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Magnetic Pumping in Accretion Disk Coronae

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Abstract Most microquasar models involve a hot plasma in a corona above the disk or at the base of the jet. The Accretion-Ejection Instability (AEI) occurring in magnetized disks leads to the growth of a spiral density wave which can explain low frequency QPOs. It has already been shown to be very efficient in extracting accretion energy from the disk and emitting it upward in the corona as Alfvén waves. Here we present a simple mechanism which also allows the AEI to excite coronal ions. This heating is due to magnetic pumping, i.e. a resonant process occurring as the magnetic field lines emerging from the disk are periodically compressed by the spiral wave. We show how it acts on a collisionless population of ions, trapped above the disk by the joint action of gravity and magnetic stresses. We discuss the efficiency of this mechanism in heating coronal particles and explaining observational evidence.

Key words: accretion: accretion disks — corona — kinetic: gyrokinetic — magnetic pumping

1 INTRODUCTION

Several different features seem to indicate that accretion disks of microquasars and AGNs are embedded in a gas of hot plasma called corona. A power-law extending to high energy (with a cut-off at about 100 keV) is for instance observed in the spectrum of microquasars in the low-hard state. In many models, this tail is interpreted as the Comptonization of soft photons from the disk by a gas of hot electrons (Sunyaev and Titarchuk 1980). The structure and heating mechanism of the corona remain unclear although numerous authors have developed different models to fit the observations. However, many models assume that the corona extends above the disk plane in a slab geometry (Malzac 2001) and some observations may indicate that it could radially extend much further than what is thought (Church and Baluncińska-Church 2004). So far, the main models do not include magnetic field at all (Chakrabarti et al. 1996, Rózańska and Czerny 1996) or they have a magnetized, but very inhomogeneous corona, with magnetic loops (Galeev et al. 1979, Merloni and Fabian 2002, Liu et al 2002). In the later models, a large fraction of the gravitational energy is often assumed to be released in the corona by reconnection events (Haardt and Maraschi 1991). They have strong magnetic fields ($\beta \sim 1$) localized in small regions but the net vertical magnetic flux through the disk vanishes. We use here a completely different geometry, grounded on the results of jet models.

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YSO-, AGN- and microquasar jets are known to be very collimated. So far, the most favored models that can explain this property in a consistent way have been MHD models of jets (Blanford and Payne 1982, Pelletier and Pudritz 1980, Ferreira 1997). Their results give some constraints on the magnetic field strength and topology:

– They require the magnetic field to be strong enough to redirect upward the matter radially accreting and to load it at the base of the jet. Namely, the magnetic pressure has to be in equipartition with the gas pressure: $\beta \sim 1$
– Then, in order to accelerate the jet and collimate it efficiently, they constrain the magnetic field to be large scale and structured with a strong poloidal component, at least close to the disk. Contrary to the models mentioned above, the magnetic field varies on long time- and length scales.

In section 2.1 and 2.2, we present two interesting elements that can be derived from these conditions for the magnetic field. On one hand, the Accretion-Ejection Instability develops in the disk and leads to the growth of a spiral MHD wave. And on the other hand, coronal particles can be trapped in an oscillating motion along the poloidal field lines. In section 2.3 we show that a resonant process can occur between the oscillating motion and the periodic perturbation by the spiral wave. This mechanism, well known in plasma physics, is called magnetic pumping. In section **??**, we discuss the efficiency of magnetic pumping to heat coronal ions.

2 MODEL

The corona density is supposed to be very low, therefore the time between two collisions might be long in comparison with other time scales. This means that kinetic effects have to be included in the analysis. In the kinetic approach the usual fluid quantities as density, velocity, pressure, temperature... have to be replaced by the distribution function $\mathcal{F}(\boldsymbol{x}, \boldsymbol{v}, t)$ that gives the number of particles at time t and position $\boldsymbol{x}$ with the velocity $\boldsymbol{v}$. Without collision, the evolution of the system is then fully determined by the Vlasov equation which stands for the usual fluid ones (continuity, Euler, energy...). It reads:

$$\partial_t \mathcal{F} + \boldsymbol{v}.\boldsymbol{\nabla}\mathcal{F} + \Gamma.\partial_{\boldsymbol{v}}\mathcal{F} = 0, \tag{1}$$

where Γ is the total specific force applied to the particles. This equation cannot be analytically solved in the most general case, and further assumptions have to be made.

In a uniform and steady magnetic field, charged particles have a circular motion perpendicular to the field lines. They rotate with the cyclotron frequency: $\Omega_c = qB/m$ (where q is the charge, m the mass and B the amplitude of the magnetic field). Forces perpendicular to the field lines, like gravity, electric fields, non uniform or non steady magnetic fields, result in perpendicular drift motions in addition to the circular one. Forces parallel to the magnetic field have the usual effect on particle acceleration. A basic assumption is to consider that all time scales of the problem are longer than the cyclotron period. In this limit, it is known that the magnetic moment $\mu = v_\perp^2/2B$ is conserved. In this case, the cyclotron motion can be neglected and particles represented by their guiding center. The motion of the guiding center is then governed by the perpendicular drift velocities and by the classical parallel forces, but to keep information about the cyclotron motion and its effects, we have to add a new parallel force called the mirror force:

$$F = \mu \partial_{\|} B\,.$$

As the Lorenz force does not provide work, the energy has to remain constant when no other force is involved: $v_\perp^2 + v_\|^2 = cst$. And since μ is invariant, $v_\|$ has to decrease when B increases. The mirror force is thus a repulsive force that pushes charged particles away from strongly magnetized regions. This force is a kinetic effect and has not been used in previous models.

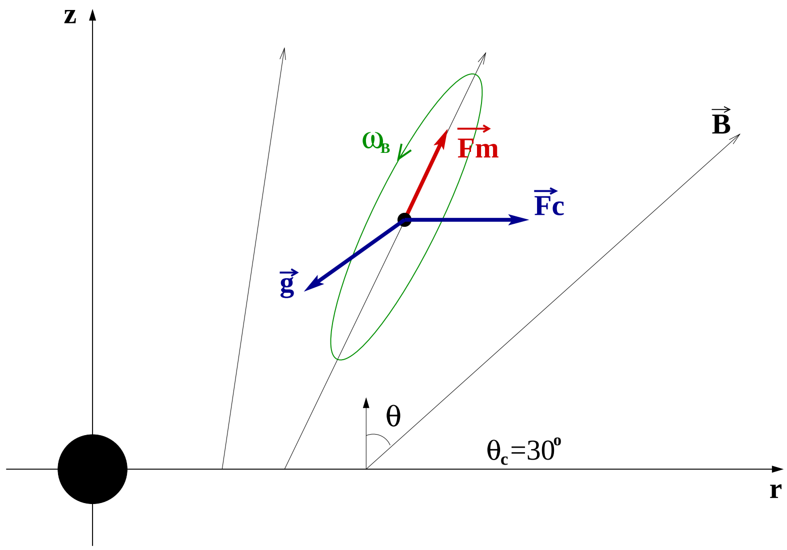

Fig. 1 Simple picture of the local geometry of field lines and forces: gravity $\boldsymbol{g}$, centrifugal force $\boldsymbol{F}_c$ and mirror force $\boldsymbol{F}_m$.

Assuming that μ is invariant makes the problem less complex, but it is still too hard to be solved. Therefore it is necessary to consider the MHD wave as a small perturbation added to the equilibrium described in introduction. All the quantities are thus written as the sum of an equilibrium part and a small perturbed one: the distribution function $\mathcal{F} = F^0 + f$, the forces $\boldsymbol{\Gamma} = \boldsymbol{\Gamma}^0 + \boldsymbol{\gamma}$, the magnetic field $B = B^0 + b$... The Vlasov equation to first order is:

$$\partial_t f + \boldsymbol{v}.\boldsymbol{\nabla} f + \boldsymbol{\Gamma}^0.\partial_{\boldsymbol{v}} f = -\boldsymbol{\gamma}.\partial_{\boldsymbol{v}} F^0.$$

This gives the perturbed distribution function, from which, we can derive the evolution of the perturbed system.

2.1 Equilibrium Periodic Motions

The first step in this perturbation method is to determine the equilibrium properties. To simplify, let us assume that the magnetic field lines are straight, oblique and purely poloidal as shown in Fig 1. In this geometry, the magnetic field decreases with radius and altitude and the mirror force has to be taken into account. As the disk itself is much denser than the corona, the effect of coronal currents on the field line topology can be neglected. The magnetic field is imposed by the disk, rotates with it and coronal particles have to move along the field lines. It means that all the perpendicular forces act together to keep particles on their field line and we can consider the particles as beads that can freely move along them (Henrisksen 1971). The parallel motion is then governed by the projections of gravity, centrifugal force and mirror force (Fig. 1).

An analytical analysis of this system can be done by assuming that particles have low altitude, precisely when:

$$h/\cos\theta \ll r,$$

where θ is the angle of the field lines with the vertical axis. This analysis leads to two different particle behaviors with respect to θ. Two domains are found, separated by a critical angle $\theta_c = 30°$:

– for $\theta > \theta_c$, the centrifugal force is stronger than gravity and the mirror force helps particles to escape. As a consequence, all particles are ejected as in the MHD case.
– For $\theta < \theta_c$, the mirror force acts against gravity and we find that some coronal particles can oscillate around an equilibrium position above the disk. If their magnetic energy is strong, this position can be high enough so that the particles freely oscillate without crossing the disk where they would otherwise collide dense matter. Particles can thus be trapped in a so called periodic bounce motion.

This motion is one of the two ingredients for magnetic pumping. To further characterize it, we can estimate its frequency. The bounce frequency varies with the angle of the field lines and, in the limit where the magnetic energy μB is low in comparison with the gravitational one, it is found to have the following dependence:

$$\omega_B^2/\Omega_K^2 = 1 - 4\sin^2\theta\,.$$

For vertical lines, this bounce frequency is the keplerian one: $\omega_B = \Omega_K$ but then the mirror force vanishes and there are no more trapped particles. And for critical lines, ω_B vanishes.

The purpose of this paper is to estimate the efficiency of magnetic pumping. As a consequence, we will now assume that the Vlasov equation, coupled with Maxwell equations and possibly radiative transfer equations has been solved for the equilibrium, giving a consistent solution for all fluid profiles (density, temperature...) and we will focus focus on the effect of the perturbation on the equilibrium.

2.2 Perturbation: Periodic excitation

In the conditions described previously for magnetic topology and strength, namely a poloidal magnetic field in equipartition with the gas pressure, we know that the Accretion-Ejection Instability (AEI) can develop in the disk and perturb the equilibrium described in the last section 2.1. The AEI leads to the growth of a MHD wave which rotates through the disk. It was worked out by Tagger and Pellat 1999 for a disk in vacuum and is now one of the best candidates to explain the low frequency QPOs (see Rodriguèz et al. 2002). This MHD wave couples the hydrodynamic and magnetic properties of the fluid and appears as a spiral wave, both in density and magnetic intensity: where the matter is denser, the magnetic field is stronger. Numerical simulations show that, in the disk, the perturbed magnetic field is about $b/B \sim 0.1$ (Caunt and Tagger 2001). Varnière and Tagger 2002 started to studying effects of the AEI for a disk surrounded by a low-density corona, as for instance the strong emission of Alfvén waves from the corotation radius to the corona. But the main point for us is that the magnetic perturbation induced by the AEI extends far above the disk (as e^{-qz} where q is the wave number in the disk plane) and can therefore affect the behavior of coronal particles. For radii smaller than the corotation radius, the spiral wave rotates more slowly than the gas. In the frame moving with the gas, coronal particles experience a periodic compression of the field lines and therefore a periodic perturbation of the mirror force with the frequency:

$$\widetilde{\omega} = \omega - m\Omega_K\,,$$

where ω is the wave frequency in the rest frame, m is the arm number of the spiral (typically, $m = 1$) and Ω_K is the keplerian frequency. The corresponding perturbed force is:

$$\gamma_{\parallel} = \mu\partial_{\parallel} b_{\parallel}\cos\widetilde{\omega}t.$$

2.3 Magnetic Pumping

The latter periodic force excites a periodic equilibrium motion. Where the frequencies are equal: $\omega_b \sim \widetilde{\omega}$, a resonance occurs. When the bounce motion is due to the mirror force, this resonance is the magnetic equivalent of Landau damping and is called transit-time damping or magnetic pumping. Indeed it pumps energy from the magnetic perturbation to the excited particles. The corresponding power is:

$$P = \int v_{\|}\gamma_{\|} f(\boldsymbol{x}, \boldsymbol{v}, t) d\boldsymbol{x} d\boldsymbol{v} \,.$$

From the last section 2.2, by performing a Laplace transform for time and a Fourier series decomposition on the bounce frequency for the parallel dimension, we find the perturbed distribution function and it eventually yields:

$$P = \sum_{n=-\infty}^{+\infty} \int n^2 \omega_B \mu^2 |b_n|^2 \delta(\widetilde{\omega} - n\omega_B) \partial_J F^0.$$

The first result is that the electrons oscillate much too fast to interact with the wave. As $\omega_B^{e-} \gg \widetilde{\omega}$, there is no resonance for the electrons and they cannot be heated directly by magnetic pumping. On the contrary, coronal ions move slowly enough for the pumping to become efficient. The magnetic perturbation exponentially decreases with altitude: $b(z) = b_{z=0} e^{-qz}$ with $qr \sim 1$. The Fourier components b_n can thus be calculated and, in the small altitude approximation ($z \ll r$), we find for the first resonance $n = 1$:

$$b_1 = \frac{1}{2} q h_c b_{z=0},$$

where h_c is the maximal altitude of coronal ions. Since the energy of the particles is mainly in their perpendicular (cyclotron) motion, we have $\langle \mu B^0 \rangle \approx T/m_i$, where the brackets note the average over the trapped ion population. We find the following heating time scale:

$$\Omega_K \tau = \left(\frac{\omega_B}{\Omega_K}\right)^2 \left(\frac{h_c}{r}\right)^{-2} \left(\frac{b_{z=0}}{B^0}\right)^{-2} \,.$$

The reference time is the keplerian period and we see that the stronger the perturbation, and the hotter the corona, the quicker the heating and so the more efficient the pumping.

3 DISCUSSION AND CONCLUSION

Going further and giving reliable figures for a consistent coronal model is a very difficult task. The corona is indeed very poorly constrained in the literature. Our goal is not to get global results yet. This would indeed need a full model with a precise description of many other features as cooling mechanisms... Here, we rather do a first attempt to estimate the typical heating time scale understand the significance of magnetic pumping.

For typical microquasars, we have $b/B \sim 0.1$, $T = 50\,\mathrm{keV}$, $h_c/r \sim .1$. The frequency ratio varies between 0 and and 1 respectively for vertical and critical field lines. This means that for critical lines, the pumping can be infinitely efficient. By choosing an angle of 25° , we find $\tau \sim 1s$. This time scale is comparable with the time for ions to collision the electrons and give their energy, but it is much smaller than the Compton cooling time. Some assumptions made to simplify the analytic resolution could increase the efficiency, when relaxed. This has to be checked, but this kinetic heating is likely to be unable to fight the Compton cooling.

Nevertheless, we can mention that if coronal ions are not collisional, the electron collision time is about the rotation time. As a consequence the kinetic approach becomes less valid for coronal electrons. Such a plasma, between the collisional and the non collisional regime, is known to experience the Braginskii viscosity 1965. For microquasar coronae, this viscosity is quite strong and could result in a high dissipation rate and thus a corona heating. This will be discussed in a future publication.

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