure gradient?
- When the critical gradient for stability is exceeded, does the coupled turbulence/ZF system respond as expected? What happens to the size and scaling of the turbulent correlation length?
- Do electromagnetic drift wave effects at finite beta cause significant transport in confined fusion plasmas? Do such effects link small-scaled drift instabilities with larger scaled Alfvén wave activity that could arise, e.g., in burning plasmas?
- Can turbulent fluctuations explain the significant difference in both magnitude and parametric dependence of electron and ion thermal fluxes? What causes electron heat transport when flow shear eliminates ion-scale turbulent transport?
- Why does the plasma density profile respond to application of stochastic edge fields, while the electron temperature profile does not change substantially in recent experiments?
- Why does confinement depend on the ion mass? Is there any evidence that the nonlinear drift wave/damped IAW mechanism proposed from laboratory experiments occurs in confinement devices?
- Does turbulence spread from unstable regions into linearly stable regions?

We conclude that, although major conceptual breakthroughs in our understanding of drift turbulence in confined plasmas have recently occurred, and that the experimental data appear to be consistent with both theoretical expectations for the coupled drift turbulence/ZF system and with the effect of strong mean shear flows on turbulence, it is still difficult to use these
observations to then quantitatively explain the macroscopic behavior of the plasma. Future studies aimed at elucidating such links would clearly be useful, and could help provide the basis to form physics-based predictions of turbulence, transport and confinement in future fusion devices.

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