Corona Energy Budget in AGN and GBHCs

Sergei Nayakshin

Department of Physics, the University of Arizona, Tucson, AZ, 85721

Abstract. Recent progress in observations and understanding of spectra of Seyfert Galaxies shows that X-rays are most likely produced by magnetic flares on the surface of the disk, similar to Solar X-ray emission. However, while the model reproduces the shape of the spectrum well, the question of the overall normalization of the X-ray component relative to the bolometric luminosity has not been previously considered. Here we show that, in gas-dominated accretion disks, the magnetic energy transport can indeed power the corona in a way consistent with observations, if magnetic field in the disk is mostly contained to strong magnetic flux tubes. However, in radiation dominated disks radiation diffusion makes the field weak/diffuse, and thus the magnetic energy transport is less efficient, i.e., in such disks the disk intrinsic emission should be the dominant component in the overall spectrum. We compare our findings to observations, and conclude that our theory can account for the often observed “steeper when brighter” behavior in AGN and the hard-soft spectral transitions in GBHCs.

1. Introduction

The two-phase patchy corona-disk model (PCD model hereafter) was suggested by Haardt & Maraschi (1991,1993) & Haardt et al. (1994) to explain the almost unique X-ray spectral index of Seyferts. Via accurate spectral calculations, Stern et al. (1995), Poutanen & Svensson (1996) & Poutanen et al. (1997) showed that the model naturally explains observed distribution of the X-ray spectral indexes, strength of the fluorescent iron line, and the presence of Compton reflection component for Seyfert Galaxies (see also Svensson 1996).

The basic physics of the PCD model is well understood qualitatively (e.g., Haardt et al. 1994, Galeev et al. 1979, Nayakshin 1998b), because it is adapted after Solar magnetic flares (e.g., Parker 1979, Priest 1982, Tsuneta 1996, Tajima & Shibata 1997). Namely, turbulent accretion disks are expected to produce magnetic fields very efficiently, since Keplerian accretion disks are differentially rotating strongly ionized plasmas. The magnetic fields tend to concentrate into magnetic flux tubes – regions of enhanced magnetic field, with magnetic pressure comparable to the ambient gas pressure (Parker 1979, Chapter 10). Because the flux tubes immersed in the fluid must be in pressure equilibrium with that fluid, the gas density inside the tubes is lower than the ambient gas density. Parker (1955) has shown that the tubes are therefore buoyant. Thus, these tubes are expelled out of the disk. Once above the disk, the magnetic fields of these tubes
can reconnect – transfer their energy to particles trapped within the tubes, producing hot active regions, or “patches” in the patchy corona model. These active regions then produce X-rays by inverse Compton upscatter ing of the disk radiation and by other emission mechanisms.

The broad-band spectrum of accretion disk in this picture consists of basically three components: (1) the X-rays produced by magnetic flares in the disk corona; (2) radiation due to reprocessing of these X-rays in the disk below the flares (which we think is the origin of the Big Blue Bump of Seyfert 1s [BBB hereafter; Nayakshin & Melia 1997, Nayakshin 1998b]); (3) the internal optically thick disk emission. Following Svensson & Zdziarski (1994), let us parameterize the fraction of the overall disk power released through the flares as \( f < 1 \). If the disk bolometric luminosity is \( L \), the intrinsic disk emission then accounts for luminosity \((1 - f) L \). Since roughly half of the X-rays produced by the flares are directed back to the disk (e.g., Haardt et al. 1994), the observed X-rays and the BBB should both produce luminosities \( L_x \approx L_{\text{bbb}} \approx (1/2)fL \). Here we attempt to address the magnitude of \( f \) in the context of the PCD model, and compare it to observations of both AGN and GBHCs.

### 2. Energy Transfer by Buoyant Magnetic Fields

Let us estimate the magnetic energy transport due to rising magnetic flux tubes and compare it with the total disk emission (energy liberated per unit area of the disk). The latter can be shown to be (e.g., Frank et al. 1992, Chapters 4,5) \( F_{\text{tot}} \approx \alpha c_s P_{\text{tot}} \), where \( \alpha \) is the Shakura-Sunyaev viscosity prescription, \( c_s \) is the sound speed in the mid-plane of the disk, and \( P_{\text{tot}} \) is the total (gas plus radiation) pressure in the disk. The time averaged magnetic energy flux is simply given by \( F_{\text{mag}} \approx v_b \langle P_{\text{mag}} \rangle \), where \( v_b \) is the average buoyant rise velocity, which cannot be larger than \( c_s \), and \( \langle P_{\text{mag}} \rangle \) is the volume average of the magnetic field pressure. Thus, the importance of the magnetic energy transport depends on the ratio \( \langle P_{\text{mag}} \rangle / P_{\text{tot}} \). The viscosity parameter was defined by Shakura & Sunyaev (1973) to be

\[
\alpha \simeq \nu_t + \frac{\langle P_{\text{mag}} \rangle}{P_{\text{tot}}},
\]

where the first term, \( \nu_t \) is the turbulent viscosity. Thus, \( \langle P_{\text{mag}} \rangle \lesssim \alpha P_{\text{tot}} \), which yields

\[
f \equiv \frac{F_{\text{mag}}}{F_{\text{tot}}} \lesssim \frac{v_b}{c_s} \lesssim 1
\]

The buoyant rise velocity depends on properties of the magnetic flux tubes and the gas around them. For flux tubes of size \( a_0 \ll H \), where \( H \) is the disk height scale, Vishniac (1995a,b) shows that \( v_b \sim \alpha c_s \), which makes \( f \) a small number. Physically, the problem arises because during its rise to the disk surface, magnetic field also contributes to the angular momentum transfer in the disk, which in turn leads to further liberation of energy in the disk. This latter “side” process in fact produces more heat than magnetic buoyancy itself, and so it would appear that the PCD model should always produce more (and much more, if \( \alpha \ll 1 \)) thermal disk emission than the corona can, which is at odds with observations. The strongest constraint here comes from GBHCs in their
hard state, whose spectra contain most of the power in the hard power law, so that $f$ is probably $\approx 0.8$. For Seyferts, the X-ray emission often accounts for a significant fraction of the overall emission, and $f$ as large as 0.5 seems to be needed.

A possible solution to this dilemma lies in the fact that estimate $\alpha \sim \langle P_{\text{mag}} \rangle / P_{\text{tot}}$ is only correct for a diffuse magnetic field, and a similar argument carefully applied to a field localized to strong magnetic flux tubes shows that these fields can contribute considerably less to the disk viscosity and thus the local heating. The reduced local heating would then explain how it is possible for real accretion disks to have $f$ approaching unity. Our point here is that a diffuse and tangled magnetic field of sub-equipartition intensity will be simply carried along with the fluid, i.e., it will take part in the differential rotation of the fluid. By resisting the stretching through magnetic field tension, the field will contribute to viscosity in the usual way. At the same time, a flux tube is an entity of its own, which manifests itself in the fact that the tube can move with respect to the fluid, e.g., be buoyant. Accordingly, the flux tube may avoid the stretching by simply not following these motions of the fluid that try to deform the tube. Thus, the flux tube may contribute to viscosity at a smaller rate than a diffuse field would do.

Suppose for simplicity the shape of the tube is that of a torus with the larger radius $a_0 \lesssim H$ and the smaller radius $a \leq 2a_0$. There is a viscous drag force $D$ on the flux tube in this case, caused by the friction as the fluid flows by the tube:

$$D \simeq C_d \rho v_d^2 a a_0/2$$

(e.g., Parker 1979, §8.7, and references there), where $C_d$ is the dimensionless drag coefficient, $\rho$ is the gas density and $v_d$ is the differential flow velocity, which is $v_d \sim c_s a_0 / H$ for a Keplerian accretion disk. For the flux tube not to be deformed by the drag force, the tube magnetic tension $T$ should exceed this force:

$$T \simeq P_{\text{mag}} 2\pi a^2 > D$$

The ratio of these two forces is

$$\frac{T}{D} \sim \frac{4\pi a}{a_0} C_d^{-1} \frac{P_{\text{mag}}}{P_{\text{tot}}} \left(\frac{H}{a_0}\right)^2$$

where we used $P_{\text{tot}} = \rho c_s^2$. The value appropriate for the drag coefficient in accretion disks is $C_d \sim 1/4$ (following Vishniac 1995, Stella & Rosner 1984, Sakimoto & Coroniti 1989, Parker 1979). Thus, equation(5) asserts that for flux tubes with magnetic field pressure comparable to the equipartition value, and the size $a_0$ smaller than the disk scale height $H$, $T > D$, and thus the tubes cannot be deformed by the flow in this case, and instead are dragged around almost as a solid body. The contribution of the flux tube to the angular momentum transfer is reduced by the same factor $\sim D/T$, since the flux tube is being stretched at a rate slower than implied by the differential flow of the ambient gas around the tube. If all the magnetic field is in the form of strong flux tubes for which the magnetic tension exceeds the drag force, then the limits on the magnetic field volume average become

$$\left(\frac{D}{T} P_{\text{mag}}\right) \lesssim \alpha P_{\text{tot}}$$
We can now estimate the ratio of the magnetic energy flux $F_m$ to the radiation energy flux as

$$\frac{f}{1-f} \simeq \frac{v_b}{c_s} \frac{1+T}{D},$$

which is much easier to reconcile with the magnetic energy flux requested by the two-phase corona-accretion disk model, since now the buoyant rise velocity can be comfortably below its absolute maximum value, i.e., the sound speed $c_s$ and yet provide magnetic energy flux exceeding the radiation flux.

A simple physical analogy here is that of a sail on a ship. When the sail is “on”, the force (due to wind) acting on the sail is many times larger than it is in the case of the sail that is folded in. The amount of this wind-sail interaction clearly depends not on the overall mass of the sail, but on the state of the sail – whether it is open and positioned properly with respect to wind or whether it is rolled in a tube. Similarly, with same volume average magnetic field one gets less or more interaction between differential flow and the field depending on whether the field is uniform in space, or is in strong flux tubes, such that most of the flow simply miss the tubes to interact with them.

3. Radiation Pressure and Properties of a Single Flux Tube

Considering magnetic fields in the previous section, we did not explicitly separate the total pressure $P_{\text{tot}}$ into the radiation pressure $P_{\text{rad}}$ and the gas pressure $P_{\text{gas}}$. Our initial neglect of the radiation pressure dynamical effects is equivalent to the assumption that radiation and particles move together, as one fluid. Such approach is valid as long as scales of interest are much larger than the photon mean free path, since in this case the radiation is essentially “glued” to particles due to the large opacity. In the opposite limit of small opacity, the gas and radiation will behave very differently with respect to magnetic fields. Whereas radiation does not directly interact with magnetic field (i.e., these two components are unaware of each other’s presence), particles in ideal MHD cannot cross magnetic field lines and their motion is constrained to the direction parallel to the field. Let us then compare the time scale for the radiation diffusion into the flux tube with a time scale important for generation and maintenance of strong magnetic flux tubes. The radiation diffusion time scale $t_d$ can be estimated as

$$t_d \sim \frac{a}{c} n'_e \sigma_T a,$$

where $n'_e$ is the particle density inside the flux tube, which we can assume to be of the order of the disk particle density $n_e$, and $\sigma_T$ is the Thomson cross section.

Turbulent motions of the fluid are believed to be the mechanism for the magnetic field amplification (e.g., Vishniac 1995a, b). Let $u_t$ be the typical turbulent velocity, and $\lambda_t$ be the turbulent length scale (corresponding to the largest eddy length scale). The gas executes turbulent motions on the eddy turn over time scale $t_t \equiv \lambda_t/u_t$. This yields

$$\frac{t_{\text{d}}}{t_t} \sim \frac{a^2}{H \lambda_t} \frac{u_t}{c_s} \frac{\tau_d c_s}{c},$$

where $\tau_d$ is the disk Thomson optical depth, and $a \lesssim \lambda_t$ (Vishniac 1995a). Further, in the standard Shakura-Sunyaev viscosity prescription, the turbulent velocity and spatial scale are parameterized by $u_t \lambda_t = \alpha c_s H$. Finally, in the
radiation pressure dominated region of the disk, the standard disk equations lead to $\tau_d c_s / c \simeq \alpha^{-1}$, for arbitrary radii and accretion rate. Therefore, one can see from Equation (8) that the ratio of the diffusion time scale to the turbulent time scale is of the order unity. Moreover, we compared $t_d$ with one eddy turnover time scale, whereas generation of the field comparable with the equipartition value is likely to take much longer time, simply because one turbulent eddy does not carry enough energy (we assume that turbulence is sub-sonic). Due to this, the diffusion of radiation into the flux tubes is much faster than the field generation process.

As a result of the efficient radiation diffusion, the radiation pressure inside a flux tube should be equal to the ambient radiation pressure, which means that the magnetic field pressure of the tube can only be as large as the gas pressure $P_{\text{gas}}$. A simple analogy here is a car tire that is pumped with a gas that easily diffuses in or out. Obviously, difference in the gas pressure inside and outside will be small which will render such a tire useless. In the case of the flux tubes, $P_{\text{mag}} \lesssim P_{\text{gas}} \ll P_{\text{tot}}$ for the radiation-dominated disks means that the tubes are now weak and easily stretched by the differential motions of the disk, and thus $(P_{\text{mag}}) \sim \alpha P_{\text{tot}}$. The tension force is now small compared with the drag force due to differential Keplerian flow, and, according to equation (7), $f/(1 - f) \simeq v_b/c_s < 1$. In physical terms, the X-ray luminosity should always be smaller than the optical- to soft X-ray luminosity in the radiation-dominated accretion disks.

4. Spectral States of AGN and GBHCs

In order for magnetic flares to be in the parameter space of the PCD model, the magnetic fields in the active regions should be sufficiently strong to make the compactness parameter $l \gg 0.01$, and the disk intrinsic flux should be much smaller than the X-ray flux. Since the field is limited to the equipartition value in the disk, one finds $l \propto \dot{m} \alpha^{-1}(P_{\text{mag}}/P_{\text{tot}})$ (e.g., Nayakshin 1998b). Numerically, it seems unlikely that accretion disks dimmer than $\dot{m} \lesssim 10^{-4}$ or so will be able to produce flares of the right compactness. Thus, we believe that to produce the typical hard X-ray spectrum, AGN must accrete above some minimum accretion rate $\dot{m}_d \sim 10^{-4}$.

The transition from the gas to radiation-dominated disks happens at $\dot{m} = \dot{m}_r$:

$$\dot{m}_r = 2.2 \times 10^{-3} (\alpha M_8)^{-1/8} (1 - f)^{-9/8}$$

For accretion rates $\dot{m}_d < \dot{m} < \dot{m}_r$, the magnetic flares satisfy PCD model constraints, and thus their X-ray spectra should be hard, (i.e., typical hard Seyfert spectrum for AGN and typical hard state spectrum for GBHCs). In addition, since magnetic buoyancy can transfer more energy out of the disk than the usual radiation flux, most of the power is contained in the X-ray power law (i.e., $f \simeq 1$) and the emission due to reflection and reprocessing of this radiation in the disk.

As the accretion rate increases above $\dot{m}_r$, the importance of X-ray production by magnetic flares decreases, i.e., the fraction $f$ is decreasing as $\dot{m}$ increases. This means that intrinsic disk emission becomes the dominant feature in the spectrum of an AGN or a GBHC. Note that the expression for $\dot{m}_r$
given by equation (9) depends on \( f \) itself, but since radiation-dominated disks are not effective in transporting energy into the corona, factor \( 1 - f \) should not be too small for such disks. In fact, if we use observational constraints on \( f \) from hard/soft spectral transitions in GBHCs, \( 1 - f \approx 1/2 \) at \( \dot{m} = \dot{m}_r \) (see §5).

Based on theoretical arguments alone, we cannot be certain about what happens to the shape (as opposed to already discussed normalization) of the X-ray spectrum from flares when \( P_{\text{rad}} \gg P_{\text{gas}} \). The uncertainty is present due to our ignorance of the numerical value of \( \alpha \) when radiation pressure exceeds the gas pressure. The standard accretion disk theory in this case is unstable to viscous and thermal perturbations (e.g., Frank et al. 1992), and so the form of viscosity law in radiation dominated disks remains a highly controversial issue. However, if we assume that the viscosity is proportional to the gas pressure only (for motivation see Lightman & Eardley 1974, Stella & Rosner 1984, Sakimoto & Coroniti 1989), the accretion disk is stable, and it turns out that the decrease in the effective value of \( \alpha \) may compensate for the fact that magnetic fields are limited to the gas pressure, so that the flares may still be in the PCD model parameter space. Namely, the effective \( \alpha \) scales as \( \alpha_g P_{\text{gas}} / P_{\text{tot}} \) in this case, where \( \alpha_g \leq 1 \) is a constant. Thus, \( l \propto (\dot{m} / \alpha) (P_{\text{gas}} / P_{\text{tot}}) = (\dot{m} / \alpha_g) \), that is, the compactness parameter does not have to decrease as the disk goes from the gas- to the radiation-dominated regime.

Depending on behavior of \( \alpha \) in the radiation-dominated disks, the X-ray spectrum from magnetic flares will either be unchanged from that of the hard state, or will become steeper because of the extra cooling of the active regions caused by the stronger intrinsic disk emission. The X-ray spectrum will steepen when the disk intrinsic flux will approach the X-ray flux from a flare. This happens for accretion rates above \( \dot{m}_{\text{soft}} \), where \( \dot{m}_{\text{soft}} \) is given by

\[
\dot{m}_{\text{soft}} = 0.06 \left( l / 0.1 \right)^{1/2} (1 - f_s)^{-1}
\]

As in equation (9), \( f_s \) is the fraction of power reprocessed via magnetic flares and depends on the accretion rate itself, but we again expect that \( 1 - f_s \approx 1 \). Further, we scaled the poorly known compactness parameter on \( 0.1 \), since X-ray reflection calculations (Nayakshin 1998a,b) point to \( l \) of this order in both AGN and GBHCs.

Thus, the theory predicts that the X-ray spectrum should become steeper when \( \dot{m} \) increases above \( \dot{m}_{\text{soft}} \sim 0.06 \). For even higher accretion rates \( (\dot{m} \gtrsim 0.2 \text{ or so}) \), the disk becomes geometrically thick, and advection of energy (in the sense of Abramowicz et al. 1988) in the black hole will change the properties of the disk further, so that we do not attempt to describe those disks due to theoretical uncertainties. We can then complete classifying accretion disks by naming the radiation-dominated state with \( \dot{m}_r \lesssim \dot{m} \lesssim \dot{m}_{\text{soft}} \) the “intermediate” one, the state with \( \dot{m}_{\text{soft}} \lesssim \dot{m} \lesssim 0.2 \) “soft”, and, finally, the one with \( \dot{m} \gtrsim 0.2 \) “very high” in analogy with the very high state of GBHCs.

5. Comparison With Observations

Typical Hard Seyferts. As discussed in §4, our model predicts that Seyferts with the typical hard X-ray spectra should accrete at accretion rates below \( \dot{m}_{\text{soft}} \).

In order to estimate \( \dot{m} \equiv L / L_{\text{Edd}} \) based on observations, we need to know AGN
masses. Although these are unknown, variability studies may provide some help. For example, the global compactness parameter $l_g$ has been estimated for a sample of Seyfert Galaxies by Done & Fabian (1989). In their estimate, they assumed that the typical size of the emitting region is given by the distance traveled by light during the shortest doubling time scale $\Delta T$ observed for a given source ($l_g = \sigma_T L_x/(m_e c^4 \Delta T)$). For an accretion disk this typical size should be of order $\sim 10 R_g$. This yields $\dot{m} \simeq (10 m_e/2 \pi m_p) l_g/2 \simeq l_g/2000$. (see also Fabian 1994). Now, in Table 1 of Done & Fabian (1989), the maximum compactness is about 200, thus maximum $\dot{m} \sim 0.1$. Moreover, 80% of the sample have $\dot{m} < 0.02$, with smallest values of the order of $10^{-4}$. These estimates do not include the BBB, which could make a significant contribution to the bolometric luminosity of Seyfert Galaxies. However, using 1375 Angstrom fluxes reported by Walter & Fink (1993) for the sources with the highest values for the compactness, we estimated $L_{uv} \sim L_x$, so that inclusion of the emission at lower wavelengths did not affect our conclusions significantly.

Sun & Malkan (1989) fitted multi-wavelength continua of quasars and AGNs with improved versions of standard accretion disk models. They found that low-redshift Seyfert Galaxies radiate at only few percent of their Eddington luminosities. Rush et al. (1996) studied soft X-ray (0.1-2.4 keV) properties of Seyfert Galaxies. Their results indicate that $\sim 90\%$ of sources in their sample have soft X-ray luminosity below $10^{44}$ erg/s (with the mean value of order $\sim 10^{43}$ erg/s). If we assume the typical Seyfert 1 spectrum above 2.4 keV, i.e. a power law with intrinsic photon index $\sim 2$ and the cutoff at several hundred keV (e.g., Zdziarski et al. 1996), then total X-ray/gamma-ray luminosity of these objects can be a factor of 2-3 higher than the soft X-ray luminosity. Nevertheless, if a typical Seyfert Galaxy has the black hole mass of $\sim 10^8$, then the average bolometric luminosity of the Rush et al. (1996) sample is at or below $\sim 1\%$ of the Eddington luminosity. For NGC5548, using results of Kuraszkiewicz, Loska & Czerny (1997), we obtained $\dot{m} \sim (4 - 16) \times 10^{-3}$. Summarizing, there is some evidence that X-ray hard Seyfert 1 Galaxies accrete at a relatively low accretion rate, i.e., from probably just below $\dot{m} = 0.1$ to very low accretion rates of $\sim 10^{-3}$.

**Steep X-ray spectrum AGN, Narrow Line Seyfert 1 (NLS1) Galaxies.** Relatively recently, it was found that a subset of Seyfert Galaxies have unusually steep soft X-ray spectra (for a review, see Pounds & Brandt 1996, PB96 hereafter, and Brandt & Boller 1998). Common properties of the group include steep spectra, rapid variability, strong Fe II emission and identification with NLS1. PB96 speculated that the most likely explanation for the steep X-ray spectrum is an unusually high accretion rate. PB96 also showed that soft X-ray (i.e., 0.1-2 keV) spectral index is strongly correlated with the width of the H$\beta$ line for a sample of Seyfert Galaxies. Wandell & Boller (1998) and Wandel 1998 suggested an explanation of the correlation based on the simple idea that steeper X-ray spectrum implies a larger ionizing UV luminosity, which translates into a larger broad line region size, and thus a smaller velocity dispersion (since Keplerian velocity is $\propto 1/R^{1/2}$). They found that masses of the narrow-line Seyfert Galaxies tend to be lower that those of typical broad H$\beta$ line Seyferts, and thus have larger $\dot{m}$. Brandt, Mathur & Elvis (1997) found that the higher energy ASCA
slopes (2-10 keV) correlate with the Hβ line as well. Thus, NLS1 galaxies often have intrinsic X-ray slope that is steeper than that of normal Seyferts.

Laor et al. (1997) found the same correlation for a sample of quasars, and suggested that NLS1 galaxies accrete at a higher fraction of the Eddington accretion rate than normal Seyferts do. They assumed that the bulk motion of the broad line region is virialized, and that the scaling of the BLR with luminosity is that found from reverberation line mapping of AGN (e.g., Peterson 1993). In this case larger luminosities \( L \) correspond to larger BLR size, and thus smaller Hβ FWHM. Now, if \( \Gamma \) is larger for higher \( \dot{m} \), then the observed relation (smaller Hβ FWHM – larger \( \Gamma \)) ensues. However, no reason for \( \Gamma \) to become larger with increasing \( \dot{m} \) was given, except for a heuristic suggestion of Pounds et al. (1995) that if the power released in the corona remained constant, then the X-ray index would become steeper with increasing bolometric luminosity (which is equivalent to assuming \( f/(1 - f) \propto (\dot{m})^{-1} \)). Our theory of accretion disk states may provide a natural explanation for this spectral steepening (see §4).

**Observational Constraints on Corona/Disk Energy Partitioning.** So far in §5, we discussed how the shape of the X-ray spectrum depends on \( \dot{m} \). Since we believe that the BBB is produced by reflection of X-rays produced by the flares off the surface of the disk, we should also see if we can test our theory based on observations of BBB. Here we will only concentrate on the overall normalization of the bump, taking up the question of its spectrum in the other contributed paper in these proceedings (Nayakshin 1998a).

Walter & Fink 1993, Walter et al. 1994 and Zhou et al. 1997 studied the BBB of Seyfert 1’s and found that this spectral feature is ubiquitous in these sources. The observed spectral shape of the bump component in Seyfert 1’s hardly varies, even though its luminosity \( L_{\text{bbb}} \) ranges over 6 orders of magnitude from source to source. However, recent work of Zheng et al. (1997) and Laor et al. (1997) showed that quasars in their (different) samples do not show the BBB component.

There are two reasons why we think that there is no inconsistency here. As discussed above, typical Seyfert Galaxies should accrete at a relatively small accretion rates, i.e., \( \dot{m} \lesssim \dot{m}_{\text{soft}} \), which corresponds to luminosity \( L \simeq 5 \times 10^{44} M_8 \) ergs/sec. In the sample of Walter & Fink (1993), very few objects have UV luminosity above few \( \times 10^{44} \) ergs/sec, whereas Zheng et al. (1997) fit to the mean spectrum in their sample gives \( L \simeq 8.5 \times 10^{45} \) ergs/sec. Similarly, almost all AGN in Laor et al. (1997) sample have \( L_{3000} > 10^{45} \) ergs/sec. Therefore, Walter & Fink (1993) sample contains Seyferts that are dimmer than sources in the other two samples by factors 10 to 100. Accordingly, sources in Walter & Fink (1993) sample may accrete at hard or intermediate state \( \dot{m} \lesssim \dot{m}_{\text{soft}} \), whereas AGN of the two other samples could accrete at the soft and/or very high state regime. Our first argument is that according to discussion in §4, the more luminous sources must be X-ray weak, i.e., \( f \ll 1 \) and thus the normalization of the BBB is \( \sim f/2 \ll 1 \), so that it disappears on the background of the disk thermal emission (1 – \( f \simeq 1 \)) in samples of Zheng et al. (1997) and Laor et al. (1997). Our second argument has to do with the shape of the reprocessed spectrum. Namely, we find (Nayakshin 1998a,b) that an ionization instability may lead to the reprocessed spectrum being a power law with basically same index as that
of the incident X-rays, up to $\sim 30\,\text{keV}$, which will blend to undetectability with the X-ray spectrum from the flares (see also the limit of the “hot medium” in Zycki et al. 1994).

Another well known observational fact for quasars is the correlation between the optical to X-ray spectral slope $\alpha_{\text{ox}}$ and the optical luminosity (Green et al. 1995). The optical to X-ray index $\alpha_{\text{ox}}$ is not a real spectral index in this energy range, but is defined as the index of an imaginary power law connecting observed optical and X-ray emission. Wilkes et al. (1994), Green et al. (1995) show that more luminous sources have larger $\alpha_{\text{ox}}$, i.e., more luminous objects have comparatively less X-ray emission. This is again in a qualitative agreement with our theory.

**GBHCs state transitions** As Nayakshin (1998a) shows, PCD model with magnetic flare parameters (most notably $l$) same as for AGN case, when re-scaled for the case of GBHCs gives harder X-ray spectra and weaker reprocessing features (as compared to AGN), explaining peculiarities of hard state spectra of GBHCs such as Cyg X-1. We thus will try to use the same logic as for AGN, only changing $M$ from $\sim 10^8$ to $\sim 10$, to understand GBHCs spectral states, that are often referred to (in order of increasing $\dot{m}$) as hard, intermediate, soft and ultra-soft state (e.g., Grove, Kroeger & Strickman 1997 and Grove et. al. 1998). Note that for $M \sim 10\,M_\odot$, the transition from the gas- to radiation-dominated regimes happens at a considerably higher $\dot{m}$, namely, $\dot{m}_r \simeq 1.7 \times 10^{-2} (1-f)^{-9/8}$. Further, a transition is in fact defined by fraction $f$ changing from large to small numbers, so we can assume $f \simeq 1/2$, such that $\dot{m}_r \simeq 0.037$. This number is remarkably close to where the state transition are observed to occur for GBHCs.

Barret, McClintock & Grindlay (1996) assembled a sample of GBHCs in a $L_{\text{hx}} - L_\text{x}$ phase space, where $L_{\text{hx}}$ is the hard X-ray luminosity in the range $20 - 200\,\text{keV}$, and $L_\text{x}$ is the X-ray luminosity in the range $1 - 20\,\text{keV}$. One of the striking results of this exercise is that GBHCs always have most of their power in
the soft X-ray component (presumably the optically thick disk emission) when they radiate at a high fraction of their Eddington luminosity. Barret et al. (1996) also plotted the evolutionary track of the GBHC transient source GRS 1124-68 on the same $L_{\text{hx}} - L_x$ phase diagram. Since we are interested in the dimensionless accretion rate, we reproduce data of Barret et al. in terms of $L_x/L_{\text{Edd}}$ and $L_{\text{hx}}/L_{\text{Edd}}$ in Figure (1), adding some data for Cyg X-1. Figure (1) shows that the spectra of GBHCs indeed have most of their power in the hard component up to $\dot{m} \sim 0.04$, and then there is a rather strong spectral transition, which confirms our finding that $f$ decreases with increasing $\dot{m}$ for $\dot{m} \geq \dot{m}_r$. This behavior unites AGN and GBHCs.

There are differences, too. Just as in AGN case, above the gas- to radiation transition, our theory predicts existence of the intermediate state ($\dot{m}_r \leq \dot{m} \leq \dot{m}_{\text{soft}}$). However, note that for GBHCs the intermediate state is squeezed in in the narrow interval between the hard and the soft states, whereas for AGN it is not the case because $\dot{m}_r$ is much lower than $\dot{m}_{\text{soft}}$ (if compactness parameter $l$ is indeed of order $\sim 0.1$ in both AGN and GBHCs, as suggested by spectral modelling). The difference is then such that the intermediate state essentially disappears in GBHCs, so that as $f$ decreases at the region $\dot{m} \sim \dot{m}_r$, the X-ray spectral index steepens, while for AGN case it still may stay hard until $\dot{m}$ reaches $\sim \dot{m}_{\text{soft}}$.

6. Discussion

We have shown that one cannot power coronae of accretion disks with diffuse magnetic fields, at least not in a way consistent with observations of hard states in Seyferts and GBHCs. In short, the problem is that these fields are transported into the corona “too slowly” ($v_b$ is realistically just a fraction of sound speed), so that during their rise, the fields contribute “too much” to the angular momentum transport in the disk and to the local disk heating, which in turn produces disk thermal radiation. It is then not possible to ever produce more power in X-rays than in the disk thermal emission (i.e., $f \ll 1$), which is inconsistent with observations.

However, if one assumes that most of the disk magnetic field is localized in the form of strong magnetic flux tubes, the field contribution to viscosity goes down substantially. This then leads to a decrease of the local disk heating, so that it becomes possible for accretion disks to release most of their energy in coronae rather than through the common thermal radiation diffusion.

We further have shown that in radiation-dominated accretion disks, diffusion of radiation into flux tubes does not allow the tube magnetic fields to reach equipartition values, and that instead the field pressure is limited to the gas pressure. Thus, we concluded that magnetic fields in the radiation-dominated disks are in the diffuse form, and thus such disks cannot produce X-rays as efficiently as the gas-dominated disks can. Using this idea, and some spectral constraints from PCD model, we also showed that the observed spectral transitions and states of GBHCs and the mounting evidence for presence of similar states in AGN accretion disks support the theory of accretion disks with magnetic flares. We believe that this, together with other successes of the theory,
provides one with optimism that magnetic flares is the physics that was missing in the standard accretion disk theory for decades.

Acknowledgments. The author is very thankful to the workshop organizers for the travel support, and to F. Melia for support and useful discussions in an early stage of this work.

References

Abramowicz, M.A., et al. 1988, ApJ, 332, 646
Barret, D., McClintock, J.E., & Grindlay, J.E. 1996, ApJ, 473, 963
Brandt, W.N., Mathur, S. & Elvis, M. 1997, MNRAS, 285, L25
Brandt, W.N., & Boller, Th. 1998, preprint, astro-ph/9808037
Done, C. & Fabian, A. C. 1989, MNRAS, 240, 81
Fabian, A.C. 1994, ApJSupplement, 92, 555
Frank, J., King, A., & Raine, D. 1992, Accretion Power in Astrophysics (Cambridge, UK: Cambridge University press).
Galeev, A. A., Rosner, R., & Vaiana, G. S., 1979, ApJ, 229, 318
Green, P. J., et al. 1995, ApJ, 450, 51
Grove, J.E., Kroeger, R.A. & Strickman, M.S. 1997, Proc. 2nd Integral Workshop, ESA SP-382, p. 197
Grove, J.E., et al., 1998, submitted to ApJ.
Haardt F., & Maraschi, L., 1991, ApJ, 380, L51
Haardt F. & Maraschi L., 1993, ApJ, 413, 507
Haardt F., Maraschi, L., & Ghisellini, G. 1994, ApJ, 432, L95
Kuraszkewicz, J., Loska, Z., Czerny, B. 1997, Acta Astronomica, v.47, 263
Laor, A., et al. 1997, ApJ, 477, 93
Lightman, A.P., & Eardley, D.M. 1974, ApJ, 187, L1
Nayakshin, S. & Melia, F. 1997, ApJ, 484, L103
Nayakshin, S. 1998a, these proceedings
Nayakshin, S. 1998b, PhD thesis, University of Arizona.
Parker, E.N. 1955, ApJ, 121, 491
Parker, E.N. 1979, Cosmical Magnetic Fields, Clarendon Press, Oxford.
Peterson, B.M. 1993, PASP, 105, 247
Pounds, K.A., Done, C., & Osborne, J.P. 1995, MNRAS, 277, L5
Pounds, K.A., & Brandt, W.N. 1996, astro-ph/9606118
Poutanen, J. & Svensson, R. 1996, ApJ, 470, 249
Poutanen, J., Svensson, R., & Stern, B. 1997, Proceedings of the 2nd INTEGRAL Workshop The Transparent Universe, ESA SP-382
Priest, E.R., 1982, Solar magnetohydrodynamics, Dordrecht, D. Reidel Publishing Co.
Rush, B., Malkan, M.A., Fink, H.H., & Voges, W. 1996, ApJ, 471, 190
Sakimoto, P.J., & Coroniti, F.V. 1989, ApJ, 342, 49
Shakura, N.I., & Sunyaev, R. A. 1973, ApJ, 24, 337
Shapiro, L.E., Lightman, A.P., & Eardley, D.M. 1976, ApJ, 204, 187
Stella, L., & Rosner, R. 1984, ApJ, 277, 312
Stern, B., et al. 1995, ApJ, 449, L13
Svensson, R. & Zdziarski, A. A. 1994, ApJ, 436, 599
Svensson, R. 1996, ApJSupplement, 120, 475
Sun, W., & Malkan, M.A. 1989, ApJ, 346, 68
Tajima, T., & Shibata, K. 1997, Plasma Astrophysics, Addison Wesley
Tsuneta, S. 1996, ApJ, 456, 840
Vishniac, E.T. 1995a, ApJ, 446, 724
Vishniac, E.T. 1995b, ApJ, 451, 816
Walter, R., & Fink, H.H. 1993, A&A, 274, 105
Walter, R., et al. 1994, A&A, 285, 119
Wandel, A., & Boller, Th. 1998, A&A, 331, 884
Wandel, A. 1998, these proceedings
Wang, T., Lu, Y., & Zhou, Y. 1998, preprint, astro-ph/9801142
Wilkes, B., Mathur, S., McDowell, J., & Elvis, M. 1994, ApJ, 185, 3402
Zdziarski, A.A., et al. 1996, A&A Suppl., 120, 553
Zheng, W. et al. 1997, ApJ, 475, 469
Zhou, Y., et al. 1997, ApJ, 475, L9
Życki, P.T. et. al. 1994, ApJ, 437, 597