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Endogenous magnetic reconnection and associated high energy plasma processes

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ABSTRACT

An endogenous reconnection process involves a driving factor that lays inside the layer where a drastic change of magnetic field topology occurs. A process of this kind is shown to take place when an electron temperature gradient is present in a magnetically confined plasma and the evolving electron temperature fluctuations are anisotropic. The width of the reconnecting layer remains significant even when large macroscopic distances are considered. In view of the fact that there are plasmas in the Universe with considerable electron thermal energy contents this feature can be relied upon in order to produce generation or conversion of magnetic energy, high energy particle populations and momentum and angular momentum transport.

Magnetic reconnection processes have received increased attention lately as their effects have been proposed for the explanation of a variety of observations from space physics, such as the high energy particle production at the edge of the Heliosphere, to high energy astrophysics.

Now, the existence of endogenous magnetic reconnection processes is proposed whose characteristic feature is that their driving factor is contained within the region where a drastic change of magnetic field topology is produced. In contrast with this, well known reconnection processes such as those represented by the Sweet–Parker model or the resistive internal kink mode [1] are driven by factors (e.g. flows or plasma pressure gradients) that are outside the reconnection region.

The reconnecting mode that is discussed in the following is of the propagating type with a characteristic phase velocity and the two-fluid theory [2] on which it is based applies to weakly collisional regimes. Another significant finding is that the identified kind of mode involves widths of the layer in which reconnection takes place that remain relevant even when large macroscopic distances, such as those of interest to space and astrophysics, are considered.

We note that, the commonly held view that the main result of reconnection events is the violent conversion of magnetic energy into thermal energy, as is the case for large solar flares, is shown not to be correct in general. For instance, the main effect of the resistive internal kink mode [1], that can be excited near the center of toroidal plasma columns and involves magnetic reconnection, is of redistributing the radial profile of the plasma pressure.

The presence of a significant electron temperature gradient is a necessary feature of the reconnecting modes analyzed in this letter. Since in the Universe there are plasmas with considerable electron thermal energy contents it makes sense to rely on this feature to generate, through magnetic field reconnection, high energy particle populations, momentum and angular momentum transport and, in any case, conversion of electron thermal energy or of magnetic energy into each other.

The simplest equilibrium configuration that can be analyzed in weakly collisional regimes is the "sheared field configuration" in a plane geometry represented by $\mathbf{B} \simeq B_0(x)\mathbf{e}_z + B_y(x)\mathbf{e}_y$ and sketched in Fig. 1. The steps needed in order to transfer, with clear limitations, the results obtained for this plane configuration to cylindrical and axisymmetric toroidal configurations have been indicated previously [2].

The (normal) modes that we consider involve magnetic field perturbations represented by $\hat{\mathbf{B}} = \tilde{\mathbf{B}}(x)\exp(-i\omega t + ik_y y + ik_z z)$. The gradients of the particle density and temperatures are assumed to be significant for $|x| = |x_0| > 0$, as we choose to analyze modes such that $\mathbf{k}\cdot\mathbf{B}(x = x_0) = k_z B_0(x_0) + k_y B_y(x_0) = 0$ and, for $B_0^2 >> B_y^2$ with $B_0 \simeq const$, $\mathbf{k}\cdot\mathbf{B} \simeq (k_y B'_y)(x - x_0)$ around $x = x_0$. In particular, we refer to modes with $k_y^2 \lesssim |\partial^2/\partial x^2|$ in the region where reconnection takes place.

The electron and ion temperatures, T_e and T_i, are considered to be nearly isotropic in the equilibrium state but to have a different evolution in the perturbed state. In particular, we take $\hat{T}_{e\perp} \neq \hat{T}_{e\parallel}$ on account of the large anisotropy of the electron thermal conductivity. Then we note that $\nabla \cdot \mathbb{P}_e = \nabla p_{e\perp} + [\mathbf{B}\cdot\nabla(\bar{\bar{p}}_e\mathbf{B})]/4\pi$

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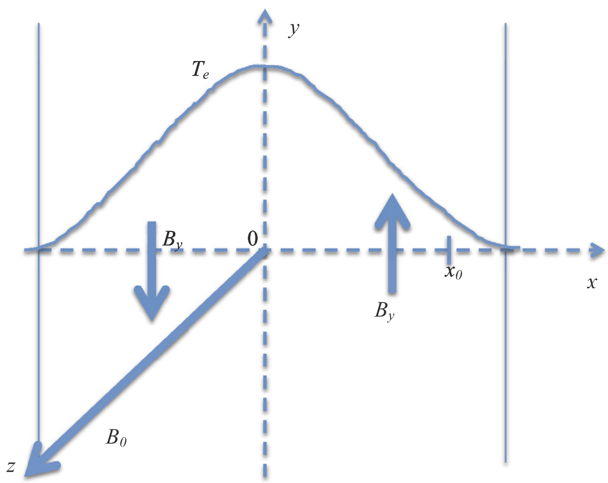

Fig. 1. Sketch of the simplest sheared field configuration.

where $\bar{\bar{p}}_e = 4\pi(p_{e\parallel} - p_{e\perp})/B^2$. Consequently $\nabla \cdot \hat{\mathbb{P}}_e \simeq \nabla \hat{p}_{e\perp} + n\mathbf{B}[\mathbf{B}\cdot\nabla(\hat{T}_{e\parallel} - \hat{T}_{e\perp})]/B^2$. We refer to weakly collisional regimes where ν_{ee}, the electron–electron collision frequency, is considerably smaller than $|\omega|$ considering that, as we shall show, $\omega \sim \omega_{di} \equiv [k_y c/(enB)](dp_i/dx)$. Clearly, $|\omega|^2 \ll \Omega_{ce}^2$ and in the considered regimes $\hat{T}_{e\perp} \simeq -\hat{\xi}_x dT_e/dx$, that is justified by the validity of the "equation of state" $d[p_{e\perp}/(Bn)]/dt = 0$. Here $\hat{\xi}_x = \hat{u}_{Ex}/(-i\omega)$ and $\hat{\mathbf{u}}_E$ represents the $\hat{\mathbf{E}}\times\mathbf{B}$ drift velocity.

Referring to $\hat{T}_{e\parallel}$ that, under the condition $\nu_{ee} < |\omega_{di}|$, is decoupled from $\hat{T}_{e\perp}$, we notice that in the limit of very large longitudinal thermal conductivity

$$(\widehat{\mathbf{B}\cdot\nabla T_{e\parallel}}) = \hat{B}_x \frac{dT_e}{dx} + i\left(k_y B_y'\right)(x - x_0)\hat{T}_{e\parallel} \simeq 0. \quad (1)$$

Clearly, Eq. (1) indicates that when $\hat{B}_x(x = x_0) \neq 0$, $\hat{T}_{e\parallel}$ would become singular [2] if $(dT_e/dx) \neq 0$, unless a transverse diffusion term is introduced in the relevant thermal energy balance equation. On the other hand the localized modes that we shall identify are characterized by eigenfunctions $\tilde{B}_x$ that are odd in $(x - x_0)$, that is $\tilde{B}_x \simeq (x - x_0)(d\tilde{B}_x/dx) \simeq -ik_y(x - x_0)\tilde{B}_y$. Clearly, for modes that have one important component of $\tilde{B}_x$ that is odd in $x - x_0$ within δ_m, where δ_m represents the width of the layer where reconnection takes place,

$$i(\mathbf{k}\cdot\mathbf{B})(\tilde{T}_{e\parallel} - \tilde{T}_{e\perp}) \simeq -\frac{dT_e}{dx}\left[\tilde{B}_x - i(\mathbf{k}\cdot\mathbf{B})\tilde{\xi}_x\right]. \quad (2)$$

We observe that Eq. (1) can be justified under the condition where $k_y^2\delta_m^4/L_s^2 > D^e_{\perp\parallel}/D^e_{\parallel\parallel}$, $D^e_{\perp\parallel}$ is the relevant transverse electron thermal conductivity, $D^e_{\parallel\parallel}$ the longitudinal electron thermal conductivity and $1/L_s \equiv B_y'/B$. Moreover, we note that for the modes of interest $\omega < k_\parallel v_{the}$ as $\omega \simeq |\omega_{di}| = k_y \rho_i v_{thi}/(2r_{pi})$, where $r_{pi} \equiv -(dp_i/dx)^{-1}p_i$, and $k_\parallel \sim k_y\delta_m/L_s$. Therefore, this condition implies $\rho_i/\delta_m < 2(r_{pi}/L_s)(v_{the}/v_{thi})$.

The adopted form of the perturbed total momentum conservation equation is

$$-i\omega n m_i \hat{\mathbf{u}}_i = -\nabla\left(\hat{p}_i + \hat{p}_{e\perp} + \frac{\hat{\mathbf{B}}\cdot\mathbf{B}}{4\pi}\right) - \frac{n}{B^2}\mathbf{B}\left[i(\mathbf{k}\cdot\mathbf{B})(\hat{T}_{e\parallel} - \hat{T}_{e\perp})\right] + \frac{1}{4\pi}\left[i(\mathbf{k}\cdot\mathbf{B})\hat{\mathbf{B}} + \hat{B}_x\frac{d}{dx}\mathbf{B}\right]. \quad (3)$$

Applying the operator $\mathbf{e}_z\cdot\nabla\times$ to Eq. (3) and considering $\hat{u}_{iz} \cong 0$, $\hat{B}_z \cong 0$, $\nabla_\perp\cdot\hat{\mathbf{u}}_i \simeq 0$, we find that the equation of interest within $|x - x_0| \sim \delta_m$ is

$$-\omega(\omega - \omega_{di})\left(\frac{d^2\tilde{\xi}_x}{dx^2} - k_y^2\tilde{\xi}_x\right) + \frac{1}{m_i n}\frac{k_y B_y}{B^2}\frac{d}{dx}\left[n(\mathbf{k}\cdot\mathbf{B})(\tilde{T}_{e\parallel} - \tilde{T}_{e\perp})\right] \simeq \frac{i}{4\pi m_i n}\left[(\mathbf{k}\cdot\mathbf{B})\left(\frac{d^2\tilde{B}_x}{dx^2} - k_y^2\tilde{B}_x\right) - k_y\frac{d^2B_y}{dx^2}\tilde{B}_x\right]. \quad (4)$$

Here we have taken $\tilde{u}_{ix} = -i(\omega - \omega_{di})\tilde{\xi}_x$ and $d^2B_y/dx^2 = (4\pi/c)(dJ_z/dx)$, J_z being the main component of the longitudinal current density, and introduce the scale distance $a_z \equiv J_z/(dJ_z/dx)$.

The longitudinal electron momentum conservation equation has a key role in the onset of magnetic reconnection as is well known. For the configuration that is being analyzed we adopt the following simplified form of this equation

$$0 \simeq -\left(\widehat{\frac{\mathbf{B}}{B}\cdot\nabla p_{e\parallel}}\right) - en(\hat{E}_\parallel + i\omega\mathcal{L}_I\hat{J}_\parallel), \quad (5)$$

where $\widehat{\mathbf{B}\cdot\nabla p_{e\parallel}}/B = ik_\parallel\hat{p}_{e\parallel} + (\hat{B}_x/B)dp_e/dx$, $\hat{p}_{e\parallel} = \hat{n}_e T_e + n\hat{T}_{e\parallel}$, and we neglect, at first, the contribution of a small but finite electrical resistivity. Instead, we have introduced [3] the inductivity $\mathcal{L}_I$ that is expected to prevail in the very low collisionality regimes that are being considered. In addition, we take $\hat{\mathbf{B}} = \nabla\times\hat{\mathbf{A}}$ and $\hat{\mathbf{E}} = -\nabla\hat{\Phi} + i(\omega/c)\hat{\mathbf{A}}$ and adopt a frame of reference where no $\mathbf{E}\times\mathbf{B}$ velocity exists in the equilibrium state, that is $d\Phi/dx = 0$. The vector potential $\hat{\mathbf{A}}$ is primarily in the z-direction, i.e., $\hat{\mathbf{A}} \simeq \hat{A}_z\mathbf{e}_z$. Then $\tilde{B}_x \simeq ik_y\tilde{A}_z$, $\tilde{B}_y \simeq -d\tilde{A}_z/dx$, $\tilde{B}_z \simeq 0$, $\tilde{E}_x \simeq -d\tilde{\Phi}/dx$, $\tilde{E}_y \simeq -ik_y\tilde{\Phi}$, and $\tilde{E}_z = -ik_z\tilde{\Phi} + i(\omega/c)\tilde{A}_z$. In particular, $\tilde{E}_\parallel \cong -ik_\parallel\tilde{\Phi} + i(\omega/c)\tilde{A}_z \cong -ik_\parallel\tilde{\Phi} + i[\omega/(ck_y)]\tilde{B}_x$. The components of the $\tilde{\mathbf{E}}\times\mathbf{B}$ drift velocity, are $\tilde{u}_{Ex} \simeq -ick_y\tilde{\Phi}/B$, $\tilde{u}_{Ey} \simeq (c/B)(d\tilde{\Phi}/dx)$ and $\tilde{u}_{Ez} = -c(B_y/B^2)(d\tilde{\Phi}/dx)$. Thus $\nabla\cdot\hat{\mathbf{u}}_E \simeq 0$, as $(k_z/k_y)(B_y/B_0) << 1$, and $\tilde{E}_\parallel \simeq \omega[\tilde{B}_x - i(\mathbf{k}\cdot\mathbf{B})\tilde{\xi}_x]/(ck_y)$. We determine $\tilde{n}_e$ entering $\tilde{p}_{e\parallel}$ from the simplified electron mass conservation $-i\omega\tilde{n}_e + \tilde{u}_{Ex}dn/dx \cong 0$, where the effects of particle transport and the contribution of $\tilde{u}_{e\parallel}$ have been regarded to be less important than that of the plasma inductivity when considering Eq. (5) at the same time. Then, $\tilde{n}_e \simeq -\tilde{\xi}_x dn/dx$ and, as $(k_z/k_y)(B_y/B_0) << 1$, $k_y^2 \lesssim |\partial^2/\partial x^2|$, we have $\tilde{J}_\parallel \simeq ic(d^2\tilde{B}_x/dx^2 - k_y^2\tilde{B}_x)/(4\pi k_y)$. Furthermore, we define

$$\mathcal{L}_I \equiv \frac{4\pi}{c^2}d_I^2, \quad (6)$$

where d_I is the "inductive skin depth". With $\tilde{T}_{e\parallel}/T_e$ given by Eq. (1), the equation that can be derived from the combination of Eqs. (5) and (6) is

$$\left(1 - \frac{\omega_{*e}}{\omega}\right)\tilde{B}_x \simeq i(\mathbf{k}\cdot\mathbf{B})\left(1 - \frac{\omega_{*e}}{\omega}\right)\tilde{\xi}_x + d_I^2\left(\frac{d^2\tilde{B}_x}{dx^2} - k_y^2\tilde{B}_x\right), \quad (7)$$

where $\omega_{*e} = -k_y[cT_e/(eB)](d\ln n/dx)$. We define $\bar{x} \equiv (x - x_0)/\delta_m$. Then Eq. (7) can be rewritten as

$$\tilde{B}_x(\bar{x}) \cong i(\mathbf{k}\cdot\mathbf{B}')\bar{x}\tilde{\xi}_x(\bar{x})\delta_m + \frac{\omega}{\omega - \omega_{*e}}\frac{d_I^2}{\delta_m^2}\left[\frac{d^2\tilde{B}_x(\bar{x})}{d\bar{x}^2} - \varepsilon_k^2\tilde{B}_x(\bar{x})\right], \quad (8)$$

where $\varepsilon_k^2 \equiv k_y^2\delta_m^2$, and we consider $|\omega| \sim |\omega - \omega_{*e}|$. We shall find that $\delta_m^2 < d_I^2$ for a realistic set of parameters.

Now we have the option to consider localized modes for which $|d^2\tilde{B}_x/d\bar{x}^2|/\tilde{B}_x \sim 1$, within the δ_m-region or extended modes for which $\tilde{B}_x(\bar{x}) \simeq \tilde{B}_{x0} = const.$ and $\tilde{B}_x \simeq \tilde{B}_{x0}[1 + (\delta_m^2/d_I^2)\varphi_0(\bar{x})]$. In the former case Eq. (8) reduces to