1. Introduction

There is clear experimental evidence that scrape-off layer (SOL) signals from Langmuir probes in tokamak devices may be adequately modelled by assuming that filament transport follows a shot-noise process [1–6]. Characteristics may be used to extrapolate SOL-widths and heat-loads on the confining vessel wall and are of significant importance for the progress of the fusion program. While understanding of particle- and heat transport by filaments has made rapid progress [7–12] (for an excellent review see [13]), there still lacks explanation on how the resulting confinement properties are influenced by the fuel mass—an important issue as fusion reactors such as ITER are intended to run on deuterium–tritium plasma mixtures, compared to most laboratory experiments, employing protium or deuterium fuel. Contrary to what is to be expected from gyro-Bohm classical and neoclassical transport estimates, experiments show that heavier hydrogen isotopes produce enhanced plasma confinement in tokamaks [14–17]. Recently, improved confinement for increased ion mass has also been observed in a reversed field-pinch configuration [18]. It has been suggested that enhanced shear-flow activity for heavier hydrogen isotopes may suppress edge turbulence favourably. Gyrokinetic theory [19] and simulations show that the ion mass may result in enhanced residual zonal-flow levels [20–22]. Gyrofluid simulations of edge turbulence have found improved confinement for heavier hydrogen mixtures with only mild dependence of zonal flow activity on ion mass [23]. In turbulence simulations of edge conditions, self-consistent development of zonal structures together with consistent boundary conditions along the magnetic field lines is necessary to capture the birth of excess density filaments (close to the separatrix) and its influence on mediating the radial propagation of these filaments, which may account for a significant part of particle losses observed in the edge region [24–27]. Seeded filament motion has been shown to depend crucially on the ion mass, with heavy ions moving slower [28]. It is well worth to study the dynamical interplay between edge/SOL turbulence and its coupling to filament propagation. In this contribution we compare the shot-noise model [1] to time-series obtained from coupled edge/SOL gyrofluid turbulence simulations carried deeply into the nonlinearly saturated state and investigate how characteristic statistical blob-parameters are influenced by the ion mass.
2. Gyrofluid model and computation

In order to compare time series from numerical simulations with the shot-noise model, it is crucial to sample enough data points for statistics to be meaningful. To this end, we chose the computationally efficient gyrofluid electromagnetic model introduced by Scott [29]. Here we employ the isothermal approximation, which is justified for L-mode outer edge and SOL plasmas. For SOL simulations, a full-F gyrofluid approach would be favourable, but in our present implementation the computational costs for long time series in well-resolved edge/SOL coupling simulation restrict us here to the delta-f limit. In this limit the model consists of evolution equations for the gyrocenter densities $n_i$ and parallel velocities $u_{\parallel i}$ of electrons and ions, where the index $s$ denotes the species with $s \in (e,i)$:

$$\frac{d}{dt} n_s = -\nabla_u u_{\parallel i} + K (\phi_s + \tau_s n_s),$$

where we have introduced the Poisson bracket as

$$\{ f, g \} = \left( \frac{\partial f}{\partial x} \frac{\partial g}{\partial y} - \frac{\partial f}{\partial y} \frac{\partial g}{\partial x} \right).$$

The plasma beta parameter

$$\tilde{\beta} = 4 \pi \rho_e \left( \frac{qR}{L_\perp} \right)^2,$$

controls the shear-Alfvén activity, and

$$C = 0.51 \frac{\nu_e L_\perp}{\epsilon_0} \frac{m_e}{m_0} \left( \frac{qR}{L_\perp} \right)^2,$$

mediates the collisional parallel electron response for $Z = 1$ charged hydrogen isotopes.

The gyrofluid moments are coupled by the polarisation equation

$$\sum_s a_s \left[ \Gamma_1 n_s + \frac{\Gamma_0 - 1}{\tau_s} \phi_s \right] = 0,$$

and Ampère’s law

$$-\nabla_\perp^2 A_s = J_s = \sum_s a_s u_{\parallel i}.$$

The gyroscreened electrostatic potential is given by

$$\phi_s = \Gamma_1 \left( \rho_s^2 k^2 \right) \hat{\phi}_k,$$

where $\hat{\phi}_k$ are the Fourier coefficients of the electrostatic potential. Note that due to the smallness of the electron gyroradius, gyroscreening is only applied to the ions. The gyroaverage operators $\Gamma_0(b)$ and $\Gamma_1(b) = 1^{1/2}(b)$ correspond to multiplication of Fourier coefficients by $I_0(b) e^{-b}$ and $I_0(b/2) e^{-b/2}$, respectively, where $I_0$ is the modified Bessel function of zeroth order and we have introduced the shorthand notation $b = \rho_s^2 k^2$. We here use approximate Padé forms with $\Gamma_0(b) \approx (1 + b)^{-1}$ and $\Gamma_1(b) \approx (1 + b/2)^{-1}$ [30].

The perpendicular $E \times B$ advective and the parallel derivative operators for species $s$ are given by

$$\nabla_\perp = \left[ \frac{\partial}{\partial x} \frac{\partial}{\partial y} - \frac{\partial}{\partial y} \frac{\partial}{\partial x} \right].$$

Flux surface shaping effects [31, 32] in more general tokamak or stellarator geometry on SOL filaments [33] are here neglected for simplicity.

Spatial scales in each drift plane are normalised by the drift scale $\rho_0 = \sqrt{\tau_e m_0/e B}$, where $T_e$ is a reference electron temperature, $B$ is the reference magnetic field strength and $m_0$ is a reference ion mass, for which we use the mass of deuterium $m_0 = m_d$. The parallel coordinate is normalised by the parallel connection length, $L_\parallel = 2 \pi q R$, where $q$ is the safety factor at a reference location inside the separatrix and $z \in [-\pi, +\pi]$, where $z = 0$ denotes the outboard midplane. The influence of the connection length on turbulence properties across the separatrix is studied in [35]. The temporal

\[ z = \pi / 2 \]

\[ z = -\pi / 2 \]

\[ z = 0 \]

\[ \nabla_\perp = \left[ \frac{\partial^2}{\partial x^2} + \frac{\partial^2}{\partial y^2} \right]. \]
scale is set by $c_0/L_\perp$, where $c_0 = \sqrt{T_e/m_0}$, and $L_\perp$ is a perpendicular normalisation length (e.g. a generalised profile gradient scale length), so that $\delta = \rho_i/L_\perp$ is the dimensionless drift scale. The temporal scale may be expressed alternatively $L_\perp/c_0 = L_\perp/\rho_i\Omega_{Te} = (\Omega_{Te})^{-1}$, with the ion-cyclotron frequency $\Omega_{Te} = c_0/\rho_i$. In the following we employ $\delta = 0.015$ such that one normalised time unit corresponds to $(\Omega_{Te})^{-1} \sim 10^{-4}$ s. Fluctuation amplitudes are normalised according to $n_i/n_0 \to n_i, \phi/T_e \to \phi, A_{n_i}/L_\perp/\delta \beta_i R \to A_{n_i}, J_{n_i}/\epsilon_0 c_0 q_R R \to \epsilon_0, n_i L_\perp/c_0 q_R \to u_0$, with the electron beta $\beta_e = 4\pi \rho_e / B^2$ in terms of the background electron pressure $p_e = n_0 T_e$. Note that this normalisation produces the factor $\delta^{-1}$ with the Poisson brackets.

The main species dependent parameters are

$$a_s = \frac{Z_s n_s}{n_e}, \quad T_s = \frac{T_e}{Z_s}, \quad \mu_s = \frac{m_s}{Z_s m_0},$$
$$\rho_\perp = \mu_s T_s, \quad \epsilon_s = \mu_s \left( \frac{q R}{L_\perp} \right)^2,$$

setting the relative concentrations, temperatures, mass ratios and FLR scales of the respective species. $Z_s$ is the charge state of the species $s$ with mass $m_s$ and temperature $T_s$. Note that the index ‘s’ denotes both electrons and ions, while the index ‘i’ represents ion species such as proton, deuteron, tritium or helium.

We employ $\beta = 1$, $\omega_B = 0.04$, $\delta = 0.015$, $(qR/L_\perp)^2 = 27,000$ and magnetic shear $\delta = 1$, corresponding approximately to $L_\perp = 4.25$ cm, $R = 165$ cm, $B = 2$ T, $T_e = 70$ eV, $q = 3$ and $L_n = 31$ m. Typical for ASDEX Upgrade conditions close to the separatrix. All background parameters are constant across the $xy$ domain. Similar parameters for this numerical set-up have been employed in [36]. The collisionality parameter is set to $C = 50$, since this is known to increase radial blob velocity [28]. Lower collisionality requires longer simulation times since fewer (and slower) blobs are produced. Comparing the structures of the density fields at the outboard midplane ($z = 0$) with the shear region ($z = -\pi$) in the simulations, suggests that the outboard midplane dynamics is detached due to high collisionality, as reviewed in [13]. A similar trend with respect to collisionality has been found in [26, 27]. The ion temperature is fixed at $T_i = 1$. Each run is initialised with the following electron density profile

$$n_0(x) = n_{sep} + (n_{core} - n_{sep}) \left( 1 - \frac{x}{x_{sep}} \right), \quad x \leq x_{sep}$$
$$n_0(x) = n_{sol} + (n_{sep} - n_{sol}) \left( 1 - \frac{x - x_{sep}}{x_{sol}} \right), \quad x > x_{sep}$$

and the evolution of the profile together with the perturbations is computed at each time step (i.e. the background gradient is free to evolve and the simulations are flux driven). $n_{sep} = 0.5$ is the density at the separatrix location $x_{sep}$, $n_{core} = 1.5$ denotes the density at the innermost boundary (at $x = 0$) and $n_{sol} = 0.2$ is the density in the far-SOL, at the outermost radial boundary at $x = x_{sol}$. The width of the radial domain is then given by $x_{sep} + 4x_{sol}$. The ion gyrocenter density (the index $i$ denotes the ion species at hand) is initialised according to

$$n_i = \Gamma_i^{-1} n_0 = \left( 1 - \frac{1}{2} \tau_i \mu_i \nabla^2 \right) n_0,$$

starting each simulation with zero vorticity$^3$.

### 2.1 Parallel boundary conditions

We distinguish between two settings for parallel boundary conditions in 3-d simulations. In the case of tokamak edge region simulations a toroidal closed-flux-surface (CFS) geometry is considered, and boundary conditions on both state-variables $n_e, \phi$ and flux variables $\nu_{ee}, u_{ss}$ read

$$f(x, y, z, z^\perp) = f(x, y, z)$$

where $N$ is the number of drift-planes and $k \in 1, ..., N$ labels each locally orthogonal coordinate system on drift-plane $k$. The perpendicular coordinates $y_k$ are shear-shifted according to [38].

Including the SOL domain, parallel boundary conditions change across the separatrix. The state variables assume zero-gradient Neumann (sheath) boundary conditions at the limiter location and the flux variables are given as

$$u_{ee}|_{\pm \pi} = p_e|_{\pm \pi} = \mp \Gamma_{ee} n_e|_{\pm \pi},$$
$$\nu_{ee} = u_{ee}|_{\pm \pi} - J_i|_{\pm \pi} = \pm \Gamma_{ee} [(\Lambda + 1) n_i|_{\pm \pi} - \phi|_{\pm \pi}],$$

at the parallel boundaries $z = \pm \pi$ respectively [35]. The Debye current $J_D \sim e_n c_0 e$ is here linearised like in [34, 35], where background ($\phi$) and fluctuating ($\nu_{ee}$) quantities are separated. The sound speed then reads $c_s^2 = (T_e + T_i)/m_i$. The ions contribute to the fluctuating part of the Debye current by $\Gamma_{ii} c_0 \Gamma_{ii} e_n c_0$, and the electron pressure $p_e = n_e T_e$ is replaced by $p_e = n_i$ [35]. This effectively leads to the same dispersion relation for the Debye sheath instability as considering vorticity and (electron) temperature dynamics. Considering a six-moment gyrofluid model, one finds strong spatiotemporal correlation among electron temperature and electron density fluctuations [39], justifying the replacement in the (isothermal) linear treatment of the parallel boundary conditions. The present edge/SOL set-up and its effects on drift wave turbulence has been presented in detail by Ribeiro et al in [35, 40].

The sheath coupling constant is $\Gamma_{id} = \sqrt{1 + \tau_i}/(\mu_i \epsilon)$, with $1/\sqrt{\epsilon} = 2\pi L_\perp / L_i$, resulting from the normalisation scheme

$^3$The vorticity in this gyrofluid model reads $\Omega = n_i - \Gamma_i n_i$ and the gyroaverage and gyroscreening operators $\Gamma_0$ and $\Gamma_1$ are Hermitian.
(hence there is no collisionality dependence in the sheath constant) and the floating potential is given by \( \Lambda = \Lambda_0 + \Lambda_i \), where \( \Lambda_0 = \log \sqrt{m_i/(2\pi k_B T_c)} \) and \( \Lambda_i = \log \mu_i/(\Gamma + \tau_i) \). Here terms with the index \( i \) apply only to the ion species. The expressions presented here are obtained by considering the finite ion temperature acoustic sound speed, \( c_i = \sqrt{(Z_i T_i + T_e)/m_i} \), instead of \( c_0 \) in [35]. This results in the additional \( \Lambda_i \), and the normalisation scheme yields the extra \( \sqrt{(\Gamma + \tau_i)/\Gamma} \) in \( \Gamma_d \).

2.2. Numerical implementation

Our code TOEFL [36] is based on the delta-\( f \) isothermal electromagnetic gyrofluid model [29] and uses globally consistent flux-tube geometry [37] with a shifted metric treatment of the coordinates [38] to avoid artefacts by grid deformation. In the SOL region a sheath boundary condition model is applied [35, 40]. The electrostatic potential is obtained from the polarisation equation (3) by an FFT Poisson solver with Dirichlet boundary conditions in the (radial) \( r \)-direction. Gyrofluid densities are adapted at the \( x \)-boundaries to ensure zero vorticity radial boundary conditions for finite ion temperature. An Arakawa–Karniadakis scheme is employed for advancing the moment equations [41–43].

2.3. Shot-noise model

The statistical model is outlined in detail in [1, 3, 4]. The underlying assumption is that filament-propagation in the SOL is comprised of uncorrelated pulses of shape \( \psi(t) \) at time \( t \) such that the particle density fluctuations \( \Phi(t) \) recorded at a single point is given by a superposition of arriving blobs

\[
\Phi(t) = \sum_{k=1}^{\mathcal{K}(t)} A_k \psi(t - t_k),
\]

where a blob \( k \) of amplitude \( A_k \) arrives at time \( t_k \). During the total time of measurement, \( T \), there arrive precisely \( \mathcal{K} \) pulses. The duration of each pulse is given by

\[
\tau_d = \int_{-\infty}^{\infty} dt |\psi(t)|.
\]

Further assuming that pulses arrive according to a Poisson distribution, it follows that the waiting times between subsequent blob events is given by the exponential distribution. The average waiting time is denoted by \( \tau_w \) and can be estimated by \( \tau_w = \mathcal{K}/T \) or by fitting an exponential function to recorded values of waiting times. The stochastic process at hand then features the intermittency parameter \( \gamma = \tau_d/\tau_w \), quantifying the degree of pulse overlap (strong overlap for \( \gamma \gg 1 \) and vanishing overlap for \( \gamma \ll 1 \)). As \( \gamma \to \infty \), the probability density function for \( \Phi \) approaches the normal distribution. Blob shapes may be modelled according to

\[
\psi(t) = \begin{cases} 
\exp(t/\tau_i) & \text{for } t < 0, \\
\exp(-t/\tau_i) & \text{for } t \geq 0,
\end{cases}
\]

where the pulse consists of a trailing wake with rise time \( \tau_i \) and a steep front with fall time \( \tau_i \) such that the whole pulse lasts \( \tau_d = \tau_i + \tau_f \). For exponentially distributed filament amplitudes, it follows [3] that the stationary probability density function for the particle density is the gamma distribution

\[
P(\Phi) = \frac{\gamma^{\gamma/2}}{\Gamma(\gamma)} \left( \Phi + \gamma^{1/2} \right)^{\gamma/2} \exp\left(-\gamma^{1/2} \Phi - \gamma \right),
\]

where the mean value \( \langle \Phi \rangle \) and standard deviation \( \Phi_\sigma \) define the normalised variable

\[
\hat{\Phi} = \frac{\Phi - \langle \Phi \rangle}{\Phi_\sigma}.
\]

and the shape parameter may be found from the skewness, \( \gamma = 4/\sigma^2 \), or the flatness, \( \delta = 6/(F - 3) \), of the raw signal. The stochastic model consequently implies a parabolic relation between skewness and flatness, \( F = 3\sigma^2/2 \), typically observed in experiments [4, 5]. Note that for a normal distribution \( S = 0 \) and \( F = 3 \). The stochastic model also predicts an autocorrelation function for the case of two-sided exponential pulses such as in equation (9) with \( \tau_i = \lambda \tau_d \) and \( \tau_f = (1 - \lambda) \tau_d \):

\[
R(\tau, \lambda) = \frac{1}{1 - 2\lambda} \left\{ \left( 1 - \lambda \right) \exp \left[ -\frac{|\tau|}{(1 - \lambda) \tau_d} \right] - \lambda \exp \left[ -\frac{|\tau|}{\lambda \tau_d} \right] \right\}.
\]

3. Numerical simulations

We chose to simulate a domain of dimensions \( L_x \times L_y \times L_z = [128 \times 256] \rho_0 \times 2\pi \) with resolution \( 128 \times 256 \times 8 \). This corresponds to a box centred at the last closed magnetic flux-surface with approximate radial width 5 cm and 10 cm length in perpendicular direction. Runs are taken to \( T = 20000 \) normalised time-units with saturation occurring around \( t = 1000 \). Statistics are taken over the saturated state. Typical blob birth near the separatrix, located at \( x = 64 \), is depicted in figure 2, showing snapshots of the density (left) and potential (right) field at time \( t = 1000 \) for deuterium ions (left).

3.1. Statistical tools

Figure 3 (left) shows the probability density function (PDF) for the electron particle density measured at the outboard midplane at \( y = 128 \) and 10 grid points off the outermost
radial boundary, i.e. at $x = 117$. The artificial non-zero vorticity condition is restricted to the two outermost radial points and the probe location is sufficiently far off the boundary to not be influenced by it, while close enough to the outermost boundary in order to mimic a far-SOL Langmuir probe. The fitted normalised gamma distribution suggested from the stochastic model provides a reasonable fit, at least for the tail of the PDF. We have employed $\gamma = 4/S^2 \approx 5.7$ in equation (10) to obtain the dotted line. Similar to the experimental findings in [5], the autocorrelation function does not decrease to zero on the shown fit interval and the actual fit employed in figure 3 (right) results from $R = c + (1 - c)R$, where $c$ is a small offset from zero and $R$ is given in equation (12). Note that the offset is not important for the quantity of interest, $\tau_d$.

The fitted duration time is found to be $\tau_d = 1.1$. Figure 4 (left) shows the corresponding conditionally averaged waveforms for filament bursts with peak amplitude larger than 2.5 times the averaged fluctuation level, where the density signal is normalised according to

$$\hat{n} = \frac{n_e - \langle n_e \rangle}{n_{e,\text{rms}}}.$$  \hspace{1cm} (13)

and the conditionally averaged amplitudes are recorded for distinct peak events such that $A = \hat{n}|\hat{n}| > 2.5$. Fitting the pulse shapes in equation (9) gives $\tau_r = 0.8$ and $\gamma_f = 1.9$. For more details on conditional averages see [44–47]. The conditionally averaged radial velocity fluctuations, $V = \partial \phi/\partial y$, using the density signal at the same radial location as reference, is shown in figure 4 (right). Note that the maximum of the conditionally averaged velocity signal essentially gives a measure of the radial blob velocity as the conditional average is triggered whenever peak amplitudes are recorded in the density signal. Average burst amplitudes and waiting times are calculated by fitting truncated exponential distributions to the complementary cumulative distribution functions (CCDF)\(^4\) in figure 5. Note that for a stochastic variable $X$ with mean $\langle X \rangle$, distributed according to an exponential distribution, the expectation value for events with amplitude larger than a threshold, $y$, is given by

$$E(X|X > y) = y + \langle X \rangle.$$  \hspace{1cm} (14)

In figure 5 the threshold is $y = 2.5$ for the blob amplitudes giving a mean value of 3.8, in accordance with the maximum of the conditionally averaged signal in figure 4. The window length in figure 4 is 50 time steps, such that the truncation threshold for the waiting time fit is $y = 50$, yielding an averaged waiting time of 138. The deviation from the fit function is likely due to few blob events at large amplitude or waiting time. Complementary simulations with smaller box size that

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\(^4\)The cumulative distribution function (CDF) of a random variable $X$ gives the probability that $X \leq x$ for some number $x$ and is related to the probability distribution function (PDF) for $X$ through $\text{CDF}(x) = \int_{-\infty}^{x} \text{PDF}(\zeta) d\zeta$. The CCDF $1 - \text{CDF}(x)$ then quantifies the probability that $X > x$. For an exponentially distributed variable with mean $\langle X \rangle$ the CCDF reads $\exp(-X/\langle X \rangle)$.
were taken to $T = 40000$ reveal that more events produce a better fit. It should be pointed out that for $\gamma \geq 1$, the conditional average does not necessarily give the correct value for the underlying rise-, fall- and waiting times of the process due to significant pulse overlap [48]. In this contribution we find similar intermittency levels throughout the ion mass scan. One may assume that corresponding errors resulting from the conditional averaging technique then are similar for all ion masses considered.

3.2. Ion mass scan

In the same fashion we now employ the presented statistical techniques to signals obtained for setting the ion mass ratio, $\mu_i = m_i/Z\mu_{gi}$, to protium ($\mu_i = 1/2$), deuterium ($\mu_i = 1$), tritium ($\mu_i = 3/2$) and singly charged helium ($\mu_i = 2$) values. Figure 6 (left) gives the renormalised density bursts obtained from computing the conditionally averaged waveform, fitting a truncated complementary cumulative exponential distribution to the recorded conditional amplitudes ($A$) and renormalising by multiplying with the respective rms value ($\sigma$) and adding the mean value ($\langle n_e \rangle$) of the raw signal. Conditionally averaged peak amplitudes show no trend with respect to ion mass, however, the renormalised bursts do in terms of reduced bursts for higher ion mass (approximately $10\%$ decreased for $\text{He}^+$ compared to $\text{H}$). The presented trend is due to a significant reduction in the signals mean value for heavier ions (about $13\%$ lower for $\text{He}^+$ compared to $\text{H}$). As the probe location (see figure 2) is just off the radial boundary, this may be taken as a signature of gradually improved confinement for heavier ions. Similarly, we record the waiting times between conditional peaks and fit a truncated complementary cumulative exponential distribution, resulting in figure 6 (right), clearly showing that bursts are less frequent for higher ion mass. While blob amplitudes appear similar throughout the mass scan, increased waiting times between arrivals at the probe location may result in reduced mean value provided that the blob sizes are comparable. In figure 7 (left) we present results from fitting the two-sided exponential pulse shape from equation (9) to the conditionally averaged waveforms. We obtain estimates for the rise time, $\tau_r$, fall time, $\tau_f$ and total pulse duration, $\tau_r + \tau_f$. Furthermore, the fits to autocorrelation functions give a second estimate of the pulse duration. Increasing ion mass produces pulses of longer rise, fall and duration according to both analyses. The right part of figure 7 shows for all runs at hand a scatter plot of calculated skewness and flatness at 6 distinct radial locations, separated by $10\rho_0$, at $y = 32, 64, 96$ and the outboard midplane throughout the SOL region of the simulation domain. According to the stochastic model, there should be a parabolic relation between skewness and kurtosis. The figure 7 (right) shows that this indeed is a reasonable suggestion for the data presented, also indicating that this relation is universal in the sense that it does not significantly depend on the ion mass as values are scattered.

Figure 5. Complementary cumulative distribution functions for burst amplitudes (left) and waiting times (right) for deuterium ions. The red dotted line denotes fits from a truncated exponential distribution.

Figure 6. Average burst amplitude for the raw density signal (left) and waiting time (right) dependence on ion mass.

5 Note that fits for the waiting time CCDF throughout the mass scan are taken over at least an order of magnitude, even for $\mu_i = 2$, where events are less frequent, this then gives reasonable statistical significance. Of course, longer time series will always yield more reliable values.
randomly, except for protium, featuring slightly increased flatness, indicating enhanced frequency of prominent burst events compared to the other runs. Skewness and kurtosis increase radially outwards throughout the SOL, however, not monotonically as might be expected from the blob paradigm [13]. Similar features for the profiles have been shown in [9]. Furthermore, the intermittency parameter obtained from PDFs is of order unity and shows no clear trend with respect to isotopic composition. Signals are characterised by burst events of intermediate overlap.

A meaningful measure for the radial blob velocity at the probe location is given by the maximum of the conditionally averaged radial velocity, $V_x$, using the density signal as a reference. The conditional average on the radial velocity signal is consequently triggered by blob events recorded in the density signal. Figure 8 shows the blob velocity dependence with respect to ion mass. Inferred radial velocities range from 990 m s$^{-1}$ for singly charged helium ions to 1600 m s$^{-1}$ for protium ions. The isotopic dependence on the zonal flow velocity is depicted in figure 9. Heavier isotopes produce zonal flows of slightly increased velocity and wavelength in the radial direction. Furthermore, it should be noted that the zonal flow for the protium simulation is not persistent in time. Computing the zonal-flow shear, $\partial^2_t \langle \phi \rangle$, where $\langle \phi \rangle$ is the flux-surface averaged potential, we find no dependence on ion mass. It should be noted that the simulation presented here can be though of as typical for an L-mode scenario, hence the zonal flow is not dominating the turbulence.

4. Discussion

It is a well established fact that reduced blob detection frequency is associated with more pronounced shearing activity to break up radial streamers and decorrelate filaments [13]. According to [49] and references therein, shearing influences the perpendicular diffusion coefficients if $\omega_E > \gamma_L$, where $\omega_E \sim E_r \sim \partial^2_t \langle \phi \rangle$ is the $E \times B$ shearing rate, and $E_r$ the radial electric field. $\gamma_L$ is the maximum linear growth rate of the system. For either drift-wave or interchange turbulence, $\gamma \sim \mu^\alpha$, with $\alpha = -1$ for drift-wave instability and $\alpha = -3/4$ for interchange instability. The simulation at hand produces approximately constant values of maximum $\partial^2_t \langle \phi \rangle$ with respect to ion mass, hence, the ratio $\omega_E/\gamma_L \sim \mu^\beta$ with $\beta > 0$, providing at least a qualitative argument that the shear flow dynamics may favourably influence the confinement improvement in terms of reduced detection frequency for heavier isotopes. Another explanation may be reduced radial propagation velocities for heavier isotopes such as found in [28] for seeded blobs, together with overall slower dynamics (increased autocorrelation time), see [23], for edge turbulence, producing fewer blobs in a given time for increased ion mass. Taking for instance the ratio of duration times from either conditional averaging or autocorrelation analysis for protium and singly charged helium (see figure 7 (left)), $\tau_d(\text{H})/\tau_d(\text{He}^+)$ is approximately $1/2$, and comparing with the radial blob velocities from figure 8, $V_x(\text{H})/V_x(\text{He}^+) \approx 8/5$, we may calculate the corresponding size ratio of characteristic structures propagating past the probe ($\sigma = V_x\tau_d$):

$$\frac{\sigma(\text{H})}{\sigma(\text{He}^+)} = \left(\frac{8}{5}\right) \left(\frac{1}{2}\right) = \frac{4}{5}.$$  

Assuming purely radial propagation at velocity $V_x$ of coherent structures of scale $\sigma$ then leads to the conclusion that the observed increase in waiting time for heavier ions is due to 25% increase in size and 38% reduction in radial velocity,
comparing singly charged helium to protium. One may crudely estimate the number of particles per time unit \( N \) carried past the probe due to isotropic blobs according to

\[ N \sim A \sigma^2 / \tau_w. \]

For similar amplitudes we postulate

\[ \frac{N(\text{H})}{N(\text{He}^+)} \sim \left( \frac{\sigma(\text{H})}{\sigma(\text{He}^+)} \right)^2 \left( \frac{\tau_w(\text{He}^+)}{\tau_w(\text{H})} \right) \approx 1.11, \]

i.e. about 11% confinement improvement. In order to determine the total particle loss due to blob transport and its dependence on ion mass, blob detection routines similar to [9] or [26] should be employed to study the detailed structure size dependence in \( x \) and \( y \) direction.

Though the gyrofluid model features approximations (isothermal, delta-f, circular flux-surfaces) that severely limit its quantitative suggestions, we are confident that the findings in this work are robust in the sense that future experimental studies of the isotope effect might show similar trends as outlined in this contribution.

5. Conclusion

We have shown that a recently established statistical model for filament propagation in the SOL of fusion plasmas may be used to describe time-series for gyrofluid turbulence simulations that include self-consistent parallel boundary conditions to mimic the edge/SOL transition. The simulations show that burst amplitudes and waiting times follow an exponential distribution, consistent with an underlying Poisson process. Probability density functions are reasonably well approximated by gamma distributions and postulated autocorrelation functions are a good description of the measured signals. A parabolic relation between skewness and flatness moments seems to be present. All of this is consistent with the assumptions behind the stochastic blob model. It is shown that resulting electron density fluctuations close to the outermost radial boundary at the outboard midplane feature an improved confinement state in terms of reduced amplitude bursts for heavier isotopes. Typical burst events are shown to be of longer duration, with longer rise and fall times for increased ion mass. Regardless of ion mass the intermittency parameter is of order unity, characteristic for pulse arrival of intermediate overlap. The detection frequency of conditional blob events is higher for lighter isotopes. Maximum values of the conditionally averaged radial velocity with the density signal as reference shows that lighter blobs propagate faster at the probe location. Blob amplitudes, skewness, flatness and intermittency seem to be universal in the sense that they do not depend on the ion mass. Further studies, not making smallness assumptions on the relative amplitude of perturbations compared to the background should be executed, preferably through a full-F 3-d gyrofluid (or gyrokinetic) computational model in order to check the statistical dependency of blob properties [12, 26, 50, 51].

Figure 9. Flux-surface averaged flow velocity time evolution. Ions are: protium (top left), deuterium (top right), tritium (bottom left) and singly charged helium (bottom right).
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