On the Particle Heating and Acceleration in Black Hole Accretion Systems

H. Li\textsuperscript{1}, S. A. Colgate\textsuperscript{1}, M. Kusunose\textsuperscript{2} and R.V.E. Lovelace\textsuperscript{3}

Abstract. The lack of our knowledge on how angular momentum is transported in accretion disks around black holes has prevented us from fully understanding their high energy emissions. We briefly highlight some theoretical models, emphasizing the energy flow and electron energization processes. More questions and uncertainties are raised from a plasma physics point of view.

1. Introduction

Figure 1 shows three (roughly) contemporaneous broad band high energy emission spectra from three galactic black hole candidates (GBHCs; Grove et al. 1998). Although it is conventional to interpret the soft black-body-like component below $\sim 10$ keV as coming from an optically thick Shakura-Sunyaev (SS) disk, the origin of the hard X-ray continuum (and its extension into soft X-rays during the low-hard state) is a constant source of debate. Extracting a physically sensible model through a maze of high quality spectral and timing data on these systems remains a great challenge.

Recently, there seems to be a renewed interest in understanding particle heating/acceleration in accretion disks. We attribute this to the observations of: possible $>0.5$ MeV emissions from Cyg X-1 and GRO J0422; the powerlaw component of GRO J1655 extending to at least 800 keV without a cutoff (Tomsick et al. 1998); and relativistic radio jets from sources like GRO J1655 and GRS 1915. Furthermore, the clearly laid-out physical requirements of ADAF models (which have enjoyed much success, see Narayan et al. 1998 for a review) also prompted further discussions on particle heating.

In this review we will mostly discuss a few models for the so-called low-hard state where the spectrum ($\nu F_\nu$) is peaking around 100-200 keV. We apologize for not able to cover all the models (see Liang 1998 for a recent extensive review). The powerlaw tail that seems to extend beyond 500 keV during the soft-high state also begs explanation, though the total energy contained in this tail is perhaps $<10\%$ of the total emission, so we will place less emphasis on them. We

\begin{flushleft}
\textsuperscript{1}Theoretical Astrophysics, T-6, MS B288, Los Alamos National Laboratory, Los Alamos, NM 87545, hli@lanl.gov, colgate@lanl.gov
\textsuperscript{2}Department of Physics, School of Science, Kwansei Gakuin University, Nishinomiya 662-8501, Japan, kusunose@kwansei.ac.jp
\textsuperscript{3}Department of Astronomy, Cornell University, Ithaca, NY 14853, rvl1@cornell.edu
\end{flushleft}
will focus on the electron energization processes of these theoretical models. We will not discuss any detailed spectral and temporal analyses (see other articles in this volume). Even so, we quickly realized that writing on this topic is a very difficult task because we find many questions and confusions with no clear and definite answers.

2. Some Models for the Origin of Hard X-rays and Gamma-rays

In all the models discussed here, the physics of angular momentum transport (or “α” viscosity) during accretion is not well understood. As a direct consequence, unfortunately, modeling energy dissipation in accretion disks has many ad hoc elements. Quite generally, the matter (surface density Σ) in accretion disk is evolved as (taken from Papaloizou & Lin 1995)

$$\frac{\partial \Sigma}{\partial t} - \frac{1}{r} \frac{\partial}{\partial r} \left[ F_1 + F_2 + F_3 \right] - S_\Sigma = 0$$  \hspace{1cm} (1)

where $F_1 \propto \partial((\nu)\Sigma r^{1/2})/\partial r$ is the local viscous transport with viscosity $\langle \nu \rangle$ (i.e., the standard α−disk viscosity or from MHD turbulence by Balbus & Hawley 1991, 1998); $F_2 \propto S_\Sigma J$ is the advective loss with $J$ being the angular momentum carried by the source/sink ($S_\Sigma$) material (i.e., magnetic flux and/or winds, Blandford & Payne 1982); $F_3 \propto \Lambda$ is the external perturbation (i.e., tidal interactions).
Three models (or their variants) are usually employed for explaining the high energy emissions, namely, the SS model, the SLE model (Shapiro et al. 1976), and the ADAF model. All of them use the local viscous transport prescription (the $F_1$ term) and the energy is also dissipated locally at the disk. In SS model disk is optically thick and geometrically thin, and the plasma is also highly collisional. The heat deposited from transporting angular momentum is successfully radiated away so that disk remains thin ($H \ll R$). In SLE and ADAF models, however, an inner, hot ($T_e \sim 100$ keV), optically thin ($\tau \leq 1$) and two-temperature ($T_i \gg T_e$) region is postulated. This region is then cooled via various radiation processes, such as thermal Compton scattering and Synchrotron.

The arguments for the existence of this hot, optically thin region might be summarized as follows: if local viscous energy dissipation only heats protons, and if there is only Coulomb coupling between electrons and protons, then when the energy input rate is high enough, the system will become unstable if the cooling via radiation is not quick enough, so the plasma has to expand and become optically thin. Here, we want to emphasize that the accreting plasma, during this transition from an optically thick, thin disk to an optically thin, quasi-spherical state, has also changed from highly collisional to essentially collisionless. This brings up several immediate questions which are related to the above “if”.

3. Open Questions

3.1. Will local viscous energy dissipation only heat protons?

Bisnovatyi-Kogan & Lovelace (1997) first discussed this issue and argued that dissipation in such a magnetized collisionless plasma predominantly heats the electrons owing to reconnection of the random magnetic field. On the other hand, Quataert (1998) and Gruzinov (1998) have argued that conditions for ADAF could be true in the high $\beta = P_{\text{plasma}}/P_{\text{magnetic}} \geq 5$ limit by calculating the linear damping rates of short wavelength modes in a hot (but nonrelativistic) plasma, in an (implicit) almost uniform magnetic field. Note that even though MHD turbulence phenomenology was used in both papers, the damping rates are valid in the linear regime for plasma waves only (see below for further discussion). But these calculations perhaps are not answering the question of how to form the optically thin region in the first place because they are damping rates in the collisionless limit. Instead, one perhaps might first evaluate the energy dissipation processes (with an understanding of $\alpha$ viscosity) in the collisional limit which is the physical state initially. These collisions ensure thermal electron and proton distributions and efficient energy exchange between them, especially at the so-called transition radius in ADAF ($10^3 - 10^4 r_s$).

If one uses Balbus-Hawley instability (see also Velikov 1959 and Chandrasekhar 1981) as the origin of the viscosity in the disk, then the gravitational energy is mostly released in large scale (longest wavelength of the magnetic field changes) and this energy will amplify the field first (instead of going into heating the particles). Once the nonlinear saturation is reached (say with magnetic energy density being 10% of the kinetic energy density of the shear flow), we are actually faced with two possibilities, namely, whether the magnetic fields will be expelled (or escape) from the disk, or they will have to dissipate locally in the
disk. Bisnovatyi-Kogan & Lovelace (1997) argued for the second possibility (but see Blackman 1998). Since we know that both the fluid and magnetic Reynolds numbers are exceedingly large in these flows, any “classical” viscous and ohmic dissipations will happen on timescales longer than the age of the universe, thus efficient magnetic reconnection has been sought as the primary candidate for energy dissipation in the disk. They further argued that current-driven instabilities in this turbulent plasma will give rise to large local $E_\parallel$, which mostly accelerate electrons. Thus, up to half of the magnetic energy input goes directly to electrons and is subsequently radiated away, and the disk will always stay thin and optically thick. The uncertainties in these arguments are nevertheless quite large since we don’t fully understand MHD turbulence, let alone its dissipation via kinetic effects. For example, it is unclear whether such reconnection sites are populated throughout the plasma so that most fluid elements encounter such regions. There has been some detailed numerical simulations with magnetic Reynolds number up to 1000 (Ambrosiano et al. 1988) in which test particles are observed to get accelerated by the induced small scale electric fields associated with reconnection sites in turbulent MHD flows. If indeed the magnetic energy dissipation is through accelerating particles by the induced electric fields (this is a big if), since electrons are the current carriers, it is hard to imagine that protons receive most of the energy.

3.2. Is there any collective process that could ensure efficient energy exchange between protons and electrons besides Coulomb?

Putting aside the uncertainties discussed above, if there is indeed an optically thin, hot, two-temperature plasma region, a pertinent question is how much energy electrons can get. This question is, unfortunately, ill-fated again because we do not know how to formulate the problem. Another way to look at it is how to identify the free energy, since most plasma instabilities require a good knowledge of the free energy as determined by the system configuration. For example, is there a relative drift between protons and electrons and can fast electrons be regarded as a beam to an Maxwellian proton core distribution? Is there temperature anisotropy parallel and perpendicular to background magnetic fields, etc.? Begelman & Chiueh (1988) have studied some plasma instabilities in detail and found plausible ways of transferring energy from ions to electrons, under the conditions that a substantial level of MHD turbulence will give a large enough proton density gradient (or curvature drifts) so that proton drift velocity can be large enough to drive certain modes unstable. The fluctuating electric field parallel to the magnetic field will then accelerate electrons. The applicability of this instability is again hampered by our lack of knowledge of the presumed MHD turbulence. Narayan & Yi (1995) argued that this mechanism does not work well in ADAF.

A conceptual difficulty is that the typical modes excited by protons (having most of the energy) are below the proton gyrofrequency $\Omega_{ci}$. This makes resonance with the electrons difficult. But a possible avenue is to have protons excite (almost) perpendicular modes (i.e., high $k_\perp$ and very small $k_\parallel$). Then the resonant conditions for electrons to resonate with these waves are easier to satisfy. More work is needed to explore these possibilities.
3.3. Could accretion disk have a magnetically dominated, hot corona like our Sun?

The formation of a “structured corona” was first proposed by Galeev et al. (1979). In this model a radial quadrupole field is wound up by differential rotation into an enhanced toroidal field. Then the helicity of the presumed convective “turbulence” converts a fraction of the toroidal flux back into poloidal field and hence produces an exponentiating dynamo that saturates by back reaction. This is the classical $\alpha - \Omega$ dynamo although not identified as same in the paper. Furthermore the saturation or back reaction limit of this disk dynamo is assumed to be the random loops of flux characteristic of the solar surface.

One important step in the above model is the requirement of vertical (thermal) convection in the $\{R,z\}$ plane. The convective motion may be driven by heat released at or near the mid-plane. Lin et al. (1993) have shown that under specific conditions of opacity and equation of state that convective instability should occur both linearly and nonlinearly, thus leading to large amplitude cells. However, the convective cells are highly constrained radially. The problems of restrictive initial conditions and the restrictive cell geometry leads one to conclude that this is not the universal mechanism needed to explain accretion disks. Colgate & Petschek (1986) showed that to drive convective cells whose displacement radially is of the order of the disk height $h$, (unrealistic) efficiency of the energy flow (a Carnot cycle of $\sim 100\%$ efficiency) is necessary to drive these convective eddies, and the cells created are also highly restrictive, tall but narrow radially (i.e., similar as Lin et al.). Thus the existence of strong convective turbulence is doubtful. The result of this lack of convective turbulence with rising plumes is to negate the origin of the helicity invoked in the structured corona model.

4. Electron Energization

Besides the possible role of magnetic reconnection in accelerating electrons which is observed in the solar corona (Tsuneta 1996), there are more standard processes which involve wave-particle interactions (see Kuijpers and Melrose 1996). Shock acceleration is not considered here. We give a quick review of the electron energization by plasma waves and turbulence.

4.1. Particle heating/acceleration – linear and quasilinear theory

Linear Vlasov equation is usually used to describe the collisionless plasma, which is a good approximation of astrophysical plasmas. Linearization of the Vlasov equation yields various dispersion relations $\omega = \omega(k)$ which describe how the system will respond to small electrostatic and electromagnetic perturbations. Since the field energy of low frequency fluctuations (i.e., $\omega < \Omega_{cp}$) is predominantly magnetic, particles generally experience strong pitch-angle scattering before they can be energized. Of fundamental importance is the wave-particle resonance, that is, given an electromagnetic fluctuation of frequency $\omega$ and wavevector $k$, a charged particle ($q,m$) is considered to be in resonance with this fluctuation when
\[ \omega - k_{\parallel}v_{\parallel} - \ell \Omega_0/\gamma = 0, \quad \ell = 0, \pm 1, \pm 2, \ldots \] (1)

where the nonrelativistic gyrofrequency \( \Omega_0 = |q|B_0/mc \), and \( v_{\parallel} \) and \( \gamma \) are the particle’s parallel velocity and Lorentz factor, respectively. When the harmonic number \( \ell = 0 \), the resonance is referred to as the Landau or Cherenkov resonance, and implies that the particle speed along the magnetic field matches the speed of the parallel wave electric or magnetic field. If \( |\ell| > 0 \), the process is called gyroresonance, and there is a matching between the wave transverse electric field and the cyclotron motion of the particle. The sign of \( \ell \) depends upon the transverse polarization of the wave and the sign of \( q \): if the transverse wave electric field and the particle rotate in the same sense about \( B_0 \) in the plasma frame, then \( \ell \) is positive. In most settings, only \( \ell = \pm 1 \) is of importance.

The key quantities are the plasma beta \( \beta = n(T_i + T_e)/(B^2/8\pi) \) factor and the temperature ratio \( T_e/T_i \). Furthermore, we have:

- **Linear theory.** The linear theory of plasma waves and instabilities is often reduced to a linear dispersion equation with a complex \( \omega \), whose imaginary part gives the growth or damping of certain modes. Recent studies by Gruzinov (1998) and Quataert (1998) belong to this case. The usual candidates for wave-particle resonances are: for \( \ell = 0 \), the transit time damping (TTD) for particle with oblique fast magnetosonic waves and Landau damping (LD) with kinetic Alfvén waves; for \( |\ell| \geq 1 \), gyroresonances between the proton/Alfvén wave and the electron/whistler wave.

- **Quasilinear theory.** A detailed physical understanding of pitch-angle scattering and stochastic acceleration is beautifully presented in Karimabadi et al. (1992), using nonlinear orbit theory with the Hamiltonian formalism. In the presence of a continuum of plasma waves, the number of resonances between the particle and waves is greatly increased to a point that the trapping width associated with one particular resonance can overlap with neighboring resonances, thus allowing particles “jump” from one resonance to another. As particles sample different resonances, they gain energy in a “ladder-climbing” fashion. Hence the description stochastic acceleration. This approach has been adopted in several studies on electron acceleration by fast-mode waves and whistler waves in accretion disk (Li et al. 1996, Li & Miller 1997). We typically find that the electron distribution is hybrid with a nonthermal tail, which is responsible for the production of \( > 500 \) keV emissions in several GBHCs. We have built a computer code which solves 3 coupled, time-dependent kinetic equations for particles, photons and waves, respectively. Namely,

\[
\frac{\partial N_e}{\partial t} = - \frac{\partial}{\partial E} \left\{ \left[ \frac{dE}{dt} + \frac{dE}{dt}_{\text{loss}} \right] N_e \right\} + \frac{1}{2} \frac{\partial^2}{\partial E^2} [(D + D_c)N_e] \quad (2)
\]

\[
\frac{\partial W_T}{\partial t} = \frac{\partial}{\partial k} \left[ k^2 D \frac{\partial}{\partial k} \left( k^{-2} W_T \right) \right] - \gamma W_T + Q_W \delta(k - k_0) \quad (3)
\]

\[
\frac{\partial n_{\text{ph}}(\varepsilon)}{\partial t} = -n_{\text{ph}}(\varepsilon) \int dE N_e(E) R(\varepsilon, E) + \int \int d\varepsilon' dE P(\varepsilon; \varepsilon', E) n_{\text{ph}}(\varepsilon') N_e(E) + \dot{n}_{\text{ext}}(\varepsilon) + \dot{n}_{\text{emis}}(\varepsilon) - \dot{n}_{\text{abs}}(\varepsilon) - \frac{n_{\text{ph}}(\varepsilon)}{t_{\text{esc}}} \quad (4)
\]
The particle distribution can be arbitrary. This allows us to determine from all the interactions whether the distribution is thermal or nonthermal. Pair production is not included so far. The Coulomb terms are also implemented for arbitrary particle distributions. Accurate Compton scattering is treated as a scattering matrix (Coppi 1992) with the full cross section. The Cyclo-Syn. process is calculated according to Robinson & Melrose (1984) which enters both as a cooling and heating term (Ghisellini et al. 1988). Syn-self absorption is also included. The radiation part of the kinetic code is tested against Monte Carlo simulations (Kusunose, Li, & Coppi 1998) and is found to be very good for $\tau \leq 3$ and for both thermal and nonthermal electron distributions.

![Figure 2](image-url)

**Figure 2.** Time evolution (from $t = 0 - 10 R/c$) of particle distribution (top), MHD wave spectral density (middle) and photon flux (bottom). The initial particle $T_e = 100$ keV and $\tau = 0.1, 0.5, 1$, respectively. The soft photons are injected with $T_s = 1$ keV and the compactness $\ell_s = 10$. The size is $\sim 30 r_g$ and $M = 7 M_\odot$. The final distributions are indicated by thick curves and dashed lines are initial distributions. Deviations from Maxwellian are obtained as MHD waves cascade to higher $k$, accelerating electrons out of the thermal bath. Cyclo-Syn. (with self-absorption) and Compton are the radiation processes considered here.
plication to optically thin environment in mind. The plasma density $n$ is varied from $\tau = 0.1 - 1$. At early times, the particle distribution softens first as shown in upper panels, due to that waves have not fully cascaded (i.e., small $\langle k \rangle$ as shown in middle panel), and losses dominate at high energies. As waves cascade over the inertial range, $\langle k \rangle$ quickly grows to a level that acceleration overcomes all losses, electrons are then energized out of the thermal background and the nonthermal hard tail forms. The photon spectra indicates that gamma-rays can be produced when $\tau < 0.5$. Furthermore, note that the nonthermal tails start to develop at $E/m_e c^2 \sim 0.13$ (corresponding to $v_A/c = 0.46$), this nicely confirms the fact that only particles with $v > v_A$ can be accelerated.

4.2. MHD turbulence, are they an ensemble of waves?

The above described calculations, both linear and quasilinear, can be broadly regarded as “dissipation” in a general MHD turbulence theory. Finding a dynamical model that might adequately describe the evolution of magnetic fluctuations (such as equation (3) above) is at best phenomenological. In the dissipation range, the physics of the couplings that connect fluid and kinetic scales is not understood at all.

A critical assumption that is employed in all the kinetic calculations is that the magnetic fluctuations that cascade from large scales to small scales could be regarded as an ensemble of kinetic waves with a well-defined dispersion relation to describe them. This view is by no means proven, though it allows us to get an estimate of the particle heating rate since the kinetic theory is significantly more advanced (see an application of such an approach to the interplanetary magnetic field dissipation range, Leamon et al. 1998). On the other hand, the dynamics of MHD turbulence has been studied using statistical theories and simulations (e.g., Kraichnan & Montgomery 1980; Shebalin et al. 1983; Matthaeus & Lamkin 1986), and has never been convincingly presented or developed within a normal-mode, perturbation-type of framework. A further complication is that most (MHD) turbulence theory is based on the incompressible fluid model, how it will “carry-over” to compressible astrophysical flow is still an open question.

4.3. MHD turbulence Truncation

Recently, the assumption of a cascade to smaller scales of MHD turbulence is criticized in dynamo theory. It has been argued that both the more rapid folding of magnetic flux as well as the smaller energy density at small scale ensures rapid saturation or back reaction by the field stress, immobilizing the small scale fluid motions expected from the Kolmogorov spectrum. This will truncate the turbulent spectrum at the back reaction scale, initially the smallest and progressively reaching the largest. Since the energy input to the turbulence is assumed to be primarily at the largest scale, this leaves one with negligible power at the small scale (Cattaneo & Vainshtein 1991; Kulsrud & Anderson 1992; Gruzinov & Diamond 1994; Cattaneo 1994). The remaining largest scale is that of the disk itself.

In general all particle instabilities presumably leading to particle heating require large local gradients in some aspect of their phase space, i.e. temperature, density, and velocities, etc. Furthermore, all the free gravitational energy must flow through these gradients. This requirement, however, will not be met if
the small scale turbulent motions are strongly damped by the back reaction of the field itself. We therefore look for a solution to this paradox in large scale magnetic structures.

5. A Sketch View

Here, we outline some plausible physical pictures about what might be happening in an accretion flow. Most these are ideas that have not been thoroughly investigated. It is also clear that there are obvious gaps which need to be filled with rigorous calculations.

5.1. Hydrodynamic transport and high-soft state

Many investigations have sought a linear instability deriving energy from the Keplerian flow to produce a growing mode leading, in the non-linear limit, to turbulence. The Papaloizou & Pringle instability (1984) seems to be the most studied instability but its relevance to Keplerian accretion disks has been questioned (Balbus & Hawley 1998).

Recently we have identified a linear instability in Keplerian disk leading to Rossby waves and presumably Rossby vortices in the nonlinear limit (Lovelace et al. 1998). This instability grows most effectively from a large radial gradient in entropy. It has the advantage that the nonlinear limit consists of co-planar, co-rotating vortices (Nelson et al. in preparation) that require only a radial gradient, not vertical gradient of entropy. The radial gradient, we believe, is astrophysically reasonable because all disks are presumably fed by matter at some outer radius by, for example, Roche-lobe overflow in low-mass X-ray binaries. If there is no angular momentum removal mechanism, the matter will accumulate until it builds up enough to trap heat, and variations in entropy would then render the onset of the above instability. We do not, however, expect this instability to lead to turbulence in the usual sense of convective turbulence. An ensemble of co-planar vortices does not lead to significant vertical flow as compared to the usual picture of convective turbulence where buoyant plumes would convect heat released at the the mid-plane to the disk surface.

We expect the angular momentum transport is done via nonlinear interactions of these vortices with the background flow, but this has to be addressed by extensive hydro simulations. The heat flow derived from the “viscosity” of the ensemble of Rossby vortices must be removed by radiation flow. We expect this not to be a problem because the radiation thickness of the disk, $\tau$, is small enough such that the effective diffusion velocity, $v_{\text{diff}} \approx c/(3\tau) >> v_\phi$. Under these conditions, the disk solution will be essentially the same as the SS disk. Thus, this picture might be applied to the high-soft state of GBHCs. Relatively speaking, magnetic fields do not play a major role during this state but some nonthermal processes (such as a weak magnetic outflow) might be responsible for the powerlaw component from 20 keV - 1 MeV.

5.2. Role of large scale magnetic fields and low-hard state

As pointed out by Blandford & Payne (1982), large scale magnetic fields can also be very effective in removing the disk angular momentum. These large
Figure 3. A linearly unstable mode in a 2D Keplerian disk with an entropy bump initially located at \( r = 3 \). Shown is the amplitude (± means moving out/in) of the radial velocity \( v_r \) which is zero initially. This is an \( m = 5 \) mode. The unstable mode is “trapped” at the entropy bump. Since the flow is nonbarotropic, additional vorticity \( \nabla \times \mathbf{v} \) is also produced around \( r = 3 \).

scale magnetic outflows could be a hydromagnetic wind (Blandford & Payne 1982), or it might be a nearly force-free helix (Poynting flux) with very little matter as discussed by Lovelace et al. (1987, 1997).

The accreting plasma from, say a companion star, is likely to be magnetized. In the advection of this flux with the mass flow, there will necessarily be a convergence and strengthening of the field. In the region where an \( \alpha \)-viscosity prevails and the field acts as a passive marker of the flow, there will be both advection and diffusion. The diffusion radially outwards depends upon probably the same diffusion coefficient which allows the diffusion of angular momentum. Hence there will be a unique relationship between advection inwards and diffusion outwards leading to the relationship, \( B_z \propto r^{-3/2} \) (see also Bisnovatyi-Kogan & Lovelace 1997 in which they argued \( B_r \propto r^{-2} \)).

If the initial field strength advected with the mass flow at the outer disk radius is large enough, then the field energy density could become comparable to the Keplerian stress at a certain radius. Bisnovatyi-Kogan & Lovelace (1997) argued that magnetic flux then has to be destroyed at the disk via reconnection. Alternatively, instead of destroying the flux, magnetic fields (presumably tied to the companion star) could be twisted such as they will remove the angular momentum of the flow and take away the released gravitational energy. So the energy dissipation (into radiation) might not be at the disk at all. The reconnection dissipation of the current supporting the torsion of the magnetic field will perhaps lead to the non-thermal emission of GBHCs. In fact, there
is ample evidence in AGNs that perhaps most of the energy release is in the outflow/jet. Such a picture could also apply to GBHCs with the hard X-ray to gamma-ray emissions produced via nonthermal processes (such as Syn. or SSC) in the magnetized outflow away from the disk.

If the initial field strength advected with the mass flow at the outer disk radius is weaker, then the amplified magnetic field (by $r^{-3/2}$) may never be greater than the Keplerian stress, thus is not the dominant channel of angular momentum transport. However, there will always be nonthermal energy release in the twisted magnetic field, which could be the powerlaw tail during the high-soft state.

**Acknowledgments.** HL wish to thank the meeting organizers for their kind invitation and financial support. Part of the work on quasilinear wave-particle interactions was done together with J. Miller. We thank J. Finn and E. Liang for many useful discussions. HL gratefully acknowledges the support of an Oppenheimer Fellowship at LANL.

**References**

Ambrosiano, J. et al. 1988, JGR, 93, 14383
Balbus, S.A. & Hawley, J.F. 1991, ApJ, 376, 214
Balbus, S.A. & Hawley, J.F. 1998, Rev. Mod. Phys., 70, 1
Begelman, M.C. & Chiueh, T. 1988, ApJ, 332, 872
Bisnovatyi-Kogan, G.S. & Lovelace R.V.E. 1997, ApJ, 486, L43
Blackman, E.G. 1998, PRL, in press
Blandford, R.D. & Payne, D.G. 1982, MNRAS, 199, 883
Cattaneo, F. 1994, ApJ, 434, 200
Cattaneo, F. & Vainshtein S.I. 1991, ApJ, 376, L21
Chandrasekhar, S. 1981, Hydrodynamic and Hydromagnetic Stability, New York: Dover
Colgate, S.A. & Petschek, A.G. 1986, Los Alamos Science, No. 13, p61
Coppi, P.S. 1992, MNRAS, 258, 657
Galeev, A.A. et al. 1979, ApJ, 229, 318
Ghisellini, G. et al. 1988, ApJ, 334, L5
Grove, J.E. et al. 1998, ApJ, 500, 899
Gruzinov, A.V. 1998, ApJ, 501, 787
Gruzinov, A.V. & Diamond, P.H. 1994, PRL, 72, 1651
Karimabadi, H. et al. 1992, JGR, 97, 13853
Kraichnan, R. & Montegomery, D. 1980, Rep. Prog. Phys., 43, 547
Kuijpers, J. & Melrose, D.B. 1996, in Plasma Astrophysics, eds. C. Chiuderi, G. Einaudi, Springer-Verlag
Kulsrud, R.M. & Anderson S.W. 1992, ApJ, 396, 606
Kusunose, M., Li, H. & Coppi, P.S. 1998, ApJ, submitted
Leamon, R.J. et al. 1998, JGR, 103, 4775
Li, H. et al. 1996, ApJ, 460, L29
Li, H. & Miller, J.A. 1997, ApJ, 478, L67
Liang, E.P. 1998, Phys. Rep. 302, 67
Lin, D.N.C. et al. 1993, ApJ, 416, 689
Lovelace, R.V.E. et al. 1987, ApJ, 315, 504
Lovelace, R.V.E. et al. 1997, ApJ, 484, 628
Lovelace, R.V.E. et al. 1998, ApJ, in press
Matthaeus, W.H. & Lamkin, S.L. 1986, Phys. Fluids, 29, 2513
Narayan, R. & Yi, I. 1995, ApJ, 452, 710
Narayan, R. et al. 1998, astro-ph/9803141
Nelson, A.F. et al. 1999, in preparation
Papaloizou, J.C.B. & Lin, D.N.C. 1995, ARAA, 33, 505
Papaloizou, J.C.B. & Pringle, J.E. 1984, MNRAS, 208, 721
Robinson, P.A. & Melrose, D.B. 1984, Aust. J. Phys., 37, 675
Shakura, N.I. & Sunyaev, R.A. 1973, A& A 24, 337
Shapiro, S. et al. 1976, ApJ, 204, 187
Shebalin, J. et al. 1983, J. Plasma Phys., 29, 525
Quataert, E. 1998, ApJ, 500, 978
Tomsick, J.A. et al. 1998, ApJ, in press
Tsuneta, S. 1996, ApJ, 456, 840
Velikhov, E.P. 1959, J. Exp. Theo. Phys., 36, 1398