Two-temperature coronae in active galactic nuclei

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ABSTRACT
We show that coronal magnetic dissipation in thin active sheets that sandwich standard thin accretion disks in active galactic nuclei may account for canonical electron temperatures of a few $\times 10^9$ K if protons acquire most of the dissipated energy. Coulomb collisions transfer energy from the ions to the electrons, which subsequently cool rapidly by inverse-Compton scattering. In equilibrium, the proton energy density likely exceeds that of the magnetic field and both well exceed the electron and photon energy densities. The Coulomb energy transfer from protons to electrons is slow enough to maintain a high proton temperature, but fast enough to explain observed rapid X-ray variabilities in Seyferts. The $\sim 10^9$ K electron temperature is insensitive to the proton temperature when the latter is $\geq 10^{12}$ K.

Key words: scattering - radiation mechanisms - MHD - magnetic fields - galaxies: active - accretion discs

1 INTRODUCTION
UV and X-ray observations of active galactic nuclei (AGN), particularly in Seyfert 1 galaxies, indicate that the gravitational binding energy of a massive black-hole is dissipated partly in a cold accretion disk and partly in a hot corona above it. Thermal Comptonization of soft UV-radiation in the corona leads to the production of a hard X-ray continuum, some of which is reprocessed by the cold disk (Haardt & Maraschi 1991, 1993). In order to explain the different ratios of X-ray and UV luminosities in different objects, it has been suggested that the corona consists of localized active regions. It is likely that these are produced by magnetic fields in the disk amplified through differential rotation. When the disk magnetic field builds up significantly, buoyancy forces the field out and above the disk, giving rise to active regions of high magnetic field. Magnetic field buoyancy can maintain both magnetically and thermally dominated disks (Galeev, Rosner & Vaiana 1979, Shibata et al. 1990; Field & Rogers 1993).

Magnetic dissipation (for example in a reconnection site) may more effectively energize protons than the electron. Since it is the electrons which radiate, it is therefore plausible that the corona has two kinetic temperatures, that of the protons, $T_p$, and that of the electrons, $T_e$. These temperatures depend crucially on the energy exchange rate from protons to electrons.

In this Letter we show that an optically thin corona above a thin disk can generate the appropriate electron temperature to explain the observations of the hard continuum of AGN if electrons and protons are coupled via two-body Coulomb collisions. Two-temperature models are usually considered in the context of low density ion supported tori and advection dominated thick disks (e.g. Shapiro, Lightman and Eardley 1976; Rees et al., 1982; Narayan and Yi 1995). In our context however, although the thin magnetically active coronal sheets are much less dense than their underlying thin dense disk, they can still have a much higher density than standard advection dominated disks. We show that this enables enough energy transfer by Coulomb electron heating to satisfy constraints of rapid X-ray variability observed in many Seyfert I nuclei (e.g. Zdziarski et al 1994, 1995). Also, the active region need not be infalling, but can in fact be wind or ejection dominated and still transfer the required energy to electrons.

To obtain this solution, we balance the energy transfer rates from the magnetic field to the protons and from the electrons to radiation. Balancing these rates does not presume that any of the component energy densities are equal. We find that the proton and magnetic energy densities are both much larger than equilibrium electron and photon energy densities.

2 MAGNETIC ENERGY DISSIPATION
Magnetic dissipation sites (e.g. reconnection and/or shocks) in astrophysical plasmas always involve large gradients in the magnetic field, and the presence of length scales much smaller than the field gradients of the non-dissipating regions. This is because the induction equation governing the magnetic field...
\begin{equation}
\frac{\partial \mathbf{B}}{\partial t} = \nabla \times (\mathbf{v} \times \mathbf{B}) + \nu_m \nabla^2 \mathbf{B},
\end{equation}

where $\nu_m$ is the magnetic diffusivity, implies flux freezing and no dissipation as long as the last term is small. The ratio of this last term to the second is the magnetic Reynolds number $R_m = vL/\eta >> 1$, where $L$ and $v$ are the magnitude of velocity and gradient length scales. The large sizes and conductivities in AGN (and astrophysical systems generally) make dissipation negligible except in regions of large field gradients, i.e. thin current sheets (e.g. Parker 1979).

The more pronounced the variability of the AGN source, the more likely the fewer the number of coronal dissipation regions: if $N$ dissipation regions operate at a time, then a fluctuation changing the number of dissipation regions, would change the luminosity $\sim (N \pm n)/N$, where $n$ is the change in number of regions. If $N >> 1$, $n \sim N^{1/2}$ so that the effective change in luminosity would be small compared to when $N \sim 1$.

A natural site of coronal dissipation by reconnection and its slow shocks, would be at the interface between the newly rising field from the disk and any pre-existing coronal field (Begelman 1990; Shibata et al. 1990). Dissipation between pre-existing loops has also been considered (Field & Rogers 1993). In general, the dissipation layer could be expected to cover a sizable cross sectional fraction of the underlying disk.

The development of small turbulent length scales in the magnetically active regions enables the transfer of field energy into particles (e.g. Larosa et al. 1996). The turbulence generates a cascade of magnetic scales and the true dissipation scale is that below which the energy drains into particles rather than into smaller magnetic structures. The transfer of magnetic to particle energy likely occurs from MHD wave-particle interactions, for example by stochastic or shock Fermi energization (e.g. Achterberg 1987; Eilek & Hughes 1991; Larosa et al., 1996). Such energization requires two types of particle scattering, both of which likely favor protons for AGN coronae as we now discuss.

Particle scattering by MHD waves is the first required mechanism. The MHD waves are essentially magnetic compressions on scales much larger than the particle gyro-radii which can “mirror” particles as they approach. For stochastic Fermi energization, the scatterers move in random directions at $\sim v_A$, the Alfvén speed. Only the “head-on” collisions between particles and scatterers transfer energy to the particles. For particles moving faster than or of order $v_A$, both head-on and trailing scatterings occur. Even though much smaller particles do see only head-on collisions, there is a critical velocity perpendicular to the magnetic field required by particle in order for them to be scattered (Larosa et al. 1996). This latter condition dominates, and more energy is gained for faster particles. If the protons have higher velocities to begin with, then proton rather than electron scattering would be favored.

The above scattering can only enhance a particle’s momentum component along the field. In order to be repeatedly scattered and gain significant energy, the particles must be pitch angle isotropized with respect to the local magnetic field. This pitch angle scattering is accomplished by resonance with plasma Alfvén waves on the scale of the particle gyro-radii (Achterberg 1987; Eilek & Hughes 1991). Such waves only exist below a critical frequency and so the particles must have significantly large gyro-radii before gyro-resonance can occur. The required pitch angle scattering also favors the protons: in a plasma for which the Alfvén speed is of order $c$, the electrons need to be quite relativistic to resonate with Alfvén waves. Such favoring of protons has been called the “electron injection” problem (Eilek & Hughes 1991).

Shock Fermi energization may also operate across slow shocks in the reconnection region (Blackman & Field 1994). This produces only head-on collisions and is thus more efficient than stochastic processes. However, similar arguments to those above would apply, and protons would be more easily energized.

The previous discussion suggests that when the protons are already at higher kinetic temperatures than the electrons, magnetic dissipation can maintain such a state in equilibrium. Setting up the initial configuration is a separate problem since the initial plasma might not necessarily have protons at a higher temperature than electrons. However, even if the initial dissipation went into electrons, they would cool very fast. The proton energy density could slowly grow. After some time, the protons could acquire enough energy to dominate the electrons, leading to the two-temperature equilibrium we describe herein. We leave a more detailed investigation for future work.

Simulations (Priest & Forbes, 1986) as well as solar (Tsuneta 1994) and the geomagnetic tail observations (Coroniti et al. 1994) indicate that magnetic reconnection occurs in the presence of slow shocks. This suggests that Petschek (Petschek 1962) type reconnection can occur in nature, distinctly from Sweet-Parker (Sweet 1958; Parker 1979) type reconnection. The main difference between the two for non(or mildly)-relativistic Alfvén speeds is that for the latter, the ratio of height to width of the the overall dissipation region satisfies is $h/r \sim v_{dis}/v_A \sim 1/R_m^{1/2}$ while in the former the ratio is $h/r \sim v_{dis}/v_A \sim 1/\log[R_m]$ where the $R_m >> 1$ is measured outside of the dissipation region. The Sweet-Parker dissipation region is generally thinner and magnetic energy conversion occurs much slower. Our calculations below, combined with observational constraints, seem to suggest that $h/r \sim 1/10$ is reasonable, possibly suggesting that the dissipation regions are Petschek-like.

Finally, note that the processes discussed above might generate non-thermal proton distributions. However, the electron energization times are also of order their Maxwellian relaxation times (Spitzer 1956). Thus the electrons can be rapidly thermalized even if the protons are initially non-thermal.

### 3 A TWO-TEMPERATURE CORONA

#### 3.1 Time scales

We now motivate the model by exploring relevant time scales using simple approximate formulae, reserving a fuller treatment for the next Section.

Thermal Compton scattering gives the power-law shape to the observed hard X-ray spectrum and exponential cutoff of Seyfert 1 nuclei, provided that the plasma is optically thin to electron scattering (i.e. $\tau \lesssim 0.4$) and has an electron temperature $kT_e \sim 2 - 5 \times 10^9$ K. The Compton cooling time scale

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\[ t_{\text{cool}} \approx \pi t_{\text{cross}}/\ell \]  
where the compactness parameter \( \ell = L/r(\sigma_T m_e c^2) \). \( L \) is the luminosity of the active region which is of size \( r \). For a source such as MCG–6-30-15 which has a luminosity of \( 10^{43} \text{erg s}^{-1} \) and varies typically on a time scale of 1000 s then \( \ell \sim 10 \), so \( t_{\text{cool}} \sim t_{\text{cross}}/3 \), where \( t_{\text{cross}} = r/c \). The electrons cool on less than a light crossing time.

The electron-proton coupling time, \( t_{\text{ep}} \) is defined from the energy transfer equation (Spitzer 1956)

\[
\frac{dT_p}{dt} = \frac{(T_p - T_e)}{t_{\text{ep}}} = \frac{2 m_e n_e \sigma_T c \ln \Lambda}{\sqrt{\pi}} t_{\text{cross}},
\]

where we have assumed that the sum of the \( \theta_e = kT_e/m_e c^2 \) and \( \theta_p = kT_p/m_p c^2 \) is of order unity. Provided that \( T_e \ll T_p \), the coupling time to attain \( T_e \)

\[ t_e = t_{\text{ep}} T_p/T_e \sim 0.1 \frac{r}{\ell} t_{\text{cross}}. \]

Thus if \( T_e \sim 10^9 \text{K} \) and \( T_p \sim 10^{12} \text{K} \), \( t_e \approx t_{\text{cross}}/4 \), if \( \tau_T = 0.4 \), which means that \( t_e \approx t_{\text{cool}} \). This last condition is required for heating to balance cooling for the electrons and determines \( T_e \).

A slab geometry enables \( T_e \sim 10^9 \text{K} \) for lower values of \( T_p \). If, for instance, the thickness \( h \) is one tenth of the width \( r \), the cooling time does not change, but the higher gas density means that the coupling time decreases by the factor \( h/r \). Moreover, if the slab is favourably oriented so that it is perpendicular to the line of sight, a luminosity variation can be observed on a time scale of \( h/c \), shorter than the one earlier assumed by the same factor. This enables occasional variation to be detected on 100 sec (see e.g. Reynolds et al 1995). Thus \( T_p \approx 30 h/r (\ell/\tau_T) T_e \).

### 3.2 Magnetic Field and Particle Energy Densities

We take \( \tau \) to be the Thomson optical depth given by \( \tau = n_e \sigma_T h \). The size \( h \) of an active region in the corona can be constrained by observations of the shortest luminosity variability time scales \( h/c \approx 100 \text{s} \) implying that \( h \sim 3 \times 10^{12} \text{cm} \). For \( \tau \approx 0.3 \), \( n \approx 1.5 \times 10^{13} \text{cm}^{-3} \).

Given that a fraction \( \eta \) of the accretion power is dissipated in the corona \((1 - \eta)\) in the underlying disk, we have

\[ \eta L = 4\pi r^2 \epsilon_c \approx 4\pi r^2 h(\Delta \varepsilon_{\text{m}}/dt), \]

where \( \epsilon_c \) and \( \epsilon_{\text{m}} \) are the photon and magnetic energies respectively and \( r \) is the radius of the effective dissipation sheet of thickness \( h \). Now, \( 4\pi r^2 h(\Delta \varepsilon_{\text{m}}/dt) \sim 4\pi r^2 v_{\text{dis}} \times B^2/8\pi \) so from eqn. 4 we have

\[ B^2/8\pi \sim \epsilon_c (c/v_{\text{dis}}). \]

We assume that the dissipation velocity \( v_{\text{dis}} \) is a fraction \( (f \lesssim 1) \) of the Alfvén speed \( v_A \), i.e. \( v_{\text{dis}} = fB/\sqrt{4\pi \rho} \). In a reconnection site though, the fraction \( f \sim h/r \). From eqn. 6 we then have

\[ B = 7.9 \times 10^3 \left( \frac{\eta L}{10^{40} \text{erg s}^{-1}} \right)^{1/3} \left( \frac{\rho}{10^{-13} \text{g cm}^{-3}} \right)^{1/6} \]

\[ \times \left( \frac{h}{5 \times 10^{12} \text{cm}} \right)^{-1/3} \left( \frac{r}{10 \text{gy}} \right)^{-1/3} \left( \frac{M}{10^3 M_\odot} \right)^{-1/3} G, \]

where \( r_g \equiv GM/c^2 \) and \( M \) is the central \((\text{black hole mass})\). The characteristic Alfvén velocity would then be \( v_A \approx 0.2c \).

The above estimate uses only the energy transfer rates. We now show that the assumption of equipartition between the magnetic energy density \( \epsilon_{\text{m}} \) and proton energy density \( \epsilon_p = \epsilon_m \) would likely underestimate the proton temperature. In the steady state we know that \( |(\Delta \epsilon_{\text{m}}/dt)| = |(\Delta \epsilon_p/\Delta t)| = |(\Delta \epsilon_p/\Delta t)| \), and the additional assumption of \( \epsilon_m = \epsilon_p \) would mean that \( t_{\text{ep}} \equiv \epsilon_p/(|\Delta (\epsilon_p/\Delta t)|) = \epsilon_m/(|\Delta (\epsilon_p/\Delta t)|) \equiv h/v_{\text{dis}} \). But using eq (3) with \( h \) in place of \( r \), this would imply \( v_{\text{dis}} \approx r_c/100 \sim c/300 \). Since the dissipation physics of section 2 suggests \( v_{\text{dis}} \sim h/r \), and variability constraints imply \( h/r \sim 1/10 \), the assumption of equipartition between the magnetic and proton energy densities would under estimate the proton temperature by a factor \( \sim 10 \) compared to what we find in the next section.

We will see that the reason for this is that \( t_{\text{ep}} \) can be one or two orders of magnitude longer than \( t_e \). Therefore, before a steady state can ensure, the proton energy density must build up to some critical value. The protons act as a growing sea of magnetic energy until they can transfer energy to electrons at the rate they gain energy from the magnetic field. Most of the energy in the active region thus resides in the hot protons. The least energy is contained in the electrons, whose energy density is less than the photon energy density \( \epsilon_c \) by the factor \( t_{\text{cross}}/t_{\text{cool}} \). The long \( t_{\text{ep}} \) also means that although a flare may rise on a crossing time, its decay time is uncertain. The luminosity might drop gradually as the proton energy is exhausted or abruptly if the region expands.

### 4 A DETAILED MODEL

#### 4.1 Heating of electrons by ions

We determine \( T_e \) by taking into account the relevant cooling processes. We solve \( (\Delta \epsilon_e/\Delta t)_+ = (\Delta \epsilon_e/\Delta t)_- \), taking the transfer for Coulomb collisions between populations of Maxwellian distributions using the Rutherford scattering cross-section (Stepney & Guilbert 1983). We show that typically, \( T_e \approx 10^9 \text{K} \) and that Coulomb interactions are fast enough to explain the observed variability. Future work should consider how sensitive the result is to the presence non-thermal protons, as the energization mechanisms of section 2 would likely provide them.

We write

\[
(\Delta \epsilon_e/\Delta t)_+ = \frac{3}{2} \frac{m_e}{m_p} n_e n_{\text{e}} \sigma_T c \frac{(kT_i - kT_p)}{K_2(1/\theta)_i K_2(1/\theta)_i} (2\ln \Lambda)
\]

\[
\times \left[ \frac{2(\theta_e + \theta_i)^2}{\theta_e + \theta_i} \right]^2 K_1 \left( \frac{\theta_e + \theta_i}{\theta_e \theta_i} \right) + 2K_0 \left( \frac{\theta_e + \theta_i}{\theta_e \theta_i} \right),
\]

where \( \ln \Lambda \approx 20 \), is the Coulomb logarithm, the \( K \)'s are the modified Bessel functions, and \( \theta_e, \theta_i \) are the dimensionless electron and ion temperatures previously defined. (Note that eqn. 6 is valid over all electron and proton temperatures).

Electrons in the corona are cooled through inverse Compton scattering off the soft photon field produced by the accretion disk, through thermal synchrotron radiation.
and its Comptonization. Hence, \((d\varepsilon_\gamma/dt)_+= q_{\text{sync}} + q_{\text{CS}} + q_{\text{Comp}}\). The dominant term is usually \(q_{\text{Comp}} = 4\theta_c \sigma_T n U_{\text{rad}}\) where \(\sigma_T\) is the Thomson cross section and \(U_{\text{rad}} = (1 - \eta)L/4\pi R^2\) in\(\text{disk}. The thermal synchrotron emission rises steeply with decreasing frequency, becomes self-absorbed below \(v_\gamma\) and gives rise to a blackbody spectrum. The cooling is strongly peaked around \(v_\gamma\) and can be approximated by \(q_{\text{sync}} = (2\pi/3)m_v \theta_v \varepsilon_v^2/\tau\) (Narayan & Yi 1995) where \(\theta_v\) is obtained by equating the Rayleigh-Jeans emission to the synchrotron one (Zdziarski 1985) The Comptonization of the synchrotron radiation \(q_{\text{CS}}\) is calculated using Dermer and Liang (1991) approximate treatment.

Fig. 1 shows the electron temperature solution as a function of the ion temperature when \((d\varepsilon_i/dt) = (d\varepsilon_\gamma/dt)_+\) is solved. The figure illustrates how Coulomb heating of electrons in the optically thin plasma of the localized regions naturally give rise to a \(T_\varepsilon\) consistent with X-ray observations of high energy spectral cut-offs in AGN. Such temperatures do not depend strongly on \(T_\varepsilon\) as long as protons are mildly relativistic. The \(T_\varepsilon\) is antecorrelated with \(\tau\); as \(\tau\) increases, the density \(\rho\) and consequently the magnetic field \(B\) increase. This implies more synchrotron and Comptonisation synchrotron cooling to balance an increased Coulomb coupling efficiency, resulting in a lower \(\theta_\varepsilon\). The temperature solution is not a strong function of the black hole mass \(M\) (see Fig. 1).

We note that for \(\theta_\varepsilon\) and \(\tau\) of interest, electron-electron scattering through Möller cross section competes with various loss mechanisms and in particular with Compton losses in our case. The question arises of whether electrons can thermalize where the electron-electron relaxation time scale \(\tau_{\text{ee}} = 4\sqrt{\pi} \theta_v^{3/2}/\ln(\Lambda/h/\sigma_T)\). Ghisellini, Guilbert & Svensson (1988) though, found that cyclo/synchrotron self-absorption as considered in our model, acts as a very efficient thermalizing mechanism as long as the magnetic energy density dominates the radiation energy density, which is a condition postulated here. We therefore expect electrons to be able to achieve a Maxwellian distribution.

### 4.2 Time scale for electron heating

In the previous section we have shown that electron-ion Coulomb coupling can explain the values of \(T_\varepsilon \sim \text{few} \times 10^{9}\)K inferred from observations. If particles are coupled only by two body collisions, the protons must pass energy to electrons fast enough to explain the rapid variability observed in X-ray sources. This requires \(\tau_{\text{cs}} \equiv \varepsilon_\gamma/(d\varepsilon_\gamma/dt) = \tau_{\text{esc}}/\varepsilon_\gamma\) has to be less than the light crossing time \(h/c\), where \((d\varepsilon_\gamma/dt)_+\) is given by eqn. (8). Now since \(n\sigma_T c = ct/h\), we have

\[
t_{\text{cs}} \approx \frac{2m_p}{3m_e} \frac{1}{\Delta T_i - T_e} f(\theta_\varepsilon, \theta_i) \frac{h}{cT} = \Psi \frac{h}{cT}
\]

where \(f(\theta_\varepsilon, \theta_i)\) is the function of \(\theta_\varepsilon, \theta_i\) in eqn. (8) and \(t_{\text{cs}}\) is expressed as a fraction \(\Psi\) of the time scale on which photons escape \((h/c)\) a dissipation region. In eqn. (12) \(\Psi \approx 1/10\) and almost invariant of \(T_i\) for \(T_i \gtrsim 10^{12}\) K (see Fig. 2), so that for a typical \(\tau \lesssim 0.4, t_{\text{cs}} \lesssim h/c\), consistent with \(t_{\text{cs}}/t_c \sim c/h \sim t_e \varepsilon_\gamma/\varepsilon_e\) equal to a few. Variability constraints are thus consistent with ion-electron Coulomb collisions being the dominant process for electron heating.

### 5 DISCUSSION

We have demonstrated that an electron temperature of a few \(\times 10^{9}\)K, determined from interpretation of X-ray observations, can be very naturally explained by an optically thin two-temperature plasma corona model with \(T_i \gg T_e\) if Coulomb scattering is the dominant source of electron heating. If magnetic dissipation in the corona preferentially energizes the ions, then \(T_i\) will naturally exceed \(T_e\). The efficient cooling of electrons then leads a \(T_e\) being locked at the required values. The result of \(T_e \sim 10^{9}\)K is robust because \(T_e\) is relatively insensitive to \(T_i\) when the latter exceeds \(10^{12}\)K.

We have also shown that the time scale for electron heating by Coulomb coupling is consistent with the constraints required by observations of fast variability in AGN. Though we have considered primarily thermalized ions, the actual dissipation may lead to non-thermal ion populations. This would not likely effect \(t_{\text{cs}}\) strongly, but the effect should be considered in future work. Regardless of the proton distribution, the electron distribution would still become Maxwellian.

 Pronounced variability suggests that the corona likely dissipates in large flares, rather than in a large number of small dissipation sites. Since the typically observed variabilities range between \(100 - 1000\)sec and not much shorter, the effective aspect ratio of the dissipation region can be constrained to satisfy \(h/\tau \sim 100/1000 = 1/10\). We can speculate that a small number of operating current sheets would mean that this aspect ratio characterizes the dimensions of a typical reconnection region itself. Such a thick aspect ratio might suggest that Petschek rather than Sweet-Parker type reconnection is occurring. In any case, the physics of dissipation could in principle constrain this value. Future work is needed to fully address the dissipation physics.

The model predicts that the proton energy probably dominates the magnetic energy density, but both are much larger than the photon and electron energy densities. The electron energy density is less than the photon energy density by a factor of the cooling time over the light crossing time. These differences in energy contents highlight the fact that balancing the rates of energy transfer does not necessarily imply equipartition between any energy densities.
Throughout this paper, we have assumed that the only coupling between ions and electrons is via Coulomb collisions which leads to a two-temperature plasma. However, Begelman & Chiueh (1989, hereafter BC) have discussed a plasma instability which may lead to thermal coupling of ions and electrons in a two-temperature flow ($T_i > T_e$) on a time scale shorter than the Coulomb coupling time. The instability may grow in regions with a high level of small scale MHD turbulence that have a plasma $\beta$ parameter $\beta \equiv 8\pi nkT_i / B^2 \gtrsim 1$. We can estimate the rate of electron heating via the BC mechanism using the specific parameters which characterise our solution and compare it to the rate of energy transfer via Coulomb collisions (eqn. 8). For simplicity we choose $T_i = 10^{12}$. According to BC, the rate of heating of electrons when their instability is fully developed is given by

$$\left( \frac{d\varepsilon}{dt} \right)_{\text{BC}} \sim f \frac{m_i v_i^3}{\lambda_D i} \left( \frac{v_A}{c} \right)^2 \left( \frac{\lambda_{c,i}}{L} \right)^9/2,$$

(10)

where $f$ is the filling factor of the plasma where the strong turbulence needed for the instability is present, $m_i$ is the ion mass, $v_i$ is the ion thermal velocity, $L$ is the local density gradient scale and $\lambda_D i$ is the ion Debye length and $\lambda_{c,i}$ its respective gyroradius (where $\lambda_{c,i} = (\beta/\theta_i) \lambda_D i$). We find that for $v_A/c \sim 1/5$ and for $L$ satisfying the condition for the BC instability to operate (eqn. 5.6 in BC paper) with our choice of $\beta$, we have that

$$\frac{(d\varepsilon/dt)_e}{(d\varepsilon/dt)_i} \lesssim 0.3 f$$

(11)

This shows that even in the occurrence of the BC instability, the Coulomb heating usually dominates for the choice of parameters of interest in our solution.

Finally, since the coronal plasma discussed herein has very high ion temperatures, e.g. $T_i \gtrsim 10^{12}$ we expect collisions between protons in the high energy tail of the distribution to produce pions $\pi^0$. These would then decay into high energy $\sim 70$ MeV $\gamma$-rays. Such $\gamma$-ray emission would constitute a testable feature for this model. Power-law or Maxwellian proton (or combination of the two) distributions could be distinguished by the different $\gamma$-ray spectra they predict (Dermer 1986).

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REFERENCES

Achterberg, A., 1987, in Astrophysical Jets and Their Engines, ed. W. Kundt, NATO series C, vol 208, (Dordrecht: Reidel).

Begelman, M.C., 1990, in Accretion Disks and Magnetic Fields in Astrophysics, ed. G. Belvedere (Dordrecht: Kluwer), p. 19.

Begelman M.C., Chiueh T., 1989, ApJ, 332, 872

Blackman, E.G. & Field, G.B., 1994, Phys. Rev. Lett., 73, 3097.

* This process cools the protons without heating the electrons and for ion-temperature of interest it would not cool the ions faster than Coulomb collisions.
