Ion and electron temperature effects in the nonlinear dynamics of 2D collisionless reconnection

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One of the main subjects of investigation in the study of magnetic reconnection is today the role of pressure effects. An anistropic pressure is one of the two mechanism, together with electron inertia, allowing magnetic lines reconnection in low-collision plasmas; an isotropic scalar pressure, despite not beeing able to relax the linking between magnetic lines and plasma motion of the ideal MHD, nevertheless affects the dynamics of the process by introducing gyromotion effects. Even if a consistent accounting of particles pressure would require a kinetic approach, pressure effects can be included also in fluid models. Here we recall some recent results about the role of FLR effects in magnetic reconnection in an isothermal slab plasma that have been recently discussed in [1], and we present some new results about the different dynamics that ion FLR effects and electron temperature effects have on the dynamics of the current layers inside the magnetic island. The model we are considering [2] is a two-field model for a collisionless plasma in a slab geometry (the x-y plane) with a strong guide field (the large, uniform component B_0 being along the z direction), where isotropic electron and ion pressures are included by means of an isothermal thermodynamic closure for both the species. The uniform and constant temperatures T_α are defined by $P_\alpha = n_\alpha T_\alpha$, $\alpha = e, i$, and bring in the model the two characteristic lengths $\rho_i = (m_i c^2 T_i / Z^2 e^2 B_0^2)^{1/2}$ (ion Larmor radius) and $\rho_s = (T_e/T_i)^{1/2} \rho_i$ (ion-sound Larmor radius). The mechanism triggering the magnetic reconnection is the electron inertia, which introduces the electron skin depth $d_e = c/\omega_{pe}$ as a third characteristic scale length. In the large Δ' regime for cold ions and electrons, d_e is the typical width of the reconnection layer up to the early nonlinear phase.The fields whose dynamics is described in the x-y plane are the electrostatic potential φ, playing the role of the stream-function of the fluid at $\mathbf{E} \times \mathbf{B}$-drift velocity, and the z component of the vector potential, $\psi = A_z$, that is the stream function of the shear magnetic field in the plane, $\mathbf{B}_\perp = \nabla \psi \times \mathbf{e}_z$. These two fields are related one each other via the ion and electron densities, which are equated thanks to the quasi-neutrality condition and expressed in terms of a nonlinear generalization of the Padé approximation of a kinetic ion response. The ion density is then described by the field U, which in the cold ion limit ($\rho_i = 0$) is proportional to the fluid vorticity. In the most general case ($\rho_i, \rho_s \neq 0$) our equations normalized to B_0, to the characteristic equilibrium length L and to the Alfvèn speed c_A evaluated on the slab at $|x| \to \infty$, can be written by introducing the brakets $[f,g] = (\nabla f \times \nabla g) \cdot \mathbf{e}_z$ as

$$\frac{\partial F}{\partial t} + [\varphi, F] = \rho_s^2 [U, \psi], \quad \frac{\partial U}{\partial t} + [\varphi, U] = \left[\psi, \nabla^2 \psi\right], \tag{1}$$

$$F = \psi - d_e^2 \nabla^2 \psi, \quad U - \rho_i^2 \nabla^2 U = \nabla^2 \varphi. \tag{2}$$

Equations (1-2) have already been integrated in [1, 3]. Both ρ_s and ρ_i have a regularizing effect over the gradients of the ion density U, whose peak values at saturation result greately reduced (almost an order of magnitude) with respect to the cold regimes. As shown in [1], this can be understood by the fact that ρ_s and ρ_i introduce in the Hamiltonian density

$$\mathscr{H} = |\nabla \psi|^2 + d_e^2 |\nabla^2 \psi|^2 + (\rho_s^2 + \rho_i^2)|U|^2 - U \nabla^{-2} U, \tag{3}$$

two terms, namely $\rho_s^2 |U|^2$ and $\rho_i^2 |U|^2$, which when integrated over the plasma volume can be interpreted as compression works, $P_e dV$ and $P_i dV$ respectively. The amount of total energy which during the reconnection can be transferred from the magnetic components (first two terms of (3)) to the plasma kinetic contributions, is essentially fixed by the initial conditions, that is by the equilibrium profile and by the d_e parameter. The appearence of the compression work terms takes a fraction of the energy that in the cold case would go to the last term of Eq.(3). In this latter case the development of strong gradients in the field U is thus favored. The $-U\nabla^{-2}U$ term too contains in principle a dependence on ρ_i, but in the $\rho_i = 0$ limit it reduces to the purely kinetic contribution $|\nabla \varphi|^2$, due to the ions moving in the plane at the $\mathbf{E} \times \mathbf{B}$-drift velocity. Even if from the energetic point of view the global impact of ρ_i and ρ_s on the reconnection process is qualitatively very similar, the presence of ρ_i in the second of equations (2) modifies the relative evolution of φ and U with respect to the cold ion case. In the limit of $-\rho_i^2 \nabla^2 U$ being dominant at l.h.s. of (2), U and φ become locally proportional, differently from what happens at $\rho_i = 0$, $\rho_s \neq 0$. In [1] this has been shown to occour at saturation. More in general, the evolution of φ and U has been shown to be quite different if equal values of ρ_s and ρ_i were compared by keeping one of the two parameters equal to zero. An equally remarkable difference has however not been observed for the evolution of the current layers $J_z = -\nabla^2 \psi$ and of the field F. The similarity in the evolution of F induced by ρ_s and ρ_i has been explained in terms of the dynamics of the Lagrangian invariants, as it was used in Ref.[4] to explain the transition of regime from the eddy-like dynamics at $\rho_s \geq d_e$, $\rho_i = 0$ discussed in [5, 6] to the fluid Hasegawa-Mima regime at $\rho_s = \rho_i = 0$ of [7]. For $\rho_s = 0$ and $\rho_i \neq 0$ it can be shown indeed that two quasi-Lagrangian invariants $\tilde{G}_\pm = F \pm \rho_i d_e U$ exist, advected by the stream lines of $\tilde{\varphi}_\pm = \nabla^{-2} U \pm (\rho_i / d_e) \psi$ as long a source term $-d_e^2 \rho_i^2 \left[U, \nabla^2 \psi\right]$ keeps negligible. These $\tilde{G}_\pm$

fields mirror the exact Lagrangian invariants $G_\pm = F \pm \rho_s d_e U$ advected by $\varphi_\pm = \varphi \pm (\rho_s/d_e)\psi$ at $\rho_s \neq 0$ and $\rho_i = 0$. The source term $-d_e^2\rho_i^2\left[U, \nabla^2\psi\right]$ brings its main contributions along the neutral line, next to the X- points and in between the X- and O-points. In the range of parameters investigated in Refs.[1, 3], however, the consequence of these contributions on the dynamics of F was less pronounced. A reason for this, based on algebraic arguments related to the weight of the chosen parameters, has been proposed in [1]. There it was also observed that, by chosing a sufficiently large ratio ρ_s/ρ_i for small enough values of both ρ_s and ρ_i, a competing effect between ion and electron temperature effects should become visible in the evolution of F before saturation (when $\rho_i^2\nabla^2 U$ dominates in the second of (2), the r.h.s. term of the equation (1) for F becomes $(\rho_s/\rho_i)^2[\varphi, F]$). In Figs.1,2 we then compare the results of new simulations run at $d_e = 0.3$ with the same numerical setup of Ref.[1] (with $N_x = 2N_y = 1048$ points) for the cases $\rho_s = 0.02$, $\rho_i = 0.1$ and $\rho_s = 0.1$, $\rho_i = 0.02$. The reason for this choice is that these two cases are symmetric both with respect to the sums $\rho_s + \rho_i$ (entering in the definition of the quasi-Lagrangian invariants for $\rho_s, \rho_i \neq 0$ -see [1]) and $\rho_s^2 + \rho_i^2$ (whose rational power enters in the linear growth rate of the instability for $\rho_s, \rho_i \geq d_e$), while they largely differ for the ratio ρ_s/ρ_i, respectively 0.2 and 5, and have relatively slow growth rates. These are numerically measured to be $\gamma = 0.133 c_A/L$ (for $\rho_s = 0.02$, $\rho_i = 0.1$) and $\gamma = 0.137 c_A/L$ (for $\rho_s = 0.02$, $\rho_i = 0.1$) and respectively normalize the times in the following. At the beginning of the nonlinear phase ($t \simeq 12$, not shown here) the effects of ρ_s and ρ_i are almost indistinguishible but at later times (Figs.1) the differences over the small scale dynamics of J_z become remarkable and mirror those of U (Figs.2) also for the large ρ_i case. We then see that in the small ρ_s/ρ_i regime at $\rho_s, \rho_i \neq 0$, FLR effects favour the onset of a fluid-like dynamics of the type encountered at small ρ_s for $\rho_i = 0$ ([7]), in contrast to the laminar dynamics characteristic of large ρ_s and $\rho_i = 0$ ([5, 6]).

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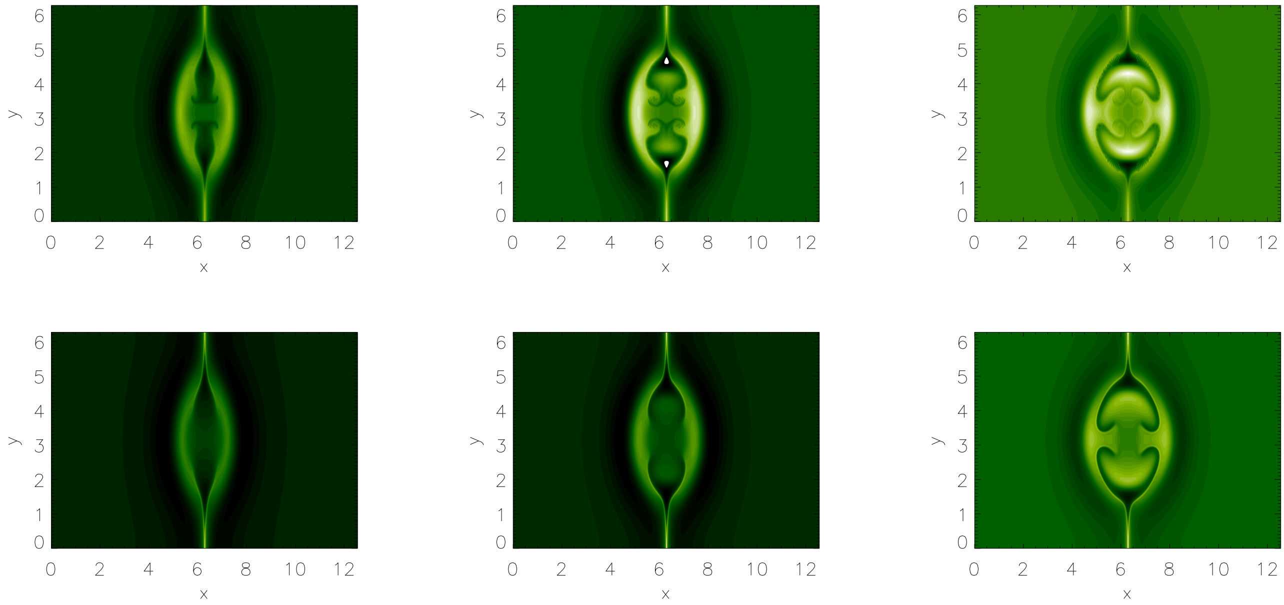

Figure 1: Contour plot of J_z at $t = 12.7, 13, 13.5$ for the case $\rho_s = 0.02$, $\rho_i = 0.1$ (top row) and $\rho_s = 0.1$, $\rho_i = 0.02$.

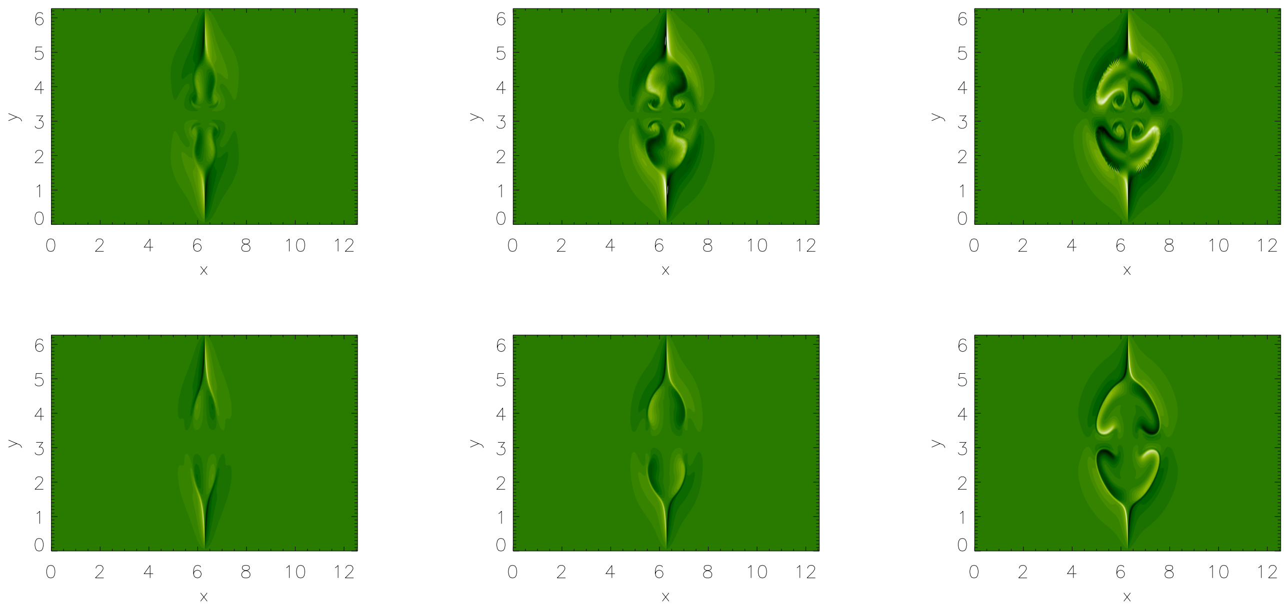

Figure 2: Contour plot of U at $t = 12.7, 13, 13.5$ for the case $\rho_s = 0.02$, $\rho_i = 0.1$ (top row) and $\rho_s = 0.1$, $\rho_i = 0.02$.