Gyrofluid vortex interaction

Alexander Kendl

Institut für Ionenphysik und Angewandte Physik, Universität Innsbruck, Technikerstr. 25, A-6020, Innsbruck, Austria

E-mail: alexander.kendl@uibk.ac.at

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Abstract

Low-frequency turbulence in magnetised plasmas is intrinsically influenced by gyroscale effects across ion Larmor orbits. Here we show that fundamental vortex interactions like merging and co-advecting in gyrofluid plasmas are essentially modified under the influence of gyroinduced vortex spiraling. For identical initial vorticity, the fate of co-rotating eddies is decided between accelerated merging or explosion by the asymmetry of initial density distributions. Structures in warm gyrofluid turbulence are characterised by gyrospinning enhanced filamentation into thin vorticity sheets.

Keywords: plasma turbulence, fusion plasmas, gyrofluid model

(Some figures may appear in colour only in the online journal)

1. Introduction

Vortices can be regarded as the basic constituents of turbulence. Vortex motion and interactions govern nonlinear structure formation, flow and convective transport properties in a variety of fluids. A particular case of interest are quasi-two-dimensional fluids, which are characterised by the possibility for formation of coherent structures and large scale (zonal) flows, a dual cascade, and the ideal conservation of enstrophy in addition to energy [1, 2]. Merging and filamentation of vortices are fundamental processes that underly these properties. Examples for quasi-2D fluids encompass stratified and rotating fluids, atmospheric and oceanic flows, or the cross-field dynamics of magnetised plasmas [3].

In magnetised space, laboratory or fusion plasmas, the fluidlike convection perpendicular to a magnetic field \( \mathbf{B} \) is governed by drifts, which describe the mean motion on top of fast charged particle gyration. In particular, the magnetic confinement of fusion plasmas is crucially determined through turbulent transport generated by drift-type instabilities, and their suppression by (turbulence driven) zonal or equilibrium flows [4, 5]. Turbulent convection in magnetised plasmas is dominated by the electric drift velocity \( \mathbf{v}_E = \mathbf{E} \times \mathbf{B}/B^2 \) in presence of a fluctuating electric field \( \mathbf{E} = -\nabla \phi \). A localised electric potential \( \phi(x, t) \) leads to vortical \( \mathbf{E} \times \mathbf{B} \) drift motion of the plasma around the potential perturbation. The vorticity \( \Omega = \nabla \times \mathbf{v} \) of \( \mathbf{E} \times \mathbf{B} \) flows can be expressed as \( \Omega = (\mathbf{B}/B^2) \nabla^2 \phi \) for constant \( \mathbf{B} \). The electric potential \( \phi \) thus has the role of a stream function for \( \mathbf{E} \times \mathbf{B} \) flows.

Various instabilities driven by pressure gradients in magnetised plasmas typically result in drift vortex structures with sizes around the drift scale \( \rho_0 = \sqrt{m_i T_e/(eB)} \), where \( m_i \) is the ion mass, \( T_e \) the electron temperature, and \( e \) the elementary charge. For warm ion plasmas with non-zero temperature ratio \( \tau_i = T_i/T_e \sim 1 \), the drift wave and vortex scales are of the order of \( \rho_0 = \sqrt{m_i T_e} \), the ion finite Larmor radius (FLR). For the fast gyration of ions with particle charge \( q_i = Ze \) around their gyrocenters (at cyclotron frequency \( \omega_i = q_i B/m_i \) ) they effectively experience a ring-averaged electric potential \( \phi_i \) rather than the potential at the gyrocenter (for low-frequency perturbations with \( \omega \ll \omega_i \) ), and contribute to a quasi-neutral spatial distribution with approximately equal electron and ion particle densities \( N_e(x) \approx N_i(x) \) on average over the gyroorbit.

The framework for a description of magnetised plasmas governed by drift motion with FLR effects is efficiently given by gyrokinetic models evolving a 5D distribution function in phase-space [6–8], or by gyrofluid models of appropriate respective fluid moments in 3D space [9–12]. In the following, the effects of finite Larmor orbits on fundamental vortex interactions are analysed within an isothermal gyrofluid model [13]. It is found that spiraling of vorticity induced by...
FLR effects in the presence of spatial asymmetry significantly alters the merger process and generates fine structured vorticity sheets. Specific initial conditions result either in strongly accelerated merging or in vortex explosion, with consequences for inverse turbulent cascade properties and zonal flows.

2. Gyrofluid model for basic vortex interactions

Here we employ a basic local (‘delta-Γ’) gyrofluid model, derived from an energetically consistent gyrofluid electromagnetic model [14] in the limit of 2D isothermal dynamics in a straight and constant magnetic field. Details on the normalisation of the model equations and their numerical implementation can be found in [15]. Gyrofluid models do not dynamically evolve the particle densities \(N_i(x, t)\) of electrons and ions \((s = e, i)\), but rather their gyrocenter densities \(n_s(x, t)\). For small density fluctuation amplitudes both are connected through polarisation [10, 13, 16, 17] of gyroorbits by \(φ\):

\[
N_i = \Gamma_i n_s + \frac{\mu_s}{\tau_i} (\Gamma_{0i} - 1) \phi.
\]

The operator \(\Gamma_i = \Gamma_{0i}^{1/2}\) models gyroaveraging of the gyrocenter density, and \(\Gamma_{0i}\) reflects the gyroscreening of the potential \(φ\) in the mass dependent \((\mu_e ≡ m_e/m_i)\) polarisation term. In wavenumber space \(\Gamma_{0i}(k_i) = I_0(k_i)\exp(-b_i)\) is expressed by the modified Bessel function \(I_0\) with \(b_i = \rho_i^2 k_i^2 = \tau_i \mu_s (\rho_i k_i)^2\). For electrons, \(\mu_e \ll 1\), and the gyro radius \(\rho_e \ll \rho_i\) is negligible, so that \(b_e \approx 0\) and \(\Gamma_{0e} = \Gamma_{1e} \equiv 1\) can be assumed at ion gyro scales. The small electron mass also leads to negligible polarisation, so that particle and gyrocenter densities \(N_e \approx n_e\) coincide. The electron and ion particle densities are connected by strict quasi-neutrality to \(N_i ≡ n_i\) at scales much larger than the Debye length.

In the absence of driving, damping or parallel coupling the evolution equation for the gyrocenter densities in a homogeneous magnetic field then reflects incompressible mass conservation, and can (for normalised \(B = 1\)) in 2D be written as

\[
\partial_t n_s + [φ, n_s] = 0,
\]

where the advection term \(\mathbf{v}_E \cdot \nabla n_s\) is expressed in Poisson bracket notation by \([φ, n_s] = (\partial_s φ)(\partial_t n_s) - (\partial_t φ)(\partial_s n_s)\), describing advection of electron density along contours of the potential \(φ\). The gyroscreening of the potential \(\phi\) and of ion gyrocenter density along the gyroaveraged potential \(\phi = \Gamma_{1i} φ\). For numerical stability a hyper-viscosity term \(-ν_4 \nabla^4 n_s\) is added to the right hand side of equation (2).

The continuity equation (2) is closed by the quasi-neutrality condition \(N_i = n_e\), which by the gyrodensity relation in equation (1) gives the polarisation equation:

\[
\frac{1}{\tau_i} (\Gamma_{0i} - 1) \phi = n_e - \Gamma_{1i} n_i.
\]

The gyro-operators can be written in the Padé approximate forms \(\Gamma_0 = [1 + b^{-1}]\) and \(\Gamma_1 = [1 + (1/2) b^{-1}]\) and using \(b_i = -\tau_i \mu_i (\rho_i \nabla_i)^2\) in real space [10]. These expressions are useful for analytical considerations. Although strictly valid for \(b \ll 1\), the Padé forms give numerically nearly indistinguishable results compared to the original Bessel forms for all here presented simulations.

3. Vortex merging for cold ions

In the cold ion limit \(\tau_i = 0\) the polarisation equation reduces to \(\nabla^2 \phi = n_e - n_i\), which for normalised \(B = 1\) describes the deviation between electron and ion gyrocenter densities in relation to a scalar \(E \times B\) vorticity \(Ω = \nabla^2 \phi\). Subtracting the ion from the electron gyrocenter density equation, equation (2) in this cold limit transforms into the classical 2D Euler equation in vorticity representation:

\[
d_i Ω = ∂_j Ω = [φ, Ω] = 0 \quad (\text{for } \tau_i = 0).
\]

Vortex interactions, in particular merging and co-advection, are in this limit akin to classical fluids, and the gyrocenter densities are just passively advected. For \(\tau_i = 0\), a vortex that is defined by a localised spatial distribution \(Ω(x)\) can be initialised by any arbitrary initial density field \(n_i(x)\) with an appropriate choice of \(n_j(x) ≡ n_s(x) - Ω(x)\): the further evolution of \(Ω(x, t)\) in time through equation (4) will not depend on the particular initial gyrocenter densities, but only on their difference.

In the following it is shown how finite ion temperature with \(\tau_i > 0\) in magnetised plasmas leads to fundamentally different behaviour of the classical vortex merger and co-advection problems, and thus the resulting spectral properties of fully developed turbulence. In particular, the warm ion vortex merger problem intrinsically depends on the initial gyrocenter density distribution.

For comparability with classical fluid merger problems the vorticity is in the following initialised as a Gaussian \(Ω(x, t_0) ∼ \exp[-(x/σ)^2]\) of width \(σ\). In the gyrofluid model this may be obtained by a difference \(δ\) between the amplitude \(a\) of the electron compared to ion gyrocenter density, in the \((x, y)\) plane perpendicular to \(B\):

\[
n_{s0} = a \exp[-((x - x_0)/σ)^2 - ((y - y_0)/σ)^2],
\]

\[
n_e = \Gamma_{1i}^{-1}(1 - δ) n_{s0}.
\]

This initialises a vorticity \(Ω = Ω_0 = (n_{s0} - n_{s0}) = δ n_{s0}\) for both cold and warm plasmas. In the merger problem two vortices with the same amplitude are placed next to each other with an initial density peak distance \(Δ_0\).

The Euler equation as a standard model for fluid flow is, like the underlying Newtonian particle motions, invariant under parity transformation \(P: (t, x, v) \rightarrow (t, -x, -v)\). The evolution of flow patterns is thus symmetric with respect to simultaneous point reflection \((x) \rightarrow (-x)\) and reversal of the flow direction \((v) \rightarrow (-v)\). This implies that also for the fluid-like (cold plasma) case \(τ_i = 0\) a change in sign of the initial (pseudovector) vorticity distribution \(Ω(x) \rightarrow -Ω(x)\) leads to a spatially anti-symmetric evolution of the vortex merger.
Figure 1. Evolution of the vorticity $\Omega(x)$ field at several times during the co-rotating vortex pair interaction for $\gamma_i = 0$: anti-symmetric with the sign of the vorticity amplitude $\delta$ (top/red: positive, bottom/blue: negative).

For initialisation via equations (5) and (6), this can either be obtained by $n_{i0} \rightarrow -n_{i0}$ (with amplitude $a \rightarrow -a$), or alternatively by leaving $a$ positive but setting $\delta \rightarrow -\delta$.

In the cold fluid-like plasma case both choices have the same result: in figure 1 such a classical merging of co-rotating vortices is illustrated at several times for $\tau_i = 0$ with positive ($\delta = +0.01$, top) and negative ($\delta = -0.01$, bottom) initial vorticity.

The vorticity field $\Omega(x, y, t)$ is shown in a $(62.5\rho_i)^2$ central section (corresponding to 500² grid points) of the $(128\rho_i)^2$ computational domain (1024² grid points). The initial vortex radius for this case is $\sigma = 6\rho_i$ with an initial peak separation $\Delta_0 = 4\sigma$. The evolution is a typical example for a 2D fluid vortex merger: vortices orbit each other by mutual advection and after a while (depending on initial separation) develop encircling vorticity and density veils, and coalesce on combined advection-diffusion time scales into a spiraling single vortex [18]. The ensembles with inverted initial vorticity evolve exactly anti-symmetrical: negative vorticity inverts the direction of co-rotation and of the final spiral arms of the merged vortex.

4. FLR effects on vortex merging

It had already been noted in the seminal work of Knorr et al [9] from 1988 on theory of FLR effects in a guiding centre plasma, that in a (reduced) exemplary simulation of vortex merging with and without FLR effects different behaviour appeared: while for zero Larmor radius the maxima remained separated, a coalescence had been observed for a FLR [9].

In the following it is shown, that such FLR effects on vortex motion and vortex interactions in warm plasmas are generally a result of a breaking of the axial symmetry of the vortex in presence of any spatial asymmetry in the initial gyrocenter density distributions, which leads to FLR induced vortex spiraling.

However, for finite $\gamma_i \geq 0$ no 'generic' merger scenario like in the fluid case can be constructed: the temporal evolution of vorticity does not only depend on the initial (for example Gaussian) distribution $\Omega(x)$, but further on the specific gyrocenter initial density distributions $n_e(x)$ and $n_i(x)$ that generate this vorticity.

For warm plasmas the densities are not any more passively advected like in the cold case with $d_i\Omega = 0$ from equation (4). The gyrofluid vorticity evolution can be understood more intuitively (compare [19]) in the long wave length limit ($b^2 \ll 1$) when the Padé form of the gyrooperators is Taylor approximated:

$$d_i\Omega \approx -\frac{n_i}{2}d_i\nabla_\perp^2 n_e + \frac{n_i}{2}[\Omega, n_i]$$  \hspace{1cm} (7)

$$ \approx -\frac{n_i}{2}d_i\nabla_\perp^2 n_e + \frac{n_i}{2}[\Omega, n_i].$$  \hspace{1cm} (8)

In the first line terms up to order $b^2$ are kept, and in the second line up to $b$. For homogeneous $\mathbf{B}$, the identity $(\nabla_\perp^2 v_E) \cdot \nabla_\perp n = [\Omega, n_i]$ has been used.

The first term on the right of equations (7) and (8) contains the (gyroviscous cancelled) diamagnetic vorticity $\Omega_d = \nabla_\perp p$ with $p \equiv \tau_i n_e$, so that the generalised vorticity $W = \Omega + \frac{1}{2}\Omega_d$ obeys $d_i W \approx \frac{1}{2}[\Omega, p]$. This can be interpreted as an FLR induced contribution to polarisation by advection of vorticity along isocontours of pressure $p$ [19].

The Poisson bracket vanishes when the isocontours of axially symmetric density and vorticity profiles coincide. For a deviation from exact axial symmetry this term gives a significant contribution to the evolution of the (generalised) vorticity. The vortex dynamics for $\gamma_i > 0$ thus depends on the...
specific initial values and momentary gradients of the gyrocenter densities.

Now the Gaussian merger is reconsidered for \( \tau_i = 1 \). Positive vorticity is initialised again by equations (5) and (6) with \( a = 1 \) and \( \delta = 0.01 \) and otherwise same initial conditions as above. The evolution of \( \Omega(x, t) \) for this case is shown in figure 2 (top row) for various times. It is observed that already in the early stages of evolution (\( t = 100 \omega_i^{-1} \)) the vortices acquire spiral arms which rapidly spin up into a radial vorticity fine structure. The merging process is around twice as fast compared to the cold plasma case. The relative separation \( \Delta/\Delta_0 \) between peak densities of the vortices is plotted as a function of time in figure 3: the first minimum of distance for the warm merger is reached at around \( t = 3000 \omega_i^{-1} \) (bold red curve) compared to the cold case at around \( t = 6000 \omega_i^{-1} \) (thin black curve). For long times (\( t > 8000 \omega_i^{-1} \)) in the second diffusive stage the further merging appears (at least for the present numerical resolution and with \( \nu_g = 10^{-5} \)) to occur on similar time scales as for the cold case.

The initialisation of negative vorticity is for \( \tau_i > 0 \) ambiguous. Negative initial density perturbation with \( a = -1 \) and \( \delta = 0.01 \) gives (similar to the cold cases of figure 1) an exactly anti-symmetric vorticity evolution: the same pattern is obtained as on the top of figure 2, with simultaneous reversal of colours (sign of vorticity) and central point reflection. (This trivial reversed case is not explicitly shown in figure 2.) Consequently the separation \( \Delta(t) \) for \( a = -1 \) follows the same (bold red) curve as for \( a = +1 \) (with \( \delta = +0.01 \) for both) in figure 3.

The situation completely changes when the same initial \( \Omega(x, t_0) \) is obtained by keeping \( a = +1 \) but setting \( \delta = -0.01 \) and thus changing the relative local differences between electron and ion gyrocenter densities. As can be seen by the second term on the right of equation (7), \( n_i \) determines the FLR effect on polarisation. This means that at the positive Gaussian (quasi-neutral) density perturbation, the ion gyrocenters are not any more shifted outwards (as \( n_i < n_e \) for \( \delta > 0 \)) but inwards (\( n_i > n_e \) for \( \delta < 0 \)). This changes all gradients of \( n_i \) and thus the dynamical FLR contribution to polarisation.

The effect of this relative polarisation reversal by \( \delta > 0 \) \( \rightarrow \delta < 0 \) on co-rotating vortices is significant: instead of a merger event, a vortex separation with increasing distance is obtained. This is illustrated as snapshots of vorticity in the bottom row of figure 2 and by the time evolution of the peak distances (blue dashed curve) in figure 3. Simulations for a range of initial separations \( \Delta_0 = 2 \) to 6 show the

![Figure 2](image1.png)

**Figure 2.** \( \Omega(x) \) during the co-rotating vortex pair interaction for \( \tau_i = 1 \). The sign of \( \delta \) as the difference between initial electron and ion gyrocenter density amplitudes determines the fate of the pair towards merging or separation. (\( t = 0 \) as in figure 1.)

![Figure 3](image2.png)

**Figure 3.** Relative vortex peak separation \( \Delta/\Delta_0 \) for initial distance \( \Delta = 4\sigma \) with \( \sigma = 6\rho_i \) for the classical fluid case (\( \tau_i = 0 \), thin black line) and the FLR cases with \( \tau_i = 1 \), and positive (bold red line) and negative (dashed blue line) initial vorticity relative to the magnetic field direction, respectively.
Figure 4. Explosive vortex repulsion out of two neighbouring ‘adiabatic blobs’ with initial Gaussian density and potential perturbations with $n_e(x, t_0) = \phi(x, t_0)$, which results in a shielded initial vorticity. $\Omega$ is shown at three times, and density $n_e$ also at $t = 5000 \omega_0^{-1}$.

Figure 5. FLR accelerated merging for ‘adiabatically’ ($n_e(x, t_0) = \phi(x, t_0)$) initialised Gaussian vorticity.
same general tendency, with expectedly faster merging for shorter separations.

5. Dependence of vortex interactions on initial conditions

It is not a priori clear which vorticity reversal method is physically more relevant. Vortices are in general not seeded, but appear dynamically mostly as a result of the specific effects of instabilities on electron and ion densities. An important mechanism for vorticity generation is the drift wave instability, driven by a nonadiabatic parallel electron response in the presence of a cross-field density gradient [4]. For low collisionality the relation between electron density and potential can however often be regarded as nearly adiabatic, following approximately a Boltzmann relation with $n_e \sim \phi$.

We can also construct a vortex merger problem for such ‘adiabatic’ vortices. The constraint $n_e(x, t_0) \equiv \phi(x, t_0)$ implies that initially $\Omega = \nabla^2 \phi = \nabla^2 n_e \sim t_0/2$. In particular, setting $\phi = n_e$ in equation (3) for a given $\phi(x, t_0)$ initialises $n_e \equiv \Gamma_{\Omega}^{-1} (1 + \frac{k^2}{2} (1 - \Gamma_0)) \phi$, which is readily evaluated in wave number space. The choice is now whether to initialise Gaussian density/potential ‘adiabatic blobs’ (which yields a shielded vorticity), or ‘adiabatic Gaussian’ vorticities. The latter can be achieved by inversion of a defined (Gaussian) $\Omega(x, t_0)$ to $\phi = n_e = \nabla^2 \Omega$.

The ‘adiabatic blobs’ case first develops small satellite vortices which are sheared off through the shielding reversed vorticity rings around the centres. Subsequent collision of the satellites between the vortices rapidly results in an explosive repulsion after formation of a vorticity tangle. The situation is depicted in figure 4.

The ‘adiabatic Gaussian’ case is shown in figure 5: the development of vorticity is more similar to the fluid-like case (compare figure 1), but again with a pronounced FLR induced filamentary fine structure and accelerated merging. The evolution of peak separation $\Delta$ is for both ‘adiabatic’ cases presented in figure 6. Both cases are again anti-symmetric after reversal of the initial amplitude.

From these examples it is obvious that the gyrofluid merger dynamics strongly depends on the initial vorticity and density distributions, in addition to parameters like initial relative vortex separation as in the fluid case.

6. Gyrospinning of asymmetric vortices

A unique effect that is here present for all warm ion gyrofluid vortices is FLR induced spinning. A related FLR spin-up has been observed before for the special case of an interchange unstable (magnetic curvature driven) ‘blob’, which is characterised by the formation of a dipolar potential on top of a monopolar (e.g. Gaussian) density or pressure perturbation [19–21]. While interchange ‘blobs’ can acquire spinning by a range of additional mechanisms [15, 22], the FLR spin-up is in the following shown to be a universal phenomenon and essentially a consequence of asymmetry. A single inviscid axially exactly symmetric vortex retains its shape. In figure 7 the spin-up of FLR spiral arms is demonstrated for a single elongated vortex ($r = 1.2 r_i$), initialised by equations (5) and (6) with $a = +1$ (left) and $a = -1$ (right).

From equation (8) the FLR polarisation contribution to the evolution of vorticity is given as $d_\nu \Omega \sim \frac{\omega}{2 \omega}$. In polar $(r, \theta)$ coordinates centred on the vortex, $\Omega, n \sim (\partial_t \Omega)(\partial_r n) - (\partial_\theta \Omega)(\partial_r n)$ at a unit radius. For any axially symmetric $\Omega(x) \sim n(x)$ the FLR polarisation vanishes.

For initial distributions close to Gaussian, $\partial_r \sim (1/\sigma)$ can be approximated. For symmetric $\Omega(t_0)$ but $n = n(\theta)$, as by elongation, with $\partial_\theta \sim i k_\theta$ a dispersion relation $\omega = (\tau_r/2 \sigma) k_\theta \text{sgn}(B \cdot \Omega)$ is obtained. A spatial density asymmetry thus spreads vorticity azimuthally with $v_r \sim (\tau_r/2 \sigma) \text{sgn}(B \cdot \Omega)$. Similarly, any radial density gradient from $n(r)$ leads to radial spreading of vorticity. The combined result is the spin-up of spiral arms as in figure 7 with orientation depending on the relative sign of magnetic field $B$ and vorticity $\Omega$. Two neighbouring vortices mutually induce initial asymmetries similar to elongation, resulting in rapid pre-merging spin-up.

As a side remark, asymmetry can enter not only via non-circular vortex initialisation, but also numerically through a coarse rectangular grid and too close boundary proximity. Grid size and resolution have to be chosen accordingly, that any grid spin-up artefacts evolve much slower than the physical time scales of interest.

7. Vorticity filamentation in vortex co-advection and turbulence

The complementary problem to merging of co-rotating 2D vortices is the straight co-advection of counter-rotating vortices. In figure 8 it is shown that FLR spin-up again significantly alters this type of vortex interaction: in a warm gyrofluid (bottom row) the vorticity filamentation slows the
joint propagation but separates the vortex cores compared to the fluid-like cold case.

In combination, merging and co-advection determine the interactions and cascade in a turbulent sea of 2D vortices. We find that in decaying turbulence initialised with a random distribution of density and vorticity fluctuations analogously to equations (7) and (8), the FLR induced spinning also leads to enhanced vorticity filamentation. Self-sustained drift-wave turbulence in inhomogeneous magnetised plasmas is in 2D effectively represented by the Hasegawa-Wakatani model [23]: the turbulent drive is maintained by a dissipative coupling term $d(\phi - n_e)$ added to the right hand side of equation (2) for electrons, which emulates parallel electron dynamics for a single parallel wave number by the parameter $d$. A comparison of the vorticity structure between $\tau = 0$ and $1$ in a saturated drift wave turbulent state for $d = 0.01$ is shown in figure 9, and demonstrates the persistence of FLR induced vorticity filamentation in fully developed turbulence.

Figure 7. FLR spin-up of vorticity spiral arms in an asymmetric vortex with positive (left) and negative (right) initial elongated vorticity distribution $\Omega(x)$.

Figure 8. Co-advection of counter-rotating vortices at two times for the fluid-like ($\tau = 0$) case (top row), and with FLR induced spin-up for $\tau = 1$ (bottom row).
Vorticity thinning has been identified as a possible explanation for the inverse energy cascade of 2D turbulence in general fluids [24, 25] and for drift wave turbulence [26], and is here shown to be strongly enhanced by FLR spin-up in a gyrofluid. The amplitude of vorticity is for $\tau = 1$ increased over the whole spectral range, while density fluctuation amplitudes are enhanced on intermediate ($\rho_0 k_i \sim 1$) scales. This can be directly seen by the larger extension of the predominant density fluctuations in the bottom right panel of figure 9.

In 3D warm gyrofluid computations of drift wave turbulence we find that the vorticity sheets are much less pronounced but still discernible. The parallel connection of fluctuations along the magnetic field lines in presence of radial zonal flow [28] and magnetic shear [27] distorts but not completely suppresses the filamentary spin-up. Results on FLR effects in 3D toroidal edge turbulence are going to be presented in a future publication. The present work has focussed on identifying and explaining the FLR induced spiral-arm spin-up in asymmetric vortices and its effects on basic vortex interactions.

8. Conclusions

In summary, spiraling of asymmetric vortices by gyroorbit effects has been identified as a novel FLR effect in magnetised plasmas. The spin-up of spiral arms in single vortices was understood as an effect of density asymmetries on the governing drift velocities through polarisation of gyroorbits. This FLR induced spiraling was shown to strongly impact all vortex interactions in quasi-2D magnetised plasma dynamics, with examples ranging from dual vortex merging, via co-advection, to fully developed (many-vortex) turbulence. The nature and morphology of drift wave turbulence, which is of overall importance in magnetised fusion plasmas, is essentially changed.

Gyrofluid and gyrokinetic simulations are able to consistently account for such FLR effects on vorticity filamentation by sufficient spatial resolution. 3D gyrokinetic and gyrofluid codes are routinely used to numerically study turbulent transport in core, edge and SOL fusion plasmas. The natural focus of interest in such simulations is usually mostly on specific instabilities, nonlinear phenomena like zonal flow emergence, transport scalings or transport reduction, and identification of comparable observables mostly of statistical nature for experimental validation. Simulations of immediate fusion relevance are thus usually operated with as much physics included as is computationally affordable, but at the same time for efficiency restricted to the most coarse grid possible while maintaining convergence.

Then again, it is an essential feature of physics to seek to reduce phenomena down to the most fundamental aspects for better understanding. Vortices may be regarded as the
‘fundamental particles’ of turbulence. The interaction of vortices underlies all nonlinear turbulent transport dynamics and is therefore at the core of understanding turbulent transport in fusion plasmas. While the single spinning and dual interactions of vortices are of course sophisticated which are not directly observable in any fusion plasma, the combined effects on turbulent transport and spectral properties will always enter into any sophisticated modelling, although they may be well disguised by other often dominant effects such as toroidicity, shearing, collisionality, magnetic flutter, temperature fluctuations, and many more.

On a marginal note, actually this little essay has had its foundations in present efforts to develop sophisticated 3D full-6-moment electromagnetic toroidal edge/SOL gyrofluid turbulence codes. In basic convergence tests for one of our evolving codes, grid refinement towards high resolutions has in some situations produced vorticity filamentations similar to those shown in figure 9. As such fine structure has never been reported in any so far published work (known to the author) on gyrokinetic or gyrofluid turbulence, the first idea was that we had encountered some numerical artefact, perhaps having to do with possible problems regarding some specific numerical handling of FLR effects. After numerical issues had been ruled out (partly by cross-verification between our two gyrofluid codes ‘FELTOR’ and ‘TOEFL’ using similar models but fundamentally different numerical methods), simple test cases were devised, and the essence of the present study has emerged.

How now actually the newly identified FLR effects affect particle and energy confinement in more complete and complex simulation scenarios is an important question, which is beyond the scope of this basic presentation. The identification and evaluation of such FLR effects in the broader context of fusion edge turbulent transport, and possible impacts on (zonal or mean) flow generation, is going to be a subject of work to come.

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ORCID iDs

Alexander Kendl @ https://orcid.org/0000-0002-4270-9160

References

[1] Kraichnan R H and Montgomery D 1980 Rep. Prog. Phys. 43 547
[2] Boffetta G and Ecke R E 2012 Annu. Rev. Fluid Mech. 44 427
[3] Horton W and Hasegawa A 1994 Chaos 4 227
[4] Horton W 1999 Rev. Mod. Phys. 71 735
[5] Scott B D, Kendl A and Ribeiro T 2010 Contrib. Plasma Phys. 50 228
[6] Hahn T S 1988 Phys. Fluids 31 2670
[7] Brizard A J and Hahm T S 2007 Foundations of nonlinear gyrokinetic theory Rev. Mod. Phys. 79 421
[8] Krommes J A 2012 The gyrokinetic description of microturbulence in magnetized plasmas Annu. Rev. Fluid Mech. 44 175
[9] Knorr G et al 1988 Phys. Scr. 38 829
[10] Dorland W and Hammett G 1993 Phys. Fluids B 5 812
[11] Beer M A and Hammett G W 1996 Phys. Plasmas 3 4046
[12] Scott B 2005 Phys. Plasmas 12 102307
[13] Scott B 2003 Plasma Phys. Control. Fusion 45 A385
[14] Scott B 2010 Phys. Plasmas 17 102306
[15] Kendl A 2015 Plasma Phys. Control. Fusion 57 045012
[16] Pfirsch D 1984 Z. Naturforsch. 39a 1 (http://zfn.mpdl.mpg.de/data/Reihe_A/39/ZNA-1984-39a-0001.pdf)
[17] Brizard A J 2013 Phys. Plasmas 20 092309
[18] Leweke T, Le Dizes S and Williamson C H K 2016 Annu. Rev. Fluid Mech. 48 507
[19] Madsen J et al 2011 Phys. Plasmas 18 112504
[20] Wiesenberger M, Madsen J and Kendl A 2014 Phys. Plasmas 21 092391
[21] Held M, Wiesenberger M, Madsen J and Kendl A 2016 Nucl. Fusion 56 126005
[22] D’Ippolito D A, Myra J R, Russell D A and Yu G Q 2004 Phys. Plasmas 11 4603
[23] Hasegawa A and Wakatani M 1984 Phys. of Fluids 27 611
[24] Chen S et al 2006 Phys. Rev. Lett. 96 084502
[25] Bruneau C H, Fischer P and Kelly H 2007 Europhys. Lett. 78 34002
[26] Manz P, Birkenmeier G, Ramisch M and Stroth U 2012 Phys. Plasmas 19 082318
[27] Kendl A and Scott B D 2003 Phys. Rev. Lett. 90 035006
[28] Lin Z, Hahm T S, Lee W W, Tang W M and White R B 1998 Science 281 1835