Delving deeper into the physics of jets in black hole binaries using an outflow-dominated model.

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Abstract

Jets are all around us. These collimated plasma outflows are found in e.g. X-Ray Binaries (XRBs), Active Galactic Nuclei (AGN), Cataclysmic Variables, protostars. For this project we used an outflow-dominated model to explore several issues related to jets in XRBs. Firstly, the peculiar black hole binary (BHB) GRS 1915+105 shows behavior that is quite distinct from what we observe in the other (∼20) known BHBs. While generally BHBs quench their jets when their luminosity becomes too high, GRS 1915+105 interestingly reveals a jet-bearing state, in spite of the fact that it accretes at near- or super-Eddington rates. How does this plateau state, as it is commonly referred to, relate to the canonical XRB “jet”-state, the so-called hard state? Secondly, from spectral energy distributions and timing analyses, the BHB GX 339–4 is observed in the hard state at different X-ray luminosity levels, depending on which state it was in previous to entering the hard state. Can we explain this behavior if we consider jets as the driving force? Thirdly, the outflow we use to interpret the above issues is not time-dependent. Hence electron cooling due to synchrotron losses were not yet implemented. Comparing the time electrons need to emit most of their energy when subjected to a certain magnetic field strength to the timescale they reside in that field, we can estimate the change in the particle distributions and, consequently, possible changes in the spectrum that results from the distribution. Are synchrotron losses significant in XRB systems? We will discuss these three questions and possible implications for jet-physics in BHBs as well as the relation to their supermassive counterparts, the AGN.

To answer the first question we analyze two broadband plateau state GRS 1915+105 data sets and compare the best-fit parameters to those found in BHBs representative of more “canonical” behavior. Although we are able to obtain statistically satisfactory results for both sets, they are revealed to be quite distinct and only one of them is deemed truly convincing on physical grounds. As expected for a black hole with such a high accretion rate, the credible dataset inhabits an extreme part of jet parameter space, portraying an extreme power input into and the magnetic domination of the jet. Also the region where particle acceleration occurs is much further removed from the jet base as usual. The issues with the “less-convincing” dataset confirm the need for a variable column density for the plateau state.

To answer the second question we analyze two GX 339–4 hard state datasets. One representative for the rise phase of the 2004-2005 outburst and one representative for the decline. We fit the rise phase data and use the results to fit the decay phase, while changing the least number of parameters possible in order to find simple physical mechanisms that are able to drive the hysteresis. In doing so we reveal three categories of possible solutions, representing some interesting trends that can be associated with physical mechanism in the literature. The main result is that the HID hysteresis appears to be related to the particle acceleration, as all three categories require an increase of the distance to the particle acceleration front.

We investigate the particle acceleration further in the third part of this thesis. After the acceleration front in our jet, a fraction of the radiating particles is continuously accelerated from a quasi-thermal distribution into a power-law distribution. At this point the synchrotron losses can become evident as a steepening of that power-law slope by unity, above a certain energy. We compare the timescales a particle has and needs to cool to estimate this break energy and find that in GX 339–4 synchrotron cooling is certainly relevant and influences the spectrum. In the new model the fraction of accelerated leptons has become a free parameter. Hence the synchrotron component emitted by each segment of the jet can now change shape between pure quasi-thermal and absorbed power-law, depending on other parameters, making it much harder to predict the total post-acceleration synchrotron spectrum. The Chapter again reveals the acceleration front distance is the most important parameter, as this parameter determines both the magnetic field as well as the time the particle has to cool and hence largely determines the shape of the synchrotron spectrum.
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In 1918 the American astronomer Heber Doust Curtis observed an optical phenomenon in the giant elliptical galaxy M87 (see Figure 1.1), using the 0.9-m Crossley reflector on the Lick Observatory in California. He described it as “a curious straight ray...apparently connected to the nucleus”. The importance of this “beam” only started to become clear decades later, when Baade & Minkowski (1954) identified this source with Virgo A (or NGC 4486): the brightest radio source in that constellation. In 1956, a weaker, extended radio halo was found with the Cambridge radio telescope (Baldwin & Smith 1956). Using an Aerobee sounding rocket M87 was also identified as a strong X-ray source (Byram et al. 1966). Even later, more than 50 years after Curtis’ first observation, scientists (e.g. Blandford & Rees 1974 and Scheuer 1974) started associating the emissions with a flow, or “jet” of relativistic particles, extending over thousands of light years and comprising ejected gaseous material from the galaxy’s core, where a supermassive (∼6.4×10⁹ M☉; Gebhardt & Thomas 2009) black hole is thought to lurk.

More recently, jets have also been observed in our own Galaxy in (much smaller) binary systems, consisting of a stellar mass compact object and a star. Fabian & Rees (1979) were able to explain Doppler shifts in optical atomic emission lines of SS 433 (Margon et al. 1979) using a kinematic model of two precessing, collimated jets with a bulk velocity of 0.26c. This makes SS 433 the first ever discovered microquasar (see Section 1.3). It is a remarkable eclipsing system: We see it almost edge-on at an inclination of 79° to the line of sight. The jets are “misaligned” and precess around the inclination axis at an angle of 20°, with a period of 162 days. This angle causes the jet to trace out a peculiar “corkscrew” pattern (see Figure 1.2). Cherepashchuk et al. (2005) modeled lightcurves obtained by INTEGRAL, obtaining masses for the optical and compact component of 30 M☉ and 9 M☉ respectively, implying an early type (∼B; Charles et al. 2007.) companion star.

Figure 1.1: The jet in M87, first described by Curtis in 1918. Credit: The Hubble Heritage Team STScI/AURA) and NASA/ESA.

Figure 1.2: VLA Image of Microquasar SS 433 (credit: Blundell & Bowler, NRAO/AUI/NSF)
and a black hole primary. These components make SS 433 a High Mass X-ray Binary (HMXB; see below).

By now we know that jets are not limited to radio galaxies or binaries, but that they are in fact an ubiquitous phenomena. They are found to range over many orders of magnitude in length (from AU to pc) and have been discovered in numerous sources, in e.g. X-Ray Binaries (XRBs), Active Galactic Nuclei (AGN), white dwarfs (WDs; appearing as supersoft X-ray sources), during the end-phases in evolution of massive stars, such as planetary nebulae, supernovae (SNe) and Gamma Ray Bursts (GRBs), and as “stellar jets”, in protostars (T Tauri) and pre-main sequence (FU Orionis) stars. It is as of jet unclear whether jets exist in Ultra-Luminous X-ray sources (ULXs), but clearly plasma outflows are of considerable importance.

1.1 Why do we study jets?

There are numerous reasons to study jets, most importantly:

- Jets are ubiquitous and studying them will therefore add to our general knowledge of the universe.

- Jets influence their launching systems, extracting large amounts energy and matter and will hence change the evolution those launching systems. Transporting appreciable quantities of matter the outflows will greatly influence the environment of the system (a phenomenon known as feedback) and because of this, the evolution of the universe on a whole.

- Although jets can be found in so many systems even simple jet launching, collimation and acceleration mechanisms are largely unknown. For the moment perhaps our best hopes of understanding jet formation lies in fully three dimensional (see e.g. McKinney & Blandford 2009) radiative (Johnson & Gammie 2003) general relativistic MHD simulations (Königl 2007) of the accretion flow. However considering the amount of computer time needed in such an approach there may still be a lot to learn using somewhat simpler, less CPU-intensive models (like the one used for this work; see Chapter 2).

- Often the structure and composition of jets are unknown, in terms of what the plasma constituents are (e.g. electrons and protons or electrons and positrons) and division of the total energy budget over the plasma components and energy densities.

- Jets provide us with means to test General Relativity (GR). E.g. Livio (1999) found evidence that the initial velocity of the jet is commensurate to the escape velocity \( v_{\text{esc}} = \sqrt{2GM/r} \), where \( G \) is the gravitational constant, \( M \) is the mass of the object responsible for the gravitational potential and \( r \) is the distance from the center of that potential) from the object that is launching the jet (in all classes of objects). So if we are able to measure jet bulk velocities accurately enough we could learn at what distance from a compact object a jet gets launched and potentially learn something about physics in the strong field regime.

1.2 What we DO know about jets

From the above there there are clearly many unknowns when it comes to jets. What do we know and how do we observe them?

Most jet formation scenarios involve systems with an accretion disk (see Section 1.3 and the caption of Figure 1.4) that is threaded by a large-scale magnetic field (Livio et al. 2003) and we may thus suspect three components are required: accretion, rotation\(^1\) and magnetic fields. Jets

\(^{1}\)Another major contender (next to the accretion disk) for supplying the rotation “required” to produce jets is the spin of the central black hole (Remillard & McClintock 2006). Black hole spin is also thought to be able to provide the power required to form a jet (Blandford & Znajek 1977).
typically occur in pairs, appearing to emerge from opposite sides of the accreting system (see Figure 1.3 and the title page for artist’s impressions).

In many instances the jets can not be seen directly and their existence is presumed from e.g. radio lobes on either side of the central object or because the radio emission is strongly polarized. Polarized emission is indicative of synchrotron emission; photons radiated by relativistic particles accelerated in a magnetic field. Synchrotron radiation is generally the strongest in radio and consequently jets have been predominantly observed at radio wavelengths. The “smoking gun” for the existence of a jet is a flat or inverted \( f_\nu \propto \nu^\alpha \), where \( \alpha \sim 0 \) radio spectrum (see Figure 2.1). The synchrotron nature of these radio spectra is further implied from the fact that the emission is distinctly non-thermal and because of relatively high brightness temperatures. The brightness temperature can be expressed in terms of \( T^\text{obs}_b \), the observed intensity at a certain frequency \( \nu \), as \( T_b = \frac{\lambda^2}{2k} I^\text{obs}_\nu \), where \( k \) is Boltzman’s constant and so a typical length scale \( \lambda \) for the radio-emitting regions can be inferred. These scales normally greatly exceed the expected scales of the other components in the accreting systems, like the orbital radius (Gallo 2009), suggesting a larger flow must be responsible for the emission.

Evidence is mounting that jet emission can extend over the entire broadband spectrum. For bright transients, e.g. Russell et al. (2006) estimate up to 90% of the near-infrared is dominated by jet production. Jets have also been resolved at energies up to X-rays (in e.g. R Aqr; Kellogg et al. 2001 and Cen A; Kraft et al. 2002) and have even been observed at TeV \( \gamma \)-ray energies (e.g. in the BL Lac object Mkn 501, Pian et al. 1998).

The higher energy photons can come from synchrotron processes, but are usually attributed to a radiative process called the inverse Compton (IC) effect, in which soft seed photons (from the accretion disk or stellar companion, see Section 2) collide with high energy particles and are consequently accelerated to increasingly high energies. In a jet this process usually occurs at the base, in the form of synchrotron self-Comptonization (SSC), where the particles that create seed photons through the synchrotron process are themselves responsible for the subsequent acceleration via the collisions. The interpretation that the observed photons originate in synchrotron and/or SSC processes is suggested by correlations between the radio, infrared and X-ray bands (see Section 1.4 and 4.1.2).

1.3 Jets in black hole systems

As mentioned jets can be found in hugely varying classes of astrophysical objects. Hence for this work we will limit ourselves to studying jets in Black Hole Binaries (BHBs), for reasons we will now explain.

Einstein’s theory of General relativity predicts that the physics related to black holes should be the same, regardless of mass. So potentially we can learn the most from jets in these systems as (through the process of accretion) black holes are observed over many orders of magnitude in mass: From stellar mass black holes in our own Galaxy (in BHBs/XRBs, see below), to supermassive black holes (in Active Galactic Nuclei; observed over a mass range from millions to billions of solar masses) at the cores of distant galaxies. As mentioned above, in the stellar mass systems, the compact object is located in a binary system and accretes from a companion star. Their extremely massive relatives accrete interstellar matter or an occasional unfortunate star. In both cases, the matter that has fallen prey to the enormous gravitational field of a compact object marks its demise with an extremely energetic and well-measurable death cry: it spirals inwards, forming an accretion disc as it does so (also see Figure 1.4), which heats up through various dissipative processes in the gas, until the resulting plasma is

![Figure 1.3: Artists impression of a black hole binary that has launched jets.](image-url)
Figure 1.4: Schematic representation of a typical LMXB geometry and various important components. The accreting compact object (either a black hole or a neutron star) is on the left and the mass donor is on the right hand side. The formation of an accretion disk can be qualitatively explained as follows: When the companion star first exceeds its Roche-lobe, a stream of gas will fall within the gravitational pull of the attractor. This gas still has appreciable angular momentum, preventing the matter from falling directly towards the central object. Unless the plasma can get rid of this angular momentum, it will remain in a (nearly) circular orbit, forming a ring. However in such a ring various dissipative processes (e.g. collisions of gas elements, shocks, viscous dissipation) allow the matter to transfer angular momentum and get rid of its kinetic energy. Assuming there are no external torques, angular momentum can only be transferred outwards by the internal torques. So as the inner material loses angular momentum and is lowered into the gravitational well, the outer rings gain in angular momentum and spiral outwards, and an accretion disk is formed. This disk remains instrumental in removing angular momentum from the gas, in the end allowing it to spiral into the event horizon (Frank et al. 2002). The central region of the disk is thought to house a Comptonizing corona. Where the transferred gas hits the accretion disk a hot spot forms. If the infall-rate of the plasma is too high, the accretor may not able to process all the matter and a collimated bipolar outflow or jet can be created, depicted on both sides of the disk. The outflow need not be collimated, but can also occur in the form of a disk wind. There are indications that these forms of outflows complement each other: If the flow-rate of one component decreases the flow-rate of the other increases, ultimately balancing the total outflow (also see Section 3.5). The X-rays coming from the central region are thought to irradiate the the companion, presumably raising its effective temperature. Irradiation of the companion star (and outer disk) is expected to make a significant contribution in bright systems, like GRS 1915+105 (Chapter 3).

so hot it emits light, primarily in the ultraviolet in AGN, and in X-rays in XRBs, where the discs are even hotter (see below).

The amount of (potential) energy that can be liberated in these systems by a particle brought in from infinity is enormous and can be estimated from $\Delta E = GMm/R$, where $G$ is the gravitational constant, $M$ and $R$ are the mass and radius of the central object, respectively and $m$ is the mass of the particle. The ratio $M/R$ defines the compactness of an object and determines the accretion efficiency, which obviously scales linearly with the black hole mass. A (non-rotating) “Schwarzschild” black hole has a radius (or event horizon) of $2GM/c^2 = 2r_g$, where $r_g$ is the so-called gravitational radius.

The innermost stable orbit (ISCO; after which the matter will fall quickly through the event
Figure 1.5: Edge-on view of the “Fundamental Plane of black hole accretion” (Merloni et al. 2003), correlating the radio loudness \( L_{\text{radio}} \), the X-ray luminosities \( L_X \) and the black hole mass \( M \). The plane can be represented as \( L_{\text{radio}} \propto L_X M^{0.8} \). The solid line is the best fit. The abbreviations in the top-left corner stand for Galactic Black Hole (GBH) and different types of AGN: LINERS and transition objects (L+T), Seyferts (Sy) and Quasi-Stellar Objects or quasars (QSOs).

horizon) in a Schwarzschild black hole lies at 6 \( r_g \), but can lie at 1 \( r_g \) for a maximally spinning, or “Kerr” black hole. Obviously this factor 6 difference in the compactness is negligible compared to the observed factor 10^8 range in black hole mass, suggesting accretion physics should roughly scale in black holes.

This idea of mass-scaling appears to be confirmed by the discovery of a Fundamental Plane (FP; Merloni et al. 2003; Falcke et al. 2004) of black hole activity (see Figure 1.5). This relatively recent development is the basis for much of the current theoretical research into accreting black hole systems. Under certain circumstances (as long as the accreting system is in a state associated with jets; see Section 1.4) and allowing for some scatter, the black hole mass, radio loudness and X-ray luminosity all appear to be correlated, i.e. the correlations yield a plane in the three-dimensional space spanned by these parameters. An X-ray/radio emission correlation is to be expected if there is a fundamental connection between the inflow and outflow (see Section 4.1.2).

How can we interpret and compare the physics governing AGN and BHBs, while they appear to be in an entirely different physical regime? Although exact physical parameters, like the size of the system or the temperature of the disk, may vary with mass (see Figure 1.6), we usually see a relatively simple mass-dependency of most parameters. For example, if there are no dramatic mass-dependent changes of the accretion geometry, the surface area through which matter with significant angular momentum accretes scales as \( M^2 \). The luminosity generated in Eddington rates^2 by a surface area scaling with \( M^2 \) is \( L \propto M \). As blackbody radiation scales with the emitting area \( A \) and temperature \( T \) as \( L \propto AT^4 \), the disk temperature scales with mass as

^2The Eddington rate gives a theoretical upper limit to the accretion rate and can easily be derived equating the outward radiation force, supplied by the photons created in the accretion process, to the inward gravitational force. Assuming hydrogen is accreted in a spherical geometry, \( L_{\text{Edd}} \sim 1.4 \times 10^{38} (M/M_\odot) \text{ erg s}^{-1} \). The Eddington mass-accretion rate can be estimated comparing \( L_{\text{Edd}} \) with the X-ray luminosity \( L_X \).
Introduction

**Table:**

<table>
<thead>
<tr>
<th>Property</th>
<th>XRBs</th>
<th>AGN</th>
</tr>
</thead>
<tbody>
<tr>
<td>Mass ((M_{\odot}))</td>
<td>few</td>
<td>(\sim 10^{6-9})</td>
</tr>
<tr>
<td>Jet-length ((\text{AU-Pc}))</td>
<td>(\sim 10^{6})</td>
<td>(\text{lyrs})</td>
</tr>
<tr>
<td>Disk radius ((\text{km}))</td>
<td>10³</td>
<td>10⁹</td>
</tr>
<tr>
<td>Disk temperature ((\text{K}))</td>
<td>(&gt; 10^{6})</td>
<td>10³</td>
</tr>
<tr>
<td>Peak radiation of disk</td>
<td>X-ray</td>
<td>UV/Optical</td>
</tr>
<tr>
<td>Typical dynamical timescale</td>
<td>ms-months</td>
<td>(10^{3-6}) yrs</td>
</tr>
</tbody>
</table>

**Figure 1.6:** Comparison between accreting supermassive black hole systems and their stellar-mass counterparts. The physics governing both types of systems are thought to scale with mass. The Table lists some fiducial values for the physical properties. Most importantly it is feasible to observe dramatic changes in microquasars within a timescale relevant for humans. In contrast, comparable evolutionary events in quasars simply have not occurred yet, ever since we first started observing them in the late 1950s. (Mirabel & Rodríguez 1998)

\(T \propto M^{-1/4}\), and hence the energy at which we observe the disk in AGN is a factor \(\sim 100\) lower in AGN, compared to XRBs.

To emphasize the similarities in black hole systems that show jets Mirabel & Rodríguez (1998) coined the term microquasar to describe the BHB GRS 1915+105 (see Chapter 3). An important difference, however, between the “light-weight” and “heavy-weight” systems, is that the time required to observe changes in accreting black hole systems increases with mass. This reveals what is probably the greatest advantage of studying XRBs: dynamical changes in XRBs take place on timescales of milliseconds to years, while analogous changes could easily take thousands of years in the most massive quasars.

So although many more extragalactic AGN \((> 2 \times 10^4; \text{Kauffmann et al. 2003})\) have been observed than XRBs \((\gtrsim 264; \text{Liu et al. 2001; Liu et al. 2006})\) or BHBs \((\sim 20 \text{ confirmed; Remillard & McClintock 2006})\) and although we could perhaps receive a larger photon flux (on the same dynamical timescale) from jets launched by supermassive black holes, we only see a “snapshot” of these systems. XRBs are however also numerous, with several types observed under several distinct accretion conditions, interchanging on the typical XRB timescales mentioned above. These so called X-ray spectral states (see Section 1.4 below) allow for a relatively systematic approach for studying them. So if the underlying physics indeed scales, we should be able to learn about supermassive black hole evolution by studying their low-mass counterparts, that (depending on their spectral state) due to their relative proximity can still be very bright. Assuming the “snapshots” of AGN also represent states, similar to the ones found in XRBs, we could use an approach comparable to the study of stellar evolution to try and understand the evolution of AGN over cosmic time, however within the span of a human life, or even over the course of a PhD. That is, providing we can obtain a large enough un-biased sample of AGN to check if the spectral states found in XRBs map onto the states that we find in AGN, if we account for mass-scaling.
Lastly, in BHBs we can study the variability in the inner-most region around a black hole using X-ray power density spectra (PDS) analysis (in a PDS the power received per frequency is plotted). While we can do the same in AGN, in XRBs the low-to-high frequency variability occurs on timescales such that measurements are still feasible in terms of time spent on a single source. For AGN only highest frequency variability analysis is possible, as dedicated all-sky monitoring devices have only been orbiting Earth for a limited time, limiting the usefulness of this method.

Neglecting possible intermediate classes, there are two types of XRBs: The High Mass X-ray Binary (HMXB) and the Low Mass X-ray Binary (LMXB; see Figure 1.4). They are named according to the mass of the secondary star. In HMXBs a more massive star is present, and mass transfer occurs via a stellar wind. In LMXBs the mass of the companion is comparable to that of the sun, and mass transfer exists because companion exceeds its Roche-lobe in size. Usually that happens via normal evolutionary process, in which such a star e.g. reaches its Giant phase. This work focuses on two LMXBs: GRS 1915+105 (see Chapter 3) and GX 339–4 (see Chapter 4 and 5). Both these LMXBs are black hole binaries (BHBs) so we will not consider neutron stars as the central compact object.

1.4 The canonical black hole spectral states and the HID

While a singular black hole is defined by only three fundamental physical properties: that of mass, spin and charge (Fender 2009), XRBs are observed in a number of “canonical” states, classified according to the X-ray and power density spectrum. I will refer to BHBs that strictly adhere to these states, as “canonical” BHBs.

Attempts at understanding the basic X-ray spectral states and their associated observations have been done using numerous models. The basis of most models is the picture of a geometrically thin but optically thick accretion disk (Shakura & Sunyaev 1973), with a temperature profile $T \approx R^{-3/4}$. Mitsuda et al. (1984) and Makishima et al. (1986) devised a non-relativistic approximation to the radiation of the “standard” Shakura & Sunyaev thin disk, resulting in a multi-color blackbody spectrum. For this multi-color disk (MCD) model a total luminosity $L_{\text{disk}} = GM \dot{M}/2R_{\text{in}}$, where $R_{\text{in}}$ is the radius of the inner edge of the disk, is assumed. In spite of their age the thin disk models accurately describe the thermal component in many (soft state and possibly intermediate state; see below) X-ray observations.

These days the general consensus is that the spectral states arise from geometrical changes in the central part of the accretion flow. This interpretation was instigated by the work from Esin et al. (1997), in which the MCD is combined with the radiatively inefficient ADAF (see Figure 1.7 and its caption) model and successfully unify four basic spectral states, attributing the changes to the truncation and refilling of the inner region of a thin disk. Later variations or expansions upon this model, are e.g. the CDAF (incorporating convection; e.g. Igumenshchev & Abramowicz 1999, Narayan et al. 2000), the Magnetically Dominated Accretion Flow (MDAF; putatively existing in a region internal to the ADAF; Meier 2005), and the ADIOS (involving disk winds; Blandford & Begelman 1999). All these models have their own merits, but all ignore radiation bands other than the X-ray regime. Therefore we will use an entirely different model, that focusses on the outflow rather than on the inflow (it does however include an MCD). It is still the only model that is able to successfully explain multi-wavelength broadband observations (see Section 2).

The canonical spectral states are usually defined according to the properties of the SEDs and PDSs of a system. However, depending on the general purpose of the work, there are many variations in the literature in nomenclature and in the exact definitions. E.g. Remillard & McClintock (2006) focus more on the energy spectral behavior in order to set up a scheme for physical modeling.

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3The highest frequency variability in AGN is commensurate to low-frequency variability in BHB.
4An object’s Roche-lobe is defined by the point in space between two objects where gravity balances out. The distance from this point to each object is determined by the mass ratio $M_{\text{BH}}/M_{\text{star}}$. The volume occupied by the equi-potential surfaces surrounding an object and going through this inner Lagrangian point (or L1) is called the object’s Roche lobe.
5As large-scale charged objects have not yet been observed in the universe, black holes are generally assumed to be electrically neutral. (Fender 2009)
Introduction

Figure 1.7: Schematic representation of the BHB geometries suspected to underly the canonical BH states, adapted from (Esin et al. 1997), incorporating some modifications according to a private communications with Dr. S. Markoff. In the original interpretation of the spectral states in view of these geometries, the accretion rate $\dot{m}$, scaled to the Eddington (mass) limit, increased towards the top. More recent empirical evidence however suggests that the states are not strictly dependent on $\dot{m}$, so the arrow on the right hand side should be ignored (private communication with Dr. T. Belloni). In this Figure, the geometrically “thin” accretion disk is represented by the red horizontal bars. Usually such a disk is referred to as a “multi-color blackbody”; in which each annulus to radiates a Planck curve in accordance with its temperature and surface area. The temperature decreases outwards as the surface area grows. In this view the different X-ray states are believed to stem from variations in the inner disk radius ($r_{in}$ in the MNW05 model; see Section 2). In the soft state the disk may even reach all the way in to the Innermost Stable Circular Orbit (ISCO). The dots represent an Advection Dominated Accretion Flow (ADAF; see e.g. Narayan & Yi 1995a, Narayan & Yi 1995b). Jets are now believed to be present not only in the Low/Hard state, but also in the Quiescent state. Note that the jets depicted here are not to scale: We expect the launching point to be near the transition radius of the disk.

while Klein-Wolt & van der Klis (2008) try to cover all the observable states in XRBs, focussing on the variability. Below we outline the canonical states according to both these works. Although when classifying observations in detail according to the state definitions in these works they may not show a one-to-one correspondence, such detail is irrelevant for this work. Moreover, until the late 1990’s (Homan & Belloni 2005), it was believed that the black hole state was determined by the instantaneous mass accretion rate $\dot{m}$, while in recent years ample evidence has been obtained that, in order to obtain a consistent classification, luminosity should not be a determinant. Hence the historic additions to the state names, associating the state with a certain luminosity regime or instantaneous mass accretion rate $\dot{m}$ (viz. high or low) have been omitted here.

When an XRB goes into outburst it usually goes through all the states listed below, before returning to quiescence. If we plot the X-ray hardness$^6$ and luminosity of a source during an outburst, we obtain a track through a hardness-intensity diagram (HID) in which each state is found to inhabit its own region (see Figure 1.8, showing qualitatively the behavior of most canonical

$^6$The hardness (or X-ray color) is the intensity-ratio of a harder and softer X-ray band, chosen so that it gives the relative importance of the thermal and power-law component in the spectrum. For most black holes (like GX 339–4; Chapter 4), the softer band is between 3-6 keV, while the harder band spans 6-10 keV. For GRS 1915+105 (Chapter 3) this system fails because the thermal (accretion disk) component is too hot, and hence for this source deviating, higher-energy bands are used (see caption of Figure 3.2).
black holes in the HID, or e.g. Figure 4.3 for a real HID for GX 339–4). In a nutshell, the physical picture thought to govern the flow through these states (starting from quiescence) is presumably similar to the following: The mass transfer progresses at a higher rate than it can be accreted onto the central object, so there is an outflow (in the form of a jet) present and the mass and temperature of the disk increase. At some point the viscosity in the disk can increase dramatically because e.g. the ionization temperature of hydrogen is reached. At that point the matter in the disk rapidly approaches the accretor and the central accretion rate and source luminosity rapidly increase, perhaps even exceeding the Eddington rate. The disk reaches all the way in and the X-ray spectrum softens while the jet is quenched and a disk wind subsumes the role of the outflow. After some weeks or months a significant fraction of the disk mass is accreted and as the disk cools accordingly the jet is launched again. Ultimately, when the disk has cooled enough the source returns to quiescence. (Fender 2009)

**Quiescent State**

Most X-ray transients spend most of their time in this low emission state (Fender 2009). The X-ray luminosity is typically $10^{30.5-33.5}$ erg s$^{-1}$ (or $L_X/L_{\text{Edd}} \lesssim 10^{-5}$ when normalized to the Eddington rate; Gallo 2009). The spectrum is hard and non-thermal, with a photon index of
1.5 \lesssim \Gamma \lesssim 2.1. The reduced emission in this state often allows for dynamical measurements of the system; if the companion star is no longer masked by flux from the accretion process and spectroscopic measurements can be done. Because of the similarities in spectra (indicating the presence of non-resolved jets) and PDSs, this state can be regarded as a low-luminosity version of the hard state (see below).

![Figure 1.9: Typical soft (red) and hard (blue) state spectra for Cyg X-1. (Gierliński et al. 1999)](image)

**Hard State (HS)**

This state is the most relevant for this work, as quite some evidence has been gathered, associating it with compact steady radio jets:

- Sources that have been in the hard state for several weeks are likely to show a correlation between the X-rays and the radio intensities \( (L_R \propto L_X^{\alpha}) \) (Corbel 2000; Corbel et al. 2003; Gallo et al. 2003) and an optically thick, flat-to-inverted radio spectrum. Russell et al. 2006 and Coriat et al. 2009 also found evidence for a similar correlation between the near infrared and X-ray bands. (also see Sections 2 and 4.1.2)

- Radio-jets have been been resolved in the hard X-ray state in some sources (e.g. Cyg X-1; Stirling et al. 2001), while the radio spectral index shows a flat/inverted spectrum.

- In GX 339–4 the radio observations showed a 2% linear polarization, while the position angle was nearly constant (Corbel et al. 2000).

When the XRB goes to the soft state (see below), a quenching of the radio emission is usually observed.

BHBs generally spend the majority of their time in the Hard State (HS), which is associated with a low accretion rate, and a hard X-ray power-law component with photon index \( 1.4 \lesssim \Gamma \lesssim 2.1 \) (typically comprising \( \gtrsim 80\% \) of 2-20 keV flux). This non-thermal radiation is believed to stem from (quasi-)thermal, outer-disk seed photons being Comptonized in the central part of the accretion flow, which could either be the inner part of the disk (the corona; see Figure 1.4), or the base of a jet. The PDSs show strong low-frequency band-limited noise (BLN). The accretion rate is typically less than a few percent of the Eddington rate, but the exact luminosity varies and often depends on the outburst history (see Section 4.1.1). The disk (if visible) temperature shows an appreciable decrease compared to the soft state, supporting the assumption of a truncated inner disk to a large radius \( R_{\text{in}} \sim 100r_g \). In some sources there may be an enhancement of the 20-100 keV flux, generally interpreted as the power-law component reflecting of the inner accretion disk (e.g. Di Salvo et al. 2001).

It should be noted here that in the literature the acronym HS is used for the High/Soft state and LS for the Low/Hard state. As we do not consider luminosity as a determinant for the canonical states we will use SS for the soft state (and HS for the hard state).
The project

Soft State (SS)

Because the X-ray spectrum in the soft state is dominated (typically \( \gtrsim 75\% \) of the 2-20 keV flux) by a soft multi-color blackbody coming from the accretion disk, it is also known as the thermal-dominated state. In this state the disk extends all the way to the ISCO. The power-law component is steep (\( 2.1 \lesssim \Gamma \lesssim 4.8 \)). There is only little rapid variability and Quasi-Periodic Oscillations (QPOs; semi-periodic behavior, evident as peaks in a PDSs) are absent or weak (rms\(<1\%\)).

Intermediate State (IMS)

In the IMS both power-law and disk components contribute appreciable to the X-ray spectrum. The variability in this state is the most complex of all, including the behavior in terms of (0.1-30Hz) QPOs. Belloni et al. (2005) subdivide the IMS into two distinct states; the Hard IMS (HIMS) and the Soft IMS (SIMS). The former shows moderately strong BLN, while the latter lacks this component, but shows power-law noise.

The IMS was first observed at a relatively high luminosity (in GX 339–4; see Chapter 4) and was consequently dubbed the Very High State (VHS; Miyamoto et al. 1991a). However since then a plethora of low luminosity observation have become available showing IMS characteristics. Generally the low luminosity IMS corresponds to the HIMS (but not always).

1.5 The project

As explained above, there are still a myriad of uncertainties when it comes to jets. Recently, however, a model (Markoff et al. 2005; hereafter MNW05) has become available that, amongst other things, has proven to be relatively successful in approximating the broadband spectrum for hard state XRBs. We describe the MNW05 model and its merits in Section 2. For this project we will explore three issues using this model:

1. Ever since the microquasar GRS 1915+105 was discovered, it is has accreted at near- or super-Eddington luminosities. The source exhibits none of the canonical black hole states, however it does exhibit a state that is associated with jets. This so-called plateau state has only been observed in this particular source and we will try to ascertain if a MNW05 jet is able to well-approximate this state in Chapter 3. If successful, we will be able to quantitatively compare the plateau state to the canonical HS.

2. Observations from a more usual XRB, GX 339–4, have shown that the canonical HS occurs at different luminosity levels, depending on if an outburst is increasing or decreasing in intensity. The track in the HID is therefore hysteretical in shape. In Chapter 4 we will see if a jet could provide an explanation for this hysteresis.

3. For Chapter 5 I will add some new physics into the jet-paradigm. In the MNW05 model the leptons are continuously accelerated into a power-law, starting at a distance \( z_{acc} \) from the jet base (see Section 2.1). The model considers the cooling rate due to synchrotron losses and compares this rate to the acceleration rate to determine the maximum energy the leptons can attain. However, according to Kardashev (1962), if cooling is significant a power-law particle distribution should show a cooling break at a certain energy, where the slope of the distribution steepens by unity. Hence, if if cooling is significant in the acceleration region, the SED predicted by the model should also change significantly. To investigate the significance of synchrotron cooling in a real BHB, we will apply the modified paradigm to GX 339–4 data from Chapter 4.

More detailed explanations of these issues and approaches to the problems are provided in the individual Sections mentioned above.
The outflow-dominated model

As explained in Section 1.4, historically the canonical states have been interpreted as due to fluctuations in the inner part of the accretion geometry. Particularly successful have been the paradigms involving a Comptonizing corona. However all these “classical” models consider only the X-ray emission and omit lower frequency wave bands, such as radio or IR. In order to meaningfully compare broadband data, for this work we will use a relatively new model (see Markoff et al. 2005; hereafter MNW05, and references therein) specifically designed for this purpose. In this paradigm the main contributions to the broadband radiation come from the jet in stead of the inflow. It is still the only model available that can fit the X-ray spectrum with the same precision as corona-only models (ibid; Nowak et al., in prep.) while also fitting the radio through IR bands from the same physical picture. In addition the MNW05 paradigm can analytically explain the slope of the Radio/X-ray correlation (see Section 1.4; Markoff et al. 2003a) as well as the spectral turnover between the IR and the X-rays (see Figure 2.1).

The MNW05 model has already been applied successfully to many Galactic sources in the “canonical” HS: the original paper features fits to Cyg X-1 and GX 339-4, while a different data set of the latter, along with observations of XTE J1118+480 are fit in Maitra et al. (2009). Further papers have explored applications to simultaneous broadband data sets from GRO J1655-40 and A0620-00 (Migliari et al. 2007; Gallo et al. 2007), with other sources in progress. The results of all these applications has been the discovery, perhaps not surprisingly, that the free parameters fit into similar ranges for all stellar mass sources.

Furthermore, as mentioned in Section 1.3 the recent discovery (Merloni et al. 2003; Falcke et al. 2004) of a Fundamental Plane of black hole accretion (see Figure 2.2) supports mass-scaling accretion physics from stellar to supermassive BHs, and thus would argue that the MNW05 should also apply to weakly accreting AGN. Confirming this, Markoff et al. (2008) successfully fit several spectral energy distributions (SEDs) from the supermassive BH M81* with parameters in the ranges found from BHBs.

2.1 Description of the model

In addition to the jet itself, our paradigm has two more components: a multi-color accretion disk (following Mitsuda et al. 1984 and Makishima et al. 1986) and a simple blackbody to account for the stellar companion (see Figure 2.3). This “standard” disk model represents a thin accretion disk with a temperature profile of \( T \propto R^{-3/4} \). The disk is divided into annuli with increasing surface area going from the inner to the outer edge. Each segment contributes blackbody radiation of a certain temperature and normalization, determined by how far removed an annulus is removed from the center of the disc and the surface area of the annulus, respectively. The disk is only a weak element in comparison to the entire spectrum, but it also offers soft seed photons for the
Figure 2.1: Schematic representation of the stratified spectrum resulting from the addition of the components emitted by idealized self-absorbed jet segments. Each subsequent jet segment grows in size with respect to its previous neighbor, due to adiabatic expansion, and the particle and magnetic energy densities decrease accordingly. Hence the peak frequency of each component decreases, resulting in a flat spectrum at the observed frequency $\nu_{\text{obs}}$ after superposition. The turnover point from optically thick to optically thin is also indicated ($\nu_{\text{SSA}}$; below this point synchrotron self-absorption is significant). We usually expect $\nu_{\text{SSA}}$ to be in the IR and there may be a link between the power-law from the IR to the X-ray once the jet becomes optically thin. (adapted from Markoff 2009)

Figure 2.2: The fundamental plane with results from Markoff et al. (2008) covering simultaneous radio/X-ray observation of the supermassive BH M81*, obtained with MNW05. The solid line indicates the best fit to the correlation, while the average scatter from the correlation is shown by the increasingly finer dashed lines. (Markoff 2009)
Compton processes (as does the stellar blackbody).

We show a schematic representation of the geometry of our jet in Figure 2.3 and summarize the physical parameters in Table 2.1. One of the most important of these parameters is the jet normalization \( N_j \). This factor scales with the accretion power at the inner edge of the accretion disk and represents the normalization of the power going into the jet (the exact power may be an order of magnitude larger, see MNW05). The power is equally divided between the internal pressure and the kinetic energy. We assume that the kinetic energy is carried by cold protons, while the leptons do the radiating. The energy involved in the internal pressure goes into the particle and magnetic energy densities, with a ratio determined by \( k \). \( k = 1 \) equals equipartition and higher values indicate magnetic dominance. The initial velocity at the base of the jet, or “nozzle” is the proper sound speed of an electron/proton plasma: \( \beta_s c \sim 0.4 c \). The radius of the

\[
\begin{array}{|c|c|}
\hline
\text{symbol} & \text{description} \\
\hline
N_j & \text{power going into jet (short of one magnitude, see MNW05) in Eddington} \\
\hline
r_0 & \text{radius of jet nozzle / base of jet in } r_g \\
\hline
T_e & \text{temperature of electrons entering the jet} \\
\hline
p & \text{particle distribution index} \\
\hline
k & \text{partition of erg. between magnetic and particle erg. densities} \\
\hline
p_{th} & \text{fraction of thermal distribution going into a power law} \\
\hline
z_{\text{acc}} & \text{location of particle acceleration front in } r_g \\
\hline
\epsilon_{\text{sc}} & \text{measure of acceleration rate (see Equation 2.5)} \\
\hline
r_{\text{in}} & \text{inner disk radius in } r_g \\
\hline
T_{\text{disk}} & \text{inner disk temperature} \\
\hline
r_{\text{out}} & \text{outer disk radius in } r_g \\
\hline
\end{array}
\]

Table 2.1: Free parameters in the MNW05 model and short description. More details on the individual parameters are in the appendix of Markoff et al. (2005).
The outflow-dominated model

jet-base is also a free parameter, \( r_0 \) (expressed in \( r_g \), or gravitational radii). The particles start of in a quasi-thermal, or (relativistic) Maxwellian distribution (see Equation 5.6 and Figure 5.3), the peak energy of which is determined by the electron temperature (\( T_e \)).

After the nozzle the jet is allowed to expand sideways freely, or adiabatically, at the initial speed causing a longitudinal pressure gradient. This pressure gradient leads to a moderate acceleration of the jet along the direction of flow; the resulting velocity profile (see Figure 2.4) is calculated from the Euler equations (e.g. Falcke 1996) and is roughly logarithmic, evening out at bulk speeds of \( \Gamma_j \sim 2−3 \). Starting at a segment of the jet at some distance \( z_{\text{acc}} \), also expressed in \( r_g \) from the base, we assume a percentage of the leptons is accelerated into a power-law, by some continuous mechanism we are as of yet unsure of (usually we assume diffusive, relativistic shock acceleration; e.g. Heavens & Drury 1988). This percentage is usually set at 75% from initial empirical results obtained with this paradigm. The slope of the power-law particle distribution is parametrized by the variable \( p \). In each jet segment after the first acceleration zone, the general shape of the distribution is assumed to remain the same (to avoid having to solve the complete Euler equation in every segment after the initial acceleration), while the total lepton density decreases according to the adiabatic expansion. This can be achieved if we assume a continuous injection of “fresh” energetic power-law distributed electrons \( N E dE \propto C E^{-p} dE \), where \( C \) is a normalization constant) in each segment after the acceleration front. Adding up the radiation produced by all the optically thick segments the “hallmark” flat radio jet spectrum is produced (see Figure 2.1). The jet is stopped at a distance \( z_{\text{max}} \) (in cm) from the base. This distance is usually set to a value large enough to sustain the flat/inverted slope through the radio data.

Next to the above free parameters there are several physical parameters that are fixed according to observations in the literature, e.g. the black hole mass, the inclination and the distance to the source or the effective temperature of the stellar companion.

### 2.2 Theory

Some example spectra of the model described above, including contributions of the different components, are shown in Figure 2.5. We already described in the caption of Figure 2.1 how the flat/inverted radio spectrum is created and that we expect a turnover point from optically thick to optically thin radiation between the IR and X-ray bands. We now further explore how the jet radiation in the highest energy bands. As mentioned in Section 1.2, the X-ray or \( \gamma \)-ray photons emitted by a jet can come from synchrotron radiation, inverse Compton (in the form of SSC of jet-base photons and Comptonization of seed accretion disk or stellar companion photons), or both, processes (Markoff et al. 2003b). In the jet the relative importance of these processes
are dependent on the maximum energy to which the particles can be accelerated. The maximum possible lepton energy is reached when the cooling rate matches the acceleration rate. Both rates are dependent on particle energy and local physical parameters (Markoff et al. 2001).

If we assume synchrotron cooling is dominant we can use the kinetic equation

$$\frac{\partial N}{\partial t} = -\frac{\partial}{\partial E}\left[\dot{E}N\right],$$

(2.1)

where $\dot{E}$ is given by the synchrotron energy losses, to describe the cooling. According to Rybicki & Lightman (1979), for an isotropic velocity distribution the angle-averaged power radiated by a single lepton, due to synchrotron radiation is

$$\dot{E} = P = \frac{4}{3} \sigma_T \beta^2 \gamma^2 U_B = \frac{4 \sigma_T U_B \beta^2 E^2}{3 m_e^2 c^3}$$

(2.2)

where $\sigma_T$ is the Thomson cross section, $m$, $\beta$ and $E$ are the lepton mass speed and energy respectively, $U_B = B^2/8\pi$, with $B$ the magnetic field strength, is the magnetic energy density.

Separating the variables we can integrate, setting the initial energy to $E_0$, to obtain the energy of a single relativistic lepton as a function of time

$$E(t) = \frac{E_0}{1 + t/\tau},$$

(2.3)

where

$$\tau = t_{\text{syn}} = \frac{3 m_e^2 c^3}{4 \sigma_T U_B \beta^2 E_0}$$

(2.4)

is the time scale of the energy loss, or the synchrotron lifetime. $t_{\text{syn}}^{-1}$ defines the cooling rate due to synchrotron losses.

On the other hand, the acceleration rate in the case of “classic” Fermi acceleration is (Jokipii 1987)

$$t_{\text{acc}}^{-1} = \frac{3}{4} \left(\frac{u_{sh} c}{c}\right)^2 \frac{c e B \epsilon_{sc}}{E \beta^2},$$

(2.5)

where $e$ is the lepton charge, $u_{sh}$ is the shock speed in the plasma frame (which, due to a lack of current understanding of the exact shock propagation mechanism, is fixed to a value $\beta_{sh} = 0.6$ in our jet) and the plasma parameter $\epsilon_{sc}$ (also a parameter in the model, see Table 2.1) is a measure of the acceleration efficiency, as it scales with the acceleration rate.

If we solve $t_{\text{acc}}^{-1} = t_{\text{syn}}^{-1}$ we get the maximum energy $E_{\text{max}}$ that a power-law lepton can attain (if synchrotron cooling is dominant).

$$E_{\text{max}} = \frac{3}{\beta^2} \frac{u_{sh} m_e c}{\epsilon_{sc}} \left(\frac{\pi e \epsilon_{sc}}{2 \sigma_T B}\right)^{0.5}$$

(2.6)

This energy determines the energy of the exponential synchrotron cut-off (around 10 keV in Figure 2.5) in the spectrum $\nu_{\text{max}}$ via the critical synchrotron frequency: $\nu_{\text{max}} = 0.29 \nu_c$, where $\nu_c \approx \frac{1}{8\pi} E_{\text{max}}^2 (eB)/(m_e^3 c^3)$ (Rybicki & Lightman 1979), so that

$$\nu_{\text{max}} \propto \epsilon_{sc} \left(\frac{u_{sh}}{c}\right)^2,$$

(2.7)

which depends on the acceleration efficiency, but clearly not on the magnetic field, jet power or location of the acceleration front. If we observe radiation above the synchrotron cut off it must be IC.

It should be noted that the above result is only true if synchrotron cooling is dominant, while in out jet also other cooling process are at work (adiabatic an IC losses), so that $E_{\text{max}}$ may in
The outflow-dominated model

fact be dependent on the magnetic field. We investigate the cooling in the jet in more detail in Chapter 5.

It should also be noted that while in Markoff et al. (2001) the values for plasma parameter $\epsilon_{sc}$ were expected to lie in the range $3.6 \times 10^{-2(2-3)}$ (for $\beta_{sh} = 0.6c$, see above and Jokipii (1987)), consistent with "classical" Fermi acceleration, since MNW05 this value is allowed to move to lower values to account for other possible acceleration processes. For our purposes, the parameter $\epsilon_{sc}$ is used as a fudge factor to account our lack of knowledge about the true acceleration physics. If patterns appear to emerge we may at a later stage be able to use this parameter to form an idea about the true acceleration process at work, and perhaps get estimates for the dependency of the acceleration rate on physical parameters.

\[ \text{synchrotron} \quad \text{post-shock} \]
\[ \text{synchrotron} \quad \text{pre-shock} \]
\[ \text{accretion disk} \]
\[ \text{inverse Compton} \]
\[ \text{total} \]
\[ \text{iron line} \]

Figure 2.5: Example fit/spectrum made for GX 339–4 (top; Chapter 4) and GRS 1915+105 (bottom; Chapter 3) using MNW05 and the software package ISIS (Interactive Spectral Interpretation System) (Houck & Denicola 2000). The individual components from the MNW05 model comprising the broadband spectrum also shown: The light-green dashed curve is the pre-shock synchrotron contribution. The dark-green dash-dotted curve is the post-shock synchrotron. The purple dotted curve represents the thermal contributions from the accretion disk (or, in the case of GRS 1915+105 the thermal contributions from a stellar blackbody below $\sim 10^{-2}$ keV and an accretion disk). The orange dashed-dotted line above $\sim 10^{-3}$ keV represent the Compton-upscattered stellar blackbody and accretion disk seed photons and the Synchrotron Self-Comptonization (SSC) of the pre-shock synchrotron. The solid grey line is the total MNW05 model spectrum originating from the jet, the accretion disk and the companion, however not forward-folded through the detector response matrices and without iron line or reflection contributions or absorption due to the interstellar hydrogen column density or the smedge model (see Section 3.4 and 4.4). The red solid lines through the data points shows the model flux including all these features. The iron line complex, modeled by a Gaussian is shown by the thick green curve near 6.4 keV.
The black hole binary GRS1915+105 and the plateau state

GRS 1915+105 is a very peculiar black hole binary, exhibiting accretion-related states that are not observed in any other stellar-mass black hole system. One of these states, however—referred to as the plateau or $\chi$ state—may be related to the canonical hard state of black hole X-ray binaries. Both the plateau and hard state are associated with steady, relatively lower X-ray emission and flat/inverted radio emission, that is sometimes resolved into compact, self-absorbed jets. In order to investigate the relationship between the plateau and the hard state, we fit two multi-wavelength observations using a steady-state outflow-dominated model, developed for hard state black hole binaries. The data sets consist of quasi-simultaneous observations in radio, infrared and X-ray bands. Interestingly, we find both significant differences between the two plateau states, as well as between the best-fit model parameters and those representative of the hard state. In particular we find that the plateau state jets must be not only extremely luminous, but also extremely magnetically dominated. We discuss our interpretation of these results, and the possible implications for GRS 1915+105’s relationship to canonical black hole candidates.

3.1 Introduction

The microquasar GRS 1915+105 was discovered on 15 August 1992, by the WATCH all-sky monitor, aboard the Russian GRANAT satellite (Castro-Tirado et al. 1992, 1994). It is a hard X-ray transient located in the constellation of Aquila, at $l = 45.37^\circ$, $b = -0.22^\circ$, and was the first stellar mass accreting black hole binary (BHB) discovered to exhibit superluminal velocities in its radio emitting-ejecta (Mirabel & Rodriguez 1994). The obvious parallels to the jets in Active Galactic Nuclei (AGN) led to this source being coined the first “microquasar”. Subsequent observations with instruments onboard the Rossi X-ray Timing Explorer (RXTE) have revealed a richness in variability distinguishing GRS 1915+105 from every other known BHB, over which astronomers are still puzzling to this day.

The longer term X-ray variability, or “dipping”, is thought to be associated with the disappearance and
regeneration of the inner accretion disk, perhaps caused by the onset of thermal-viscous instabilities (Belloni et al. 1997b,a). Other models have interpreted the spectral changes as resulting from the disappearance of the corona (Chaty 1998), or from the dissipation of magnetic energy via magneto-hydrodynamical plasma processes (e.g. Tagger et al. 2004).

Beyond the dipping behavior, Belloni et al. (2000) were able to classify all variability patterns stretching over more than a year into only twelve classes (Klein-Wolt et al. 2002 and Hannikainen et al. 2005 later identified two more classes), based on their color-color diagram and light curves. Ten of these twelve classes can be understood as the interplay of two or three of three basic states, designated as state A, B and C. The remaining two classes, $\phi$ and $\chi$, do not show state transitions and appear exclusively within states A and C respectively. State A displays the highest flux and the softest spectrum while state C displays the lowest flux and the hardest spectrum. Although state B never lasts for more than a few hundred seconds, $\phi$ and $\chi$ episodes can persist for days or short intervals of $< 100$ s.

Aside from the existence of so many distinct accretion states, GRS 1915+105 appears similar to other BHBs. Thus there has been much discussion (e.g. Reig et al. 2003) about the extent to which any of the states found in GRS 1915+105 correspond to the "canonical" states (see Section 1.4) found in the average BHB.

Similar to the other BHBs, GRS 1915+105 also displays periods of relatively hard, steady X-ray emission, but in contrast to the HS, only for about half the observation time (Trudolyubov 2001). Belloni et al. (2000) identify these intervals of decreased variability with long C state episodes. The subset of state C/ class $\chi$ observations that are particularly radio-bright are referred to as the plateau state (ibid.), but elsewhere have variously been referred to as the radio-loud, radio-plateau or radio-loud low/hard X-ray state (Muno et al. 2001), the type II hard steady state (Trudolyubov 2001) and $\chi_{RL}$ (Naik & Rao 2000).

The plateau state was first described by Foster et al. (1996) and its description was later more refined in Fender et al. (1999). The state can assert itself in a period as short as a day and is characterized by a decrease in X-ray flux density and an increase in radio flux density to a steady value, typically $\sim 10 - 100$ mJy. Ample evidence supports the fact that the plateau state X-ray and radio luminosities are indeed correlated, with an increasing delay from X-ray to infrared (IR) to radio emission (e.g. Klein-Wolt et al. 2002). As in the HS the plateau radio spectrum is optically thick, and AU-scale self-absorbed compact jets have been spatially resolved using VLBI (Dhawan et al. 2000). Another marker of the plateau state comes from timing analyses, where 1-10 Hz type-C quasi-periodic oscillations (QPOs) are present (Rao et al. 2000). The exact frequency of these QPOs appears to be inversely correlated with the radio flux (Rau & Greiner 2003).

In line with these arguments, it is tempting to identify the plateau state as GRS 1915+105’s analog of the HS. However, although the plateau state and the HS share many similarities, it does display some distinct properties that cannot be ignored. For instance, while the BHBs in the HS usually have a luminosity of $\lesssim 10\% L_{Edd}$, the average luminosity observed in the plateau state is $\sim L_{Edd}$. Moreover, the plateau X-ray photon index is never very hard, with $\Gamma \sim 1.8 - 2.5$ (Fender & Belloni 2004). Finally, Reig et al. (2003) argue that the canonical HS is never seen in GRS 1915+105 because a power-law tail is always present in the plateau state X-rays. The origin of such tails in the softer states of BHBs is still under debate, but it is not generally associated with the HS.

In this work, we seek to make a more quantitative comparison between the plateau state found in GRS 1915+105 and the HS found in more typical black holes, in the context of an outflow-dominated model that has successfully described broadband data from several BHBs. In Section 3.2 we describe the multi-wavelength data and the reduction methods. We discuss what physical parameters we use for modeling and why in section 3.3 and the model itself, together with the results in Section 3.4. We put our work into context with the findings of others in Section 3.5, before drawing our final conclusions in Section 3.6.
3.2 Observations and data reduction

Based on the lightcurves and its position in the Hardness Intensity Diagram (HID) diagram (see Figures 3.2 & 3.3), GRS 1915+105 was in a class $\chi$ state on July 8th 1999 (dataset A) and April 13th 2005 (dataset B). The presence of a type-C QPO on both occasions (see Figure 3.2, right panel) corroborates this fact. During the 1999 episode, quasi-simultaneous radio and infrared observations were performed using UKIRT and GBI. The 2005 dataset also holds observations of

| Dataset A 1999 July 8 / MJD 51367 |
|-------------------|-----------------|---------------|-----------|
| Band     | Instrument | MJD start     | UTC       | Duration  |
| X-ray    | RXTE       | 51367.2737    | 06:34:13   | 14435 sec |
| IR       | UKIRT      | (H) 51367.4821| 11:34:10   | 80 min    |
|          |            | (K) 51367.8216| 07:56:59   | 64 min    |
| Radio    | GBI        | 51367.371     | 09:02:53   | 30 sec    |

| Dataset B 2005 April 13 / MJD 53473 |
|-------------------|-----------------|---------------|-----------|
| Band     | Instrument | MJD start     | UTC       | Duration  |
| X-ray    | RXTE       | 53473.054     | 01:09:05   | 6769 sec  |
| IR       | CTIO       | (K) 53472.3216| 07:43:06   | 25 min    |
| Radio    | Ryle       | 53473.098     | 02:21:07   | 219 min   |

Table 3.1: Observations included in dataset A and B, with RXTE ObsID 40403-01-09-00 and 90105-07-02-00 respectively. The $H$ and $K$ bands, obtained using the Cooled Grating Spectrometer at UKIRT, are at 1.455-2.094 and 1.906-2.547 $\mu$m, respectively ($J$ band measurements were also done, but these are unusable due to high interstellar extinction). The $K$ band observation in dataset B was done (more than half a day before ObsID 90105-07-02-00) with the Andicam at CTIO. The GBI measurements are 2.25 & 8.3 GHz and the flux densities used are 30 sec vector scan averages. The Ryle telescope operates at 15 GHz and the flux density used is the average of $\sim$13 ks worth of five minute integrations.

Figure 3.2: Left Panel: HID comparing our observations (closed circles) to the plateau states from Belloni et al. (2000) (data courtesy of T. Belloni; open stars are $\chi_1$, closed stars are $\chi_3$). The soft, higher-luminosity red dot is ObsID 90105-07-02-00, the lower is ObsID 40403-01-09-00. The plot shows the Crab corrected hardness and flux, calculated by taking the ratio and the sum of the RXTE PCU2 rates in two bands: channel 0-35 ($\sim 2 - 15$ keV) and 36-49 ($\sim 16 - 21$ keV). Each band is divided by the Crab rate in the same band, on the same day, simulating a photon index of 2.1 and utilizing a normalization of 10 photons cm$^{-2}$ s$^{-1}$ keV$^{-1}$ at 1 keV. Right panel: Normalized (following Leahy et al. 1983) power density spectra showing the type-C QPOs in the X-rays for both datasets, indicating the observations are indeed plateau state. (Data courtesy of P. Soleri)
3.2.1 X-Ray: RXTE data reduction

We use data from two instruments on board the RXTE: PCA (Jahoda et al. 2006) and HEXTE (Rothschild et al. 1998). The X-ray data have been reduced using HEASOFT, version 6.3.1., applying standard extraction criteria: a pointing offset of $< 0.01^\circ$ from the nominal source position, and a source elevation of $> 10^\circ$. The exclusion time for the South Atlantic Anomaly (SAA), however, was only 10 min, as GRS 1915+105 is a very bright source. For the same reason we disregarded the top layer and imposed a maximum “electron ratio” of 2, to take into account contamination by source photons. All HEXTE data products were deadtime corrected. Correction for the PCA deadtime was also carried out; PCU 4 was inactive for half of the 1999 observation, while during the other half, both PCU 1 and 4 were inactive. The 2005 observation was done with PCU 0 and 2 only. Due to uncertainties in the first four PCA channels, we only include PCA data above 3 keV. Since HEXTE provides reliable data for energies $\gtrsim 20$ keV we ignored the PCA data above 22 keV. For the PCA, standard2f mode data were used. The PCA background was modeled using the pca_bkgd_cmbrightv1e_efv20051128 model. Only HEXTE data above 20 keV were used due to the uncertainty of the response matrix below these energies. At high energies, the spectrum was ignored above 200 keV.

3.2.2 IR: UKIRT and CTIO

Details on the UKIRT data reduction may be found in the original publication Harlaftis et al. (2001) Because of uncertainties in the dereddening of the data, we set the IR error bars to 10 %. The UKIRT IR spectrum included in dataset A was also re-binned (see Figure 3.11), because the spectral resolution is well beyond what we need for continuum fitting (see Section 3.4).
The 12.40 mag K band point in dataset B was taken using the CTIO 1.3m telescope using the ANDICAM detector (Neil et al. 2007). It was dereddened assuming a hydrogen column density of $N_H = 4.7 \times 10^{22}$ cm$^{-2}$ (Chaty et al. 1996; see below). Using (as in Harlaftis et al. 2001) $E(B-V)=9.6$ and $R_K=0.32$ (Howarth 1983), we calculate a total K band absorption $A_k = 3.072$ mag. Setting the zero point flux to 667 Jy (Cox 2000) and the error (again) to 10 %, this yields an unabsorbed flux of 123.9 ± 12.4 mJy.

### 3.2.3 Radio: GBI and Ryle Telescope

Dataset A includes two GBI data points and dataset B includes a data point from Ryle Telescope. Lightcurves from these telescopes are shown in Figure 3.3.

For the GBI data, a flux density calibration procedure similar to that reported in Waltman et al (1994) has been employed here. The flux densities of 0237-233, 1245-197, and 1328+254 were determined using observations of 1328+307 (3C 286). The flux density of 3C 286 was based on the scale of Baars et al. (1977), and the assumed values were 11.85 Jy at 2.25 GHz and 5.27 Jy at 8.3 GHz. On MJD 51367.371 the GBI flux densities were 25 ± 4 mJy at 2.25 GHz and 30 ± 6 mJy at 8.3 GHz. Similar behavior was shown during a plateau state of April 2000 (Ueda et al. 2002).

The Ryle Telescope operates at 15 GHz and it observed GRS 1915+105 (simultaneously with ObsID 90105-07-02-00) from MJD 53473.098 to MJD 53473.417. The data were reduced following Pooley & Fender (1997); observations of Stokes I+Q were interwoven with those of a nearby phase calibrator (B1920+154). The flux-density scale was set by reference to 3C48 and 3C286, and should be consistent with that defined by Baars et al. (1977).

### 3.3 Constraints on input physical parameters

As mentioned in Chapter 2, we determine values for model physical parameters from observations and freeze them for the duration of the fitting process, whenever possible. In the following subsections, we first discuss how we obtained these values (presented in Table 3.2), and then briefly summarize the models used and our fitting methods before presenting our results.

#### 3.3.1 Distance and hydrogen column density

The distance to GRS 1915+105 is still a matter of some debate. A first estimate was derived in Mirabel & Rodriguez (1994). From a core ejection they find a maximum distance of 13.7 kpc, under the assumption that the ejection was intrinsically symmetric. Rodriguez et al. (1995) attempted to derive a more accurate distance estimate by determining the kinematic distance from 21 cm absorption spectra of atomic hydrogen along the line of sight to GRS 1915+105 during a radio outburst, and find that it could be as far away as 12.5±1.5 kpc. Later measurements of the $^{12}$CO($J = 1-0$) spectrum by Chaty et al. (1996) are consistent with this distance. They also constrain the visual extinction to be $A_v = 26.5 \pm 1.7$ mag, corresponding to a total hydrogen column density of $N_H = 4.7 \pm 0.2 \times 10^{22}$ cm$^{-2}$.

If GRS 1915+105 resides at 12.5 kpc, it must be accreting above the Eddington limit ($L_{\text{X-ray}} \sim 2.9 \times 10^{39}$ erg s$^{-1}$; McClintock & Remillard 2003, giving $L_{\text{X-ray}}/L_{\text{Edd}} \sim 1.5$, using a mass of 14 $M_\odot$; see next section). A slightly lower maximum distance of 11.2±0.8 kpc, derived from proper motion studies of radio ejecta during four major events (Fender et al. 1999), would still give a ratio $L_{\text{X-ray}}/L_{\text{Edd}} \sim 1.3$.

However, this distance could be too large. Chapuis & Corbel (2004) measure $^{12}$CO($J = 1-0$) velocity spectra of clouds along the line of sight, as well as of two nearby HII regions, and re-evaluate the hydrogen column density to $N_H = 3.5 \pm 0.3 \times 10^{22}$ cm$^{-2}$. This column gives a visual extinction of $A_v = 19.6 \pm 1.7$ mag, reducing the number intervening molecular clouds, and arguing for a smaller distance to GRS 1915+105 of 9.0±3.0 kpc (at 9 kpc the accretion rate estimate drops below Eddington), based on evidence that GRS 1915+105 lies behind the molecular hydrogen cloud.
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G 45.45+0.06. This cloud is at $\sim 7$ kpc (Feldt et al. 1998), however Chapuis & Corbel (2004) suggest that this cloud belongs to a huge complex, located at a distance of 6 kpc, providing the above lower limit to the distance. However, the high column density required (Mirabel & Rodríguez 1994; Rodríguez et al. 1995) suggests that another molecular cloud should be between us, therefore Fender et al. (1999) adopted the conservative distance of 11 kpc.

Considering the assumptions made in Chapuis & Corbel (2004) and the resulting size of the errorbars, we have decided to follow Fender et al. (1999) in adopting a more conservative distance of 11 kpc for this work. This distance is still consistent with the larger hydrogen column density of $4.7 \times 10^{22}$ cm$^{-2}$ from Chaty et al. (1996), that we use in our fits, as this value allows for enough hydrogen clouds to exist between us and the object to account for the observed column density.

Because a distance of 11 kpc implies super-Eddington luminosities for GRS 1915+105, we also fit the data using the minimum distance of 6 kpc, in order to explore the importance of distance on the modeling conclusions.

3.3.2 GRS 1915+105 Black Hole Mass

In order to constrain the mass of the compact object in a binary via the mass function

$$f(M) = \frac{PK_2}{2\pi G} = \frac{M_X \sin^3 i}{(1 + q)^2},$$

(3.1)

the mass-ratio of the compact source to the donor $q = M_d/M_X$ (or the mass of the donor), the orbital period $P$, the half-amplitude of the velocity curve of the donor $K_2$ and the inclination $i$ must all be known. Greiner et al. (2001a) found an orbital period $P_{\text{orb}}$ of 33.5 days and a velocity amplitude $K = 140 \pm 15$ km s$^{-1}$, giving a mass function $f(M) = 9.5 \pm 3.0$ M$_\odot$. Assuming the binary plane is the same as that of the accretion disk, the orbital inclination can be estimated from the orientation of the jets. As no constant precession has been observed, the jets can be assumed perpendicular to the accretion disk and orbital plane. The exact inclination is however still open to debate, as it is determined from the brightness and velocities of both the approaching and receding ejecta. Using

$$\theta = \tan^{-1} \left[ 1.16 \times 10^{-2} \left( \frac{\mu_a \mu_r}{\mu_a - \mu_r} \right) D \right],$$

(3.2)

where $\mu_a$ and $\mu_r$ are the velocities of the approaching and the receding blobs respectively, and a distance $D$ of 11 kpc, Fender et al. (1999) calculated an inclination of $i = 66^\circ \pm 2^\circ$. Harlaftis & Greiner (2004) were able to deduce a mass ratio $M_d/M_X = 0.053 \pm 0.033$ from the rotational broadening of photospheric absorption lines of the donor star. Together these values give a mass for GRS 1915+105 of $14.0 \pm 4.4$ M$_\odot$ (Greiner et al. 2001a). We use these as our fiducial values for the distance and mass. For the 6 kpc distance we have to recalculate the inclination and the mass: Using equation 3.2 (and again the velocities from Fender et al. 1999) we find a smaller inclination of $i = 50^\circ$, and, using equation 3.1 (and again the mass ratio from Harlaftis & Greiner 2004) an increased black hole mass of $23M_\odot$.

3.3.3 Properties of the donor star

A rough identification of the companion was performed by Greiner et al. (2001b), by analyzing absorption line features in the near infrared (NIR). They conclude that the companion is a class III K-M giant and compare it successfully to a K2 III star. If the star is indeed a K2 III giant, its temperature should be $T_d = 4455 \pm 190$ K (and its mass should be $M_d = 1.2 \pm 0.2$ M$_\odot$; Alonso et al. 1999). We use this temperature for the simplistic blackbody spectral component representing the companion in our spectral fits.
Table 3.2: Fixed physical parameters used in the fitting process, obtained from the literature. We omit the error bars, as they are not used in the fits.

<table>
<thead>
<tr>
<th>parameter</th>
<th>value</th>
<th>units</th>
<th>Reference</th>
</tr>
</thead>
<tbody>
<tr>
<td>column density</td>
<td>4.7</td>
<td>$10^{22}$ cm$^{-2}$</td>
<td>a</td>
</tr>
<tr>
<td>mass</td>
<td>14 and 23</td>
<td>M$_\odot$</td>
<td>b</td>
</tr>
<tr>
<td>inclination</td>
<td>66 and 50</td>
<td>$^\circ$</td>
<td>c</td>
</tr>
<tr>
<td>distance</td>
<td>11 and 6</td>
<td>kpc</td>
<td>c, d</td>
</tr>
<tr>
<td>donor temperature</td>
<td>4455</td>
<td>K</td>
<td>e</td>
</tr>
</tbody>
</table>


The mass ratio found by Harlaftis & Greiner (2004) implies a companion mass of $M_d = 0.81$ M$_\odot$. For stars of this mass range, the mass loss rate $\dot{m}$ by stellar winds cannot account for the high accretion luminosity found in GRS1915+105, confirming the accretion proceeds via Roche-lobe overflow (Greiner et al. 2001b).

3.4 Modeling and results

3.4.1 Spectral fitting method

For all spectral fits we use the program ISIS (Houck & Denicola 2000) compiled with XSPEC version 12.3.1x libraries (Arnaud 1996). The model discussed below is forward-folded through the X-ray detector response matrices, but applied directly to the radio through IR data. To account for the additional uncertainties in the PCA response matrix, a 0.5% systematic error has been added in quadrature to all PCA data. As the relative calibration of the PCA and HXTE instruments is not certain, the normalization factor for the PCA data is set to unity and tied to the radio and IR normalizations during the fits, while the HEXTE data normalization is left unfrozen.

As the focus of our paper is on modeling the non-thermal spectrum, and the exact stellar model needed for GRS 1915+105 is uncertain, we use a simple blackbody to model the IR, with the temperature fixed as discussed in the previous section. This component serves to account for the excess IR/optical flux level due to the star, so that the overall model normalization is correct.

All fits are done using the following components: (1) The MNW05 steady-state outflow-dominated model that includes a multi-color blackbody accretion disk and a single blackbody for the companion star; (2) an additive Gaussian line profile, with a line energy left free to vary between 6 and 7 keV, and line width $\sigma$ free to vary between 0 and 2 keV; Models (1)+(2) are either convolved with Compton reflection from a neutral medium (reflect; Magdziarz & Zdziarski 1995) or multiplied with the smeared edge model (smedge; Ebisawa 1991; using the "standard" index for the photoelectric cross-section of $\sim −2.67$), that allows for relativistic lines, and multiplied with a photo-electric absorption model (phabs) to account for the interstellar medium. As the strength of an absorption feature (or edge) is related to the strength of the according emission line, the line width is in principle also related to the absorption edge width. Thus when using the smedge model, we tie the width of the edge to the width of the Gaussian. For the reflect model we assume the viewing angle to correspond to the jet inclination. Furthermore, although the reflect model already includes an absorption edge, we also tried fitting the data with both the smedge and reflect models, to see what effect the inclusion of both a relativistic line and and Compton reflection has on the statistics.

3.4.2 Spectral fitting and results

After some initial experimenting, including either the smedge or reflect model, or both, we arrived at a set of recurring ranges in the models with satisfactory initial statistics for a broadband data set ($\chi^2_{\text{red}} \sim 2 − 3$). To quantify which model is best suited, we homogeneously fit the data
Table 3.3: Parameter ranges found in canonical black holes GX339-4 (Markoff et al. 2005; Maitra, Markoff & Brocksopp 2009), XTE J1118+480 (Maitra, Markoff & Brocksopp 2009), Cyg X-1 (Markoff et al. 2005), GRO J1655-40 (Migliari et al. 2007) and A0620-00 (Gallo et al. 2007), during the HS and best-fit parameters for the GR15+105 plateau state. The error bars have been resolved at 90 percent confidence level. We failed to resolve the error bars for parameters listed in italics. \(N_j\) is the jet normalization, \(r_0\) the nozzle radius, \(T_e\) the temperature of the leptons as they enter the jet, \(p\) the spectral index of the radiating particles, \(k\) the ratio between magnetic and electron energy densities, \(z_{acc}\) location of the particle acceleration region in the jet, \(f_{scat}\) the luminosity and temperature of the accretion disk, \(\sigma_{HT}\) the relative normalization between HEXTE and PCA, \(A_{line}\), \(E_{line}\) and \(\sigma\) are the normalization, the energy and width of the iron line feature, and \(E_{edge}\) and \(\tau_{max}\) are the energy and the normalization of the model iron line edge (the iron edge width is the same as that of the Gaussian) and \(\Omega/2\pi\) is the reflection fraction from the reflect model.

<table>
<thead>
<tr>
<th>variable</th>
<th>units</th>
<th>derived in canonical</th>
<th>BHs</th>
<th>MJD 51367</th>
<th>GRS1915+105</th>
<th>MJD 53473</th>
</tr>
</thead>
<tbody>
<tr>
<td>(N_j)</td>
<td>(10^{-1}L_{Edd})</td>
<td>0.0034 - 0.71</td>
<td>4.88(\pm)0.30</td>
<td>5.65(\pm)0.01</td>
<td>10.00(\pm)0.01</td>
<td>2.22(\pm)0.01</td>
</tr>
<tr>
<td>(r_0)</td>
<td>(GM/c^2)</td>
<td>3.5 - 20.2</td>
<td>20.4(\pm)1.1</td>
<td>20.3(\pm)0.1</td>
<td>6.49(\pm)0.00</td>
<td>4.94(\pm)0.06</td>
</tr>
<tr>
<td>(T_e)</td>
<td>(10^9 K)</td>
<td>2.0 - 5.23</td>
<td>0.94(\pm)0.03</td>
<td>1.00(\pm)0.01</td>
<td>0.80(\pm)0.01</td>
<td>0.90(\pm)0.00</td>
</tr>
<tr>
<td>(p)</td>
<td></td>
<td>2.1 - 2.9</td>
<td>2.27(\pm)0.02</td>
<td>2.14(\pm)0.02</td>
<td>1.87(\pm)0.00</td>
<td>2.57</td>
</tr>
<tr>
<td>(k)</td>
<td></td>
<td>1.1 - 7</td>
<td>557(\pm)57</td>
<td>225(\pm)1</td>
<td>286(\pm)0.1</td>
<td>95.2(\pm)1.6</td>
</tr>
<tr>
<td>(z_{acc})</td>
<td>(GM/c^2)</td>
<td>7 - 400</td>
<td>6612(\pm)54</td>
<td>4844</td>
<td>6612(\pm)54</td>
<td>7431</td>
</tr>
<tr>
<td>(\epsilon_{sc})</td>
<td>(10^{-4})</td>
<td>1.6 - 299</td>
<td>3.1(\pm)0.5</td>
<td>2.2(\pm)0.1</td>
<td>0.59(\pm)0.02</td>
<td>1.9</td>
</tr>
<tr>
<td>(r_{in})</td>
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<td>2.55(\pm)0.39</td>
<td>2.79(\pm)0.02</td>
<td>4.43(\pm)0.05</td>
<td>0.63(\pm)0.07</td>
</tr>
<tr>
<td>(L_{disk})</td>
<td>(10^{-1} L_{Edd})</td>
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<td>1.07</td>
<td>1.14</td>
<td>1.43</td>
<td>0.73</td>
</tr>
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<td>(A_{HEXT})</td>
<td>keV</td>
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<td>0.85(\pm)0.05</td>
<td>0.83(\pm)0.01</td>
<td>0.69(\pm)0.01</td>
<td>0.88(\pm)0.02</td>
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<tr>
<td>(\Omega/2\pi)</td>
<td></td>
<td></td>
<td></td>
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<td></td>
<td></td>
</tr>
<tr>
<td>(E_{line})</td>
<td>keV</td>
<td>-</td>
<td>35.7</td>
<td>77(\pm)3</td>
<td>49.1</td>
<td>30(\pm)0.1</td>
</tr>
<tr>
<td>(\sigma)</td>
<td>keV</td>
<td>-</td>
<td>6.30(\pm)0.10</td>
<td>6.00(\pm)0.11</td>
<td>6.06(\pm)0.03</td>
<td>6.32(\pm)0.07</td>
</tr>
<tr>
<td>(\chi^2/DoF)</td>
<td></td>
<td></td>
<td>173/149</td>
<td>222/150</td>
<td>215/149</td>
<td>155/150</td>
</tr>
</tbody>
</table>

\[\dagger\] Value is pushing lower or upper boundary and is therefore not well constrained.

\[b\] Derived from model values for \(r_{in}\) and \(T_{disk}\).
Figure 3.4: Multi-wavelength (left) and X-ray band only (right) best fit ($\chi^2_{\text{red}} \sim 1.16$) model spectrum (cf. Table 3.3, smedge column), using MNW05+Gaussian+smedge.

Figure 3.5: Multi-wavelength (left) and X-ray band only (right) best fit ($\chi^2_{\text{red}} \sim 1.48$) model spectrum (cf. Table 3.3, reflect column), using MNW05+Gaussian+reflect.

Figure 3.6: Multi-wavelength (left) and X-ray band only (right) best fit ($\chi^2_{\text{red}} \sim 1.44$) model spectrum (cf. Table 3.3, smedge2 column), also employing MNW05+Gaussian+smedge, but with a reduced magnetic dominance ($k \sim 30$ in stead of $\sim 550$, see table 3.3).
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Figure 3.7: Multi-wavelength (left) and X-ray band only (right) best fit result ($\chi^2_{\text{red}} \sim 1.01$) spectrum (cf. Table 3.3, 6kpc column), for the 1999 data set, using MNW05+$\text{Gaussian}+\text{smedge}$ and employing a fixed distance of 6 kpc.

from the same starting point, i.e. using the obtained initial ranges, we start the fitting routine with a high tolerance, and progressively reduce the tolerance, logarithmically halving it every step. For both data sets we perform this procedure, with and without the infrared, for all three model combinations. The $\text{smedge}$ model emerged as statistically favored in all attempts. Adding a smedge line to the $\text{reflect}$ model improves the statistics of a reflection model somewhat, but as mentioned above, this is not truly self-consistent and the improvement appeared hardly significant. Therefore, we proceed to fit the data in greater detail with the $\text{smedge}$ model.

The results of all fits are presented in Figures 3.4–3.10, with respective parameters listed in Table 3.3. Below we discuss the individual fits to both datasets in more detail.

We fix the jet length to $10^{16}$ cm, sufficient so that the slope through the radio data points is continuous. The spectral index of the optically thick radio-NIR synchrotron emission is determined in part by the internal jet plasma parameters such as the electron temperature, but it is most sensitive to the Doppler beaming factor (calculated from the inclination of the jets). The closer the jet axis aligns with the line of sight, the flatter the observed spectrum.

All fits to the two GRS 1915+105 plateau state datasets reveal a key difference when compared to canonical BHBs. In order to avoid too much excess in the IR over the companion star BB, and to provide the best IR fit, the initial particle acceleration region in the jets $z_{\text{acc}}$ needs to be at a distance of at least $\sim 2000 \, r_g$. Therefore the post-acceleration region in the jets does not dominate the X-ray emission below $\sim 10$ keV, in contrast to canonical BHBs in the hard state where acceleration is found to begin on the order of $10s$ of $r_g$.

In general we obtain values for electron energy index $p \sim 2.2$ that result in a significant synchrotron contribution to the X-rays below $\sim 50$ keV, with a ratio of synchrotron to inverse Compton flux of $\sim 0.1$. The very soft X-ray flux above $\sim 20$ keV, compared to HS BHBs, is best fit by a dominant inverse Compton contribution from the base of the jets. With a much harder value of $p$ we could conceivably fit more of the X-ray emission via synchrotron emission, but only if the exponential decay shape plays a significant role (see Figure 3.10 for such a fit to dataset B). For very soft values of the spectral index, synchrotron emission will not contribute significantly, but fits without any synchrotron contribution to the X-ray at all are not statistically favored.

The exact spectral shape and normalization in the post-acceleration synchrotron component (fitting the radio data) are very dependent on the value of $p$. When calculating a local statistical minimum, the fitting routine naturally favors the plethora of X-ray data points over the few radio data points, in finding the confidence limits for $p$, explaining non-negligible systemic residuals in the radio.

The best fit models for dataset A are shown in Figures 3.4 - 3.7 (corresponding to Table 3.3, $\text{smedge}$, $\text{reflect}$, $\text{smedge2}$ and 6kpc column respectively; see Figure 2.5 for explanation of the
components in the Figures). The fits are in general statistically good, with a large contribution to the residuals coming from the inadequacy of a simple black body to fit the IR data (see Figure 3.11). Completely removing the IR data from the best fit smeedge an improvement in the \( \Delta \chi^2_{\text{red}} \) of \( \sim 0.45 \) is obtained, to \( \chi^2_{\text{red}} = 71/116 = 0.71 \).

Three fits employing a smeedge model and one with the reflect model are shown. The smeedge and smeedge2 fits differ from each other primarily in the distribution of the energy budget, most importantly a reduced value for the magnetic dominance \( k \). Furthermore pair production will be less important in the smeedge2 fit (see Table 3.4 and Section 3.5). In contrast, the reflect fit is almost certain to suffer from too much pair production. Another possible objection to the validity of this model is the broadness of the Gaussian iron line feature. The feature is so wide that it appears to account for continuum flux, however this could not be avoided.

From the arguments outlined in the previous section it is clear that the true distance to GRS 1915+105 could be much smaller than the 11 kpc adopted for most of our fits. In an attempt to explore the effect of distance, we performed one fit on dataset A with the distance to GRS 1915+105 reduced to 6 kpc (see Figure 3.7 and Table 3.3, 6kpc column). The modifications result in an improved \( \chi^2_{\text{red}} \sim 0.15 \) and some values are closer to what we have come to expect for the HS in other black holes. Although the values for the particle distribution index \( p \) and \( f \) are not constrained, the jet luminosity \( N_j \) and the partition parameter \( k \) are closer to values found for other BHBs. At \( \sim 0.22L_{\text{Edd}} \) and \( \sim 90 \) respectively they are only a fraction of the values found for the 11 kpc fits. Also the smaller nozzle radius \( r_0 \sim 5 \) is more typical of the usual HS value. One possible interpretation of these results is that GRS 1915+105 might in fact be closer than the conservative value usually taken in the more recent literature.

For dataset A we also explored the possibility of obtaining a fit with a softer spectral index than our best-fit result (smedge), employing a relatively soft spectral index of \( p = 3 \) (\( \Gamma = 2 \)). However the flux in the radio is determined by the normalization of the post-shock synchrotron radiation and is very sensitive to the exact value of the spectral index: softening the spectral index causes a significant reduction in radio/IR model flux. To compensate for this decrease, we must tune other parameters to increase the broadband spectral flux. This can be achieved by increasing the jet luminosity \( N_j \), the electron temperature \( T_e \) or the partition ratio between magnetic and gas pressure, \( k \). The first two cause the normalization of the Comptonization component to overshoot the X-ray data, in normalization and energy respectively. These issues can in turn be absorbed by increasing the radius of the base of the jet \( r_0 \), however smaller values for \( r_0 \) appeared to be statistically favored. The remaining option is increasing the magnetic domination. In our best fit results we already find rather extreme values for this physical parameter (\( k \sim 500 \)) and increasing this value even further was therefore not pursued.

The best fit models for data set B are shown in Figures 3.8 - 3.10 (corresponding to Table 3.3, column compare, eltemp and synch respectively; see Figure 2.5 for explanation of the components in the Figures).

Fitting dataset B with similar parameter values (relating to the energy budget and assumed geometry) as dataset A yields very poor statistics (\( \chi^2_{\text{red}} \sim 40 \). Even using both the smeedge and reflect models, despite the aforementioned inconsistency, does not improve matters much. Furthermore the constant that determines the HEXTE normalization (\( A_{\text{HXT}} \)) is pushing on the preset lower boundary (\( =0.8 \), due to the steepness of the X-ray data. The fitting optimization routines attempt to lower the predicted hard X-ray emission by reducing all the data points in the 22-200 kev range. The fact that this happens indicates that the same approach as to dataset A is not able to provide a sufficient combination of components to fit the must softer dataset B. Trying to model the steep X-ray spectrum, mainly employing the exponential decay of the multicolor disk fails, because a higher disk contribution offers more (soft) seed photons. These photons are up-scattered and create a Compton tail in the hardest part of the spectrum where they result in too much flux. The compare fit is shown for comparison with dataset A, using similar values for the parameters, however it has extremely poor statistics and therefore no conclusions are drawn on the basis of this fit.

We therefore find another solution that reduces the flux of Comptonized high-energy photons by reducing the electron temperature \( T_e \). A lower temperature electron distribution would on average
Figure 3.8: Multi-wavelength (left) and X-ray band only (right) best fit result ($\chi^2_{\text{red}} \sim 39$) spectrum, for the 2005 data set, using similar values as employed for dataset A. Note the decreased resolution on the residual axis. Furthermore the residual of lowest energy X-ray point is left out so the others can be better seen in the plot.

Figure 3.9: Multi-wavelength (left) and X-ray band only (right) best fit result ($\chi^2_{\text{red}} \sim 1.0$) spectrum, for the 2005 data set, using a reduced low electron temperature $T_e \sim 4 \times 10^9$ K (cf. Table 3.3, ELTEMP column).

Figure 3.10: Multi-wavelength (left) and X-ray band only (right) best fit result ($\chi^2_{\text{red}} \sim 1.15$) spectrum, for the 2005 data set, using the synchrotron cut-off to fit the X-ray (cf. Table 3.3, SYNCH column).
not up-scatter the disk seed photons to equally high energy. We find that if we allow $T_e$ to evolve freely, we end up with a very good fit ($\chi^2_{\text{red}} \sim 1$) with a credible HEXTE normalization $A_{\text{HXT}}$, however the final temperature is rather low. At $\sim 4 \times 10^9$ K the bulk of the electrons would be sub-relativistic. The use of a relativistic Maxwell-Juttner distribution, as done in the MNW05 model, would no longer be justified, as the majority of the particles in the thermal distribution would be at $\gamma = 1$. However this approach affects a very "peaky" and steep pre-shock synchrotron component. This distribution of seed jet-base photons allows for a Comptonized photon distribution that is steep enough to fit the X-ray data up to 50 keV. Above 50 keV the Comptonized disk photons harden the model spectrum just enough to fit the entire range of HEXTE data.

An alternate possibility for fitting dataset B - with an increased electron temperature - is found when we increase the ratio of synchrotron to inverse Compton emission. With synchrotron dominating below $\sim 30$ keV (see Figure 3.10 and Table 3.3, synch column) and cutoff at about 1 keV, the subsequent exponential decay can approximate the X-ray features rather well, when combined with an accretion disk peaking at the same energy (Figure 3.10). In general we feel that, given the unknowns in the exact shape of the particle distribution around the cutoff, it is undesirable to rely on this in general during the fits. However it is worth noting that it provides a very good description of the data ($\chi^2_{\text{red}} \sim 1.15$). Fitting with the synchrotron cutoff requires a very hard synchrotron spectrum ($\Gamma \sim 1.3$) but delivers a credible HEXTE normalization constant $A_{\text{HXT}}$.

Normally such a hard spectral index would be expected only in ultra-relativistic sources (e.g. Heavens & Drury 1988). This spectrum could, however, be reconciled with the steeper $p$ found in canonical BHBs by assuming that both share such a hard injected power-law, and considering the effect of cooling on the spectrum. We do not consider that explicitly here, but are investigating such issues in a separate work. Because of the diminishing magnetic field strength along the jets, the cooling (dominated by synchrotron losses) is negligible if a maximum energy particle does not reside in a jet segment long enough to radiate most of its energy under the influence of the magnetic field. The break energy where the particle distribution steepens in "particle space" shifts according to $E \propto B^{-2} \propto r^2$ (Kardashev 1962). Hence cooling becomes less and less important with each segment, and the largest contribution to the higher frequencies from synchrotron radiation is thus from the segment where the acceleration begins. Because this zone occurs around two orders of magnitude further out in GRS 1915+105 compared to BHBs, this effect could explain why we are seeing the uncooled injected spectrum rather than the cooled spectrum, expected to be steeper by unity.

Another point to note is that for the synch fit, a strict upper boundary was placed on the partition parameter $k$. When we remove this upper boundary, an improved fit of $\chi^2_{\text{red}} \sim 0.9$ results, but for a very high value of $k = 6381$. While this slightly improves the fit, such a large $k$ is less internally consistent with the underlying model, which is still based upon hydrodynamical bulk acceleration (but see upcoming work by Polko et al., in prep.).

### 3.4.3 Stellar companion spectrum

Figure 3.11 suggests that the IR data of the companion in dataset A could be fit reasonably well by a single slope 2 power-law ($F_\nu \propto \nu^\alpha$, where $\alpha \sim 2$) of the Rayleigh-Jeans regime. However for a K2 giant with a temperature of 4455 K we clearly see a systemic discrepancy, as the blackbody slope is already flattening below the H band ($> 0.6$ eV). Hence the temperature needed to fit the IR slope exceeds that of the class K-M III giant classification done by Greiner et al. (2001b). An increase in blackbody temperature and normalization should yield a better result. While the temperature of 4455 K is the predominant value for a solitary class 2 star of this type, clearly this is not the case for the companion of GRS 1915+105.

Although the physical properties of stars under these conditions are still largely undetermined, Kaper (2001) resolved discrepancies in the effective temperatures of donor stars in high mass X-ray binaries, of 10-25% higher than expected from their spectral classification. Increasing the temperature of our blackbody by 25% to 5500 K results in an improvement in fit by $\Delta \chi^2_{\text{red}} = 0.2 - 0.3$. However, letting the temperature and normalization of the blackbody in our best-fit
The black hole binary GRS1915+105 and the plateau state

Figure 3.11: Infrared spectrum of the 1999 data set, data only (top) and best fit model (bottom). Before modeling, the infrared data is rebinned and the error set to 10%. Note that the bottom plot is logarithmic. The dashed and dotted lines in the bottom plot are the contribution of the post-shock synchrotron and the companion blackbody respectively, while the total model flux is represented by the solid line.

result SMEDGE free to vary yields an extremely high blackbody temperature of $2.14 \times 10^5$ K but improves the $\chi^2_{\text{red}}$ to 81/165 = 0.49. Such an exceedingly high temperature indicates that other IR contributions than photospheric emission from the secondary are likely. Chaty et al. (1996) already concluded the same from the rapid variations and spectral shape they observed in this band. In particular irradiation of the companion star and disk for a source as bright as GRS 1915+105 should be quite significant (e.g. Gierliński et al. 2008, 2009), but the former is not accounted for in detail in our modeling, while the latter is not accounted for at all. Other possibilities include optically thick free-free emission (although usually not observed at such high energies), or optically thin free-free radiation from an X-ray driven accretion wind (Fuchs et al. 2003; Castro-Tirado et al. 1996). Figure 3.11 shows the IR could be fit even better using a broken power-law, with a turnover from $\alpha \sim 1$ to $\alpha \sim 2$ at $6 \times 10^{-4}$ keV, suggesting e.g. a combination of a 4-5$\times 10^3$ stellar black body and outer edge of the accretion disk at $2\times 10^5$ K could also work. However for near-Eddington accretion rates, the accretion disk will peak in the X-rays and therefore its IR contribution will be negligible. Moreover, if we want the disk to subsume the role of the secondary in fitting the $\alpha \sim 2$ slope, its luminosity would have to be several times the Eddington limit.

<table>
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<th>log($n_{\text{pa}} \times r_0/c$)</th>
<th>log($n_{\text{pp}} \times r_0/c$)</th>
</tr>
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<td>16.0</td>
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</tr>
<tr>
<td>SYNCH</td>
<td>11</td>
<td>15.2</td>
<td>12.4</td>
<td>15.0</td>
</tr>
</tbody>
</table>

Table 3.4: Relative importance of pair production processes at the base of the jet, for the statistically good fits in table 3.3. $n$ is the lepton number density per cm$^{-3}$, $n_{\text{pa}}$ and $n_{\text{pp}}$ are the pair annihilation rate and production rate per cm$^{-3}$s$^{-1}$, respectively. In order to compare the production rates to the lepton density, they have been multiplied by an estimate for the residence time, given by $r_0/c$. If $n \gg \text{Max}(n_{\text{pa}}, n_{\text{pp}})$, and $n_{\text{pa}} \gg n_{\text{pp}}$ we should be safe from pair production. From these two considerations, the former is the most important.
3.5 Discussion: GRS 1915+105 and the canonical states

Clearly, and perhaps surprisingly, there are large differences between datasets A and B, both in the $\chi$ state of GRS 1915+105. While dataset A can be fit fairly well by the same model as for BHBs, dataset B is marginal, thus they do not seem to represent a standard class of spectral behavior. Interestingly, Rau & Greiner (2003) also find significant variety between the individual $\chi$ states they observed in over 4 years. While the classification done by Belloni et al. (2000) is based on lightcurves and color-color diagrams, the sub-indices 1-4 do not refer to a phenomenological classification scheme, but only denote their temporal sequence. Hence there remains the possibility that the $\chi$ state still encompasses multiple physical regimes that we may be probing here via spectral fitting.

Alternately, dataset B could be in some kind of transitional state, despite being classified as a typical $\chi$ state. In GRS 1915+105 complete state changes can happen on very short timescales in comparison to the canonical BHs (seconds). Such rapid changes generally occur only between states A and B, however. While state C episodes can also last for short periods of $<100$ seconds, the current datasets involve only the longer (a day or more) intervals. It is clearly possible that during dataset B observations GRS 1915+105 was already in a transitional state and therefore the model, which assumes a steady outflow, is no longer applicable. Either way, both of the statistically convincing dataset B fits (SYNCH and ELTEMP) have other problems in their physical interpretation, as discussed above, and it is possible that the model simply cannot apply to observations at such high luminosities (see Figure 3.2).

For all of the fits presented, two parameters consistently stand out by nature of their larger values when compared to applications of the MNW05 model to BHBs: the location of initial particle acceleration in the jets, $z_{acc}$, and the energy partition parameter $k$. As mentioned before, we typically find $z_{acc}$ to fall 1-2 orders of magnitude closer in towards the base of the jets in fits of other BHBs in the HS. Similarly, while the canonical BHBs often display mild magnetic domination over the gas pressure, with $k \approx 10$, we find at least an order of magnitude increase is needed to fit GRS 1915+105 plateau states. Finally, the overall powers required (as indicated by the jet normalization parameter $N_j$, see MNW05 for an explanation) are strikingly larger than the maximum observed in other BHBs. Canonical BHBs displaying these luminosity levels in $L_{Edd}$ would long have shut down their jets, yet somehow GRS 1915+105 seems to be operating on a different energy scale.

These three parameters may well be related. Magnetic fields are by now known to play a major role in accelerating and collimating relativistic plasma outflows, and current simulations of jet formation favor rather high values of $k$ (usually expressed as low values of the plasma $\beta$; McKinney 2006; Komissarov et al. 2007; McKinney & Blandford 2009). It may be that $k$ is positively correlated with the overall power input into the jets. Evidence for such a scaling has already been observed in individual fits to broadband data the LLAGN1 M81* over the course of a yearlong campaign (see, e.g. Markoff et al. 2008, Figure 22). The overall magnetic field strength and configuration will certainly influence the formation of particle acceleration structures, however the dependence on total power and other parameters has not yet been thoroughly explored (but see Polko et al., in prep.).

A higher level of magnetic domination can also account for why the fits favor an electron temperature $T_e$ that is a factor of two or more lower than the the bottom of the range ($\sim 2 \times 10^{10}$ K) found for the canonical microquasars. A stronger magnetic field relative to the radiating plasma at the base of the jets in our model is synonymous with the same being true at the inner edge of the accretion flow, as we assume they are directly related. A relatively stronger magnetic field would result in relatively more lepton cooling in the inner regions, and consequently a lower equilibrium temperature. Such lower temperature conditions are also consistent with an ADAF (see Section 1.4), where the accretion rate is high enough for the resulting opacity to trap the energy liberated in the gas. This energy consequently stays shielded from observation as it is dragged into the central black hole. ADAFs cool via the leptons, so although the ion temperature may be close to virial (or even above, depending on the exact jet-launching processes), the allowed electron temperatures are of order $10^{9}-10^{10}$ K, fully consistent with our obtained values. Interestingly, Meier
(2005) suggests that a magnetically dominated accretion flow (MDAF; also see Section 1.4), could exist interior to the ADAF at radii \( \lesssim 100 r_g \) and would even be required for the production of jets.

In contrast, the geometry of both the base of the jets and the cool accretion disk parameters preferred in the fits is entirely consistent with the range found in the canonical BHBs. The radius of the base of the jets is 3–20 \( r_g \) exactly as seen in fits of both BHBs as well as LLAGN. Similarly, we obtain best fit values for the inner cool accretion disk radius \( r_{in} \sim 2.5 - 4.5r_g \), which is consistent with such a high accretion rate, as well as a spinning black hole as suggested by McClintock et al. (2006). What is not consistent, and differs compared to other BHB fits, is that the jet base radius in almost all fits is larger (sometimes only marginally) than \( r_{in} \). Since in our model the jets should be launched by the inner, hot accretion flow, this is not internally self-consistent and by such criteria, only fits \textsc{smedge2} and \textsc{eltemp} would survive. As the MNW05 model does not yet self-consistently solve the jet launching physics from a corona, which is still an unsolved problem in the field, we have not made this a hard constraint.

Although in most BHBs such a small \( r_{in} \) would be indicative of a soft state, this is not true for GRS 1915+105. Due to the high accretion rate, a geometry where the accretion disk reaches the innermost stable orbit is almost expected. The fact that the 6KPC fit has an inner radius of less than 1 \( r_g \) is most certainly due to the oversimplifications mentioned above and similar inner radii were also obtained for the canonical BHs (see e.g. MNW05). In fact, such small radii appear to be a common occurrence, also using other models: analyzing 4 years of \( \chi \) state observations Rau & Greiner (2003) obtained inner radii of far less then 1 \( r_g \), from the normalization of the disk contribution using the \textsc{diskbb} model. Although this model is known to underestimate the inner disk radius by a factor of 1.7-3, due to Doppler blurring and gravitational redshift, the obtained radii are still unrealistically low. The disk temperatures of 1–4 keV they find are higher than ours (0.7-0.9 keV), but again the \textsc{diskbb} model is known to produce factor 1.7 overestimates (Lee et al. 2002). Comparing the two plateau states to each other, the required disk contribution is much higher in dataset B. For dataset A, however, disk normalizations are high but comparable to those for the canonical BHBs, suggesting that, as is true for the HS, the classification of individual plateau states should not be based on the luminosity.

In agreement with the findings of many other authors (e.g. Sobolewska & Życki 2003; Rau & Greiner 2003) who have studied the \( \chi \) state, we find that the inverse Compton component is always dominant over the direct accretion disk contribution in the X-ray regime for our datasets. Although in our case the Comptonization is mainly due to SSC rather than purely from thermal accretion disk seed photons. Only in one dataset B fit (\textsc{synch}) are the X-rays not dominated by the Comptonization, but the validity of this fit may be questionable on other grounds as discussed above.

As brought up as a potential issue by Maitra et al. (2009) and Malzac et al. (2009), for high luminosities conditions at the base of the jets, which is quite compact, may comprise high enough photon densities that pair production can become important. As these processes are not yet implemented in the MNW05 model, we check the pair production and annihilation rates (\( \dot{n}_{pp} \) and \( \dot{n}_{pa} \) respectively) at the base of the jets using the methods described in Maitra et al. (2009), and references therein. The results are shown Table 3.4. If the lepton number density \( n \) is not larger than \( \dot{n}_{pp} r_0/c \) and \( \dot{n}_{pa} r_0/c \), where \( r_0/c \) is roughly the \textit{residence time} for the leptons in the jet base region, we need \( \dot{n}_{pp} < \dot{n}_{pa} \) to be able to neglect pair production.

Although our \( T_e \sim 10^{10} \) is a factor of 5 lower than the limit deemed problematic for GX339-4, the extreme luminosity of GRS 1915+105 means that pair production could be an issue for some of our fits. The source of potential pair production is a high-energy tail above \( \sim 0.5 \) MeV, due to the Comptonization of thermal accretion disk photons. The amount of flux in this tail is directly dependent on the normalization and temperature of the accretion disk component, which itself is mainly constrained by the value of \( N_H \) and the predicted flux in the tail region. Higher values of \( N_H \) result in higher accretion disk fluxes to compensate for the increased absorption in the soft X-ray band.

While we fixed the value for \( N_H \) in our fits for consistency with the IR reduction (see below), an increasing number of works are concluding that the column density is actually variable, ranging in values from \( N_H \sim 216 \times 10^{22} \text{ cm}^{-2} \) and potentially linked with an intrinsic warm absorber at
the highest values Belloni et al. (2000); Klein-Wolt et al. (2002); Lee et al. (2002); Yadav (2006). We find that the highest values of the $N_H$ range are not consistent with our model, as it causes too great a decrease below 3 keV, resulting in too much flux in the high-energy tail. However we cannot leave $N_H$ free to vary in our fits because of the need for self-consistency in the multi-wavelength analysis: we would need to re-reduce the IR data for each new value of $N_H$. We chose a moderate value of $4.7 \times 10^{22} \, \text{cm}^{-2}$ in order to be consistent with prior fits of the various $\chi$ substates (e.g. $2 \times 10^{22} \, \text{cm}^{-2}$; Belloni et al. 1997a) and the assumptions of Chatty et al. (1996). For dataset B this value is likely too high. Letting $N_H$ free to vary (without correcting the IR dereddening accordingly) settles on the lower bound allowed of $2 \times 10^{22} \, \text{cm}^{-2}$, with an improvement of $\chi^2_{\text{red}}$ to $\sim 26$, mainly due to the resulting reduction in the Comptonized high-energy tail. Therefore our conclusions is that our results are highly dependent on the value of $N_H$, and that this value is likely different for the two datasets, although for the reasons described above we have chosen to use a single value. Similarly, the uncertainty in the exact value means for either means that too much pair production can be avoided in particular by lower values of $N_H$, that seem to be statistically favored if we allow $N_H$ to vary.

The variation in absorption may be due to a variation in a disk wind. Such a wind would be expected from the fact that GRS1915+105 accretes near or above its Eddington limit (Lee et al. 2002). The wind density should be a strong function of the disk luminosity (Yadav 2006). Recently Neilson & Lee (2009) wrote an interesting article in which they devise a simple jet-quenching mechanism, driven by a disk wind. They recognize a P Cygni profile in the spectrum (although evidence for such a profile was not yet found in Lee et al. 2002) and thus infer the existence of a such a wind. They also find evidence that the (commensurate) mass loss alternates between the wind and the jet. At sufficiently high luminosities, the jet would be suppressed, as the intense disk radiation field redirects the accretion flow into a disk wind. This mechanism would allow GRS1915+105 to maintain an outflow/inflow equilibrium, independent of the spectral state (and regulate its accretion rate through feedback with the environment). The exact coupling between the disk and the jet would be mediated by Comptonization and photoionization processes.

Interestingly, Titarchuk et al. (2009) recently found that the red-skewedness of the iron Kα line complex observed in RXTE and XMM spectra from BH (viz. GX339) and NS and CV could develop in an illuminated disk wind. Normally the assumption is that the 6.4 keV iron line emerges when the inner part of the disk is illuminated by hard X-rays released through the accretion processes or the jet. The red-skewedness is then explained from GR effects (e.g. stemming from a maximally spinning BH; Laor 1991). However Titarchuk et al. (2009) find that the reduction in energy of the line photons could just as well be explained by electron scattering in a disk wind shell.

In the other works (see section 2) the MNW05 model is often convolved with the reflect model, to account for accretion disk contributions. Although it is also possible to fit the dataset A spectrum using the reflect model, the results for the smedge model are better. This can perhaps be understood from the fact that the reflect model causes a spectral hardening, while X-ray emission in the plateau state is generally softer than that in the canonical HS (Fender & Belloni 2004). At $\Omega/2\pi \sim 0.31$, the obtained reflection component normalization in our reflect fit is larger than the $\lesssim 20\%$ MNW05 thought possible in order for Compton processes at the base of the jet to take place. The value is however consistent with those obtained for GX339-4 in Maitra et al. (2009). Fitting state B and C observations, Feroci et al. (1999) find even higher values for the reflection fraction in GRS1915+105, of $\Omega/2\pi \sim 0.5$ (Fender & Belloni 2004). When fitting dataset B with a reflection component in stead of a smedge model the normalization becomes even more extreme and the radiation hitting the reflector would be of the same order as that directly viewed by a distant observer. This result is similar to that obtained by Rau & Greiner (2003), who also found extremely high values for the reflection parameter that were inconsistent with other measurements. However such large values are thought to be inconsistent with mildly relativistic beamed coronae or jet models (Migliari et al. 2007; section 3.3.4).

For this work we found that mimicking the reflection component using a relatively broad iron line in conjunction with a relativistically smeared edge was statistically favored. The smedge model is however only a phenomenological model and the parameters can not be regarded as
physical quantities. This approach to fitting GR 1915+105 has been used successfully before by Done et al. (2004), but they fix the width of the Gaussian and the smearing to 0.5 keV and 7 keV respectively and make no statement about the obtained normalizations $\tau_{\text{max}}$. Gierlinski et al. (2001) also fix the width of the smearing to 0.5 keV fitting GRO J1655-40 and obtain normalizations that match ours, but they are exploring a soft state. Reis et al. (2008) fit GX339-4 hard state RXTE data, but in contrary to our approach they do not tie the widths of the Gaussian and the edge but leave them completely free to vary, yielding edge energies and normalizations compared to ours, but a somewhat higher edge width. The width of the Gaussian in this latter work is however comparable to ours. In summary, if we interpret $\tau_{\text{max}}$ as the (phenomenological) optical depth, we find values for GR 1915+105 comparable to the canonical black holes. The same is true for the edge energy. It is however impossible to meaningful compare the widths of the iron line and the edge, due to distinctively different approaches used in the above works.

### 3.6 Conclusions

Despite the fact that the MNW05 model was originally intended for application to hard states in canonical BHBs only, it appears to well approximate the plateau state in GR 1915+105. However it does not produce convincing results in every instance. While some of the parameter values obtained are quite extreme, the results for the dataset A appear credible and consistent with what we have found in canonical BHBs. Dataset B however presents more difficulties, and a more solid determination of the distance and absorption column would go far to help us understand the difference between these two plateau states.

Clearly this work confirms that GR 1915+105 is in a much more extreme range of parameter space for an outflow-dominated model, requiring near- or super-Eddington accretion rates, maximal jet powers and high levels of magnetic domination. Although our results confirm previously noted plateau state issues, such as the need for a variable $N_H$, our model for the first time incorporates the entire broadband and has allowed the comparison between the jet producing plateau and hard states. While the baseline geometry seems similar, the plateau states of GR 1915+105 are not low-luminosity as with HS BHBs, and settle on a range where the acceleration of particles occurs much further out in the jets, which can be two orders of magnitude further out of equipartition in the direction of magnetic domination. While these two effects may be linked, the model explored here cannot self-consistently address this, but in another work we are exploring the links between physical parameters and the location of particle acceleration fronts (Polko et al., in prep.). A slightly lower electron temperature is also found compared to other BHBs, which can be interpreted in the context of the higher cooling rates found at GR 1915+105’s extreme luminosity.

The main consequence of these differences is that the synchrotron component from the outer jets no longer dominates the soft X-ray band, although a non-negligible X-ray synchrotron flux of $\sim 10$ % the inverse Compton flux below $\sim 50$ keV seems required for the statistically favored fits. In comparison, the MNW05 model applied to canonical BHBs favors synchrotron emission dominating the flux at least up to 10 keV. Interestingly our results are thus qualitatively similar to the results found from the blazar sequence (e.g. Ghisellini et al. 1998), where higher powers correspond to a decrease in the frequency range where synchrotron power peaks. It is clear that time-dependent effects such as cooling breaks (see Chapter 5) demand further exploration.

The remarkable differences between two individual plateau state observations does raise questions about whether the $\chi$ substates are distinct enough to classify all plateau characteristics. While both datasets explored here bear all the characteristics of the plateau state, the fits with a single model show more variations in free parameters than found even between different sources in the HS of canonical BHBs. Thus the current classification scheme based solely on X-ray colors and timing properties may need to be expanded based on broadband attributes.

Our results support a conclusion that, although expressing quite different properties than the HS in canonical BHBs, GR 1915+105 plateau states can still be described by the same broadband model with a steady outflow tied in power to the accretion inflow. However the BHB
model is clearly forced into very extreme ranges, which themselves provide some new clues about the relationship between accretion rate, jet production and particle acceleration.

Although GRS 1915+105 is one of the most extensively studied BHBs over the last decades since its discovery, we are far from understanding this source. At some point in the future (somewhere between 2012 and 2092, assuming the source went into outburst when it was first discovered; Deegan et al. 2009) GRS 1915+105 will invariably retreat into quiescence, and should finally yield better insight into the connection between its current accretion properties and those at sub-Eddington rates.
The GX 339–4 hard state and HID hysteresis

4.1 Introduction

4.1.1 Hysteresis

As shown in Section 1.4, the canonical microquasars generally trace out a q-shaped pattern in the HID (see Figure 1.8). Looking at the HID for GX 339–4 specifically (Figures 4.3 and 4.6), we see that – depending on whether the X-ray flux is increasing or decreasing – this BHB also traces out two parallel horizontal tracks, showing that after this source emerges from quiescence, the transition from hard-to-soft state occurs at a luminosity a factor 10–100 higher than that of the soft-to-hard state transition. Clearly the history of the source influences its evolution, effecting this track in the HID. There is another way to interpret this feature, when we consider that in most canonical BHs we expect a source to display lower luminosity levels in the hard state, than in its respective soft state (although this is not always true, see Section 1.4). We could thus say that in GX 339–4 we observe a lag between luminosity and spectral state\(^1\). For these reasons the phenomenon described here is referred to as hysteresis (Zdziarski & Gierliński 2004).

The goal of this Chapter is to see if we can find a single physical mechanism driving the hysteresis using the MNW05 paradigm.

4.1.2 The black hole binary GX 339–4

GX 339–4 is one of the most extensively studied XRBs. Its behavior is much more “canonical” than that of GRS 1915+105 although it sometimes skips a quiescence phase at the end of an outburst, and it was one of the first BHBs for which a complete set of state transitions was observed. The source was discovered in 1972 by OSO 7 (Markert et al. 1973) and lies near the Galactic plane, \(4^\circ.3\) south of the Galactic equator. It was observed mostly in the HS (although some transitions

\(^1\)Several other LMXBs also exhibit similar behavior, e.g. GRS 1758-258, 1E 1740.7-2942 (Smith et al. 2002; del Santo et al. 2005), XTE 1650-500 (Rossi et al. 2004), XTE J1550-564 (Kubota & Done 2004), XTE J1859+226, XTE J2012+381 (Maccarone & Coppi 2003), H1743-322 (Remillard & McClintock 2006), and putatively GRS 1915+105 (Klein-Wolt et al. 2002; Meyer-Hofmeister et al. 2005). In contrast, the "poster-child" for HMXBs, Cyg X-1, shows a one-to-one correspondence between spectral state and luminosity (because the latter stays approximately the same).
The GX 339–4 hard state and HID hysteresis

Figure 4.2: The GX 339–4 X-ray lightcurve (top panel) and $S_{6-10\,\text{keV}}/S_{3-6\,\text{keV}}$ hardness ratio (bottom panel) from the second Proportional Counter Unit aboard the RXTE satellite. The plot starts January 1st 2004 (MJD 53005). Clearly the relative errors are quite large at first. Hence we have not included these data points in the HID, Figure 4.3. We can see GX 339–4 going into outburst and peaking in luminosity almost exactly a year later, before starting the decay phase. The dates of the rise and decay datasets are indicated by the black arrows.

to the IMS were observed; Miyamoto et al. 1991b) until a protracted spell of quiescence started in 1999 (see e.g. Kong et al. 2000). Since 2002 the source has gone into outburst almost every other year, again spending most of its time in the hard state. Therefore this system has enormous potential for furthering our understanding of low-rate accretion. Not surprisingly it has thus been studied in all wavelength bands, from radio to $\gamma$-rays (see e.g. Markoff et al. 2003a and references therein).

Most Galactic BHBs are transients that allow for the detection of the companion star during quiescent phases. The donor in GX 339–4 however eludes direct spectroscopic observations\(^2\) even during its faintest episodes (Shahbaz et al. 2001). Because of the lack of spectral features of the secondary, GX 339–4 data can be regarded as relatively uncontaminated, as the observed flux can be assumed to be largely due to the radiative processes involved in causing the various canonical BH spectral states. Hence this system is considered an ideal source for studying those states.

GX 339–4 was the first source found to exhibit a positive correlation between the radio and the X-ray flux in the HS (Hannikainen et al. 1998). The correlation confirmed the classification of GX 339–4 as a microquasar, as it is thought to stem from coupling between accretion and ejection,\(^2\)

\(^2\)The upper limit on the optical luminosity derived from the faint episodes does however allow for classification of the system as an LMXB and mass transfer is thus expected to occur through Roche-lobe overflow.
i.e. the disc and the jet / the inflow and the outflow (see e.g. Corbel et al. 2000). The radio/X-ray correlations were quantified by various authors (e.g. Corbel et al. 2003; Markoff et al. 2003a), using a previous version of the MNW05 paradigm, suggesting the X-ray flux is due to optically thin synchrotron radiation from a compact jet\(^3\). The results were generalized by Heinz & Sunyaev

\(^3\)Gallo et al. (2004) first confirm a compact jet directly from ATCA/radio observations, resolving a large-scale steady jet (see Figure 4.1), that grew with an apparent velocity greater than 0.9c, to an extent of 12 arcsec.
The GX 339–4 hard state and HID hysteresis

Table 4.1: Observations representative of the rise and a decay phase of a GX 339–4 hard state. The RXTE data are ObsID 90118-01-07-00 and 91095-08-07 respectively. The optical (V and J bands) and infrared (J and H bands) were done with the SMARTS/ANDICAM at 1.3m CTIO telescope. The ATCA measures at 1.34, 2.37, 4.8, 6.21 and 8.64 GHz (see Table 4.2 for details).

<table>
<thead>
<tr>
<th>Band</th>
<th>Instrument</th>
<th>MJD range</th>
</tr>
</thead>
<tbody>
<tr>
<td>X-ray</td>
<td>RXTE</td>
<td>53082±0.06</td>
</tr>
<tr>
<td>IR/optical</td>
<td>CTIO</td>
<td>53082.31±0.05</td>
</tr>
<tr>
<td>Radio</td>
<td>ATCA</td>
<td>53082.79±0.29</td>
</tr>
<tr>
<td></td>
<td></td>
<td>53489±0.06</td>
</tr>
<tr>
<td></td>
<td></td>
<td>53489.33±0.07</td>
</tr>
<tr>
<td></td>
<td></td>
<td>53489.65±0.04</td>
</tr>
</tbody>
</table>

(2003) to include X-rays from radiatively inefficient processes. The correlations ultimately led to the discovery of the Fundamental Plane of black hole accretion (see Section 1.3).

Additional correlations have also been revealed, between the X-ray and various other bands: Corbel & Fender (2002) already found evidence for jet emission in the (near-)infrared band and correlations between this band and the X-rays were later confirmed by e.g. Homan et al. (2005) and Coriat et al. (2009). The latter also found correlations between the optical and the X-ray, which they suggest can be mainly attributed to thermal emission from the accretion disk.

4.2 Observations and data reduction

At the beginning of 2004 GX 339–4 went into outburst (see Figure 4.2, top panel). Over a few months the XRB evolved towards higher X-ray intensity (along the rightmost vertical of the HID, Figure 4.3), to ~ 200 counts s\(^{-1}\). We will refer to this episode as the rise phase of the outburst, during which the source was in the HS for an extended period (see Figure 4.2, bottom panel). Subsequently the hardness started to decrease steadily and the intensity diminished with a factor ~ 2 for some time, only to increase even more later on. At this point the HID goes through a loop (top left corner of the HID), where the hardness temporarily increases again and the outburst climaxes at a top X-ray intensity of ~ 700 counts s\(^{-1}\) about a year after it started, before starting its decline in 2005 (after MJD ~53475; Figure 4.2). In the next phase the BHB approaches its softest point at a luminosity of a factor 7 less. After crossing the putative jet line (see Figure 1.8) again, the source reenters the HS and the outburst starts its decay. This dataset representing this episode has almost the same hardness ratio as the rise dataset (cf. Figure 4.2). The outburst ends when the BH has descended into quiescence (see Section 1.4) again.

In this Chapter we analyze two quasi-simultaneous datasets, representing a rise phase HS (March 18th 2004) and a decay phase HS (April 29th 2005). For both occasions the HS is confirmed by the PDS and position in the HID (Figure 4.3). For the rise observation the HS assessment is further investigated in Section 4.5.1, for reasons explained there.
Observations and data reduction

4.2.1 X-ray: RXTE data reduction

We use data from two instruments on board the RXTE: PCA (Jahoda et al. 2006) and HEXTE (Rothschild et al. 1998). The PCA originally comprised an array of five identical gas-filled Proportional Counter Units (PCUs), that each have a collecting area of 1200 cm$^2$ and are sensitive to X-rays between 2.5 and 60 keV. The PCUs are periodically switched/disabled to increase their lifetime (Jahoda et al. 2006). From all five, PCU2 is best calibrated and hence was used for the observations in this section (Coriat et al. 2009). The PCA spectra were extracted using standard 2 mode data. Due to uncertainties in the first four PCA channels, we only include PCA data above 3 keV. Above energies $\geq$ 20 keV HEXTE provides reliable data and we ignore the PCA data above 22 keV. At high energies, the HEXTE spectrum was ignored above 200 keV. The HEXTE spectra were extracted from standard mode Archive and were dead-time corrected. The X-ray data reduction was performed using HEASOFT 6.4, applying standard extraction criteria for the pointing offset from the nominal source position, source elevation, exclusion time for the South Atlantic Anomaly (SAA), and the “electron ratio”.

Table 4.2: Radio (ATCA) and IR/Optical (CTIO) observations for the rise and a decay phase.

<table>
<thead>
<tr>
<th>Band</th>
<th>Frequency (GHz)/Band</th>
<th>Rise Flux (mJy)</th>
<th>Decay Flux (mJy)</th>
</tr>
</thead>
<tbody>
<tr>
<td>Radio</td>
<td>1.34</td>
<td>2.70±0.25</td>
<td></td>
</tr>
<tr>
<td></td>
<td>2.37</td>
<td>3.76±0.11</td>
<td></td>
</tr>
<tr>
<td></td>
<td>4.80</td>
<td>5.15±0.05</td>
<td>2.88±0.08</td>
</tr>
<tr>
<td></td>
<td>6.21</td>
<td>5.39±0.05</td>
<td>3.32±0.10</td>
</tr>
<tr>
<td></td>
<td>8.64</td>
<td>4.98±0.11</td>
<td></td>
</tr>
<tr>
<td>IR</td>
<td>$H$</td>
<td>24.10±0.83</td>
<td>9.83±0.33</td>
</tr>
<tr>
<td></td>
<td>$J$</td>
<td>29.17±1.56</td>
<td>10.81±0.57</td>
</tr>
<tr>
<td>Optical</td>
<td>$V$</td>
<td>24.37±2.71</td>
<td>10.75±1.19</td>
</tr>
<tr>
<td></td>
<td>$I$</td>
<td>37.72±6.96</td>
<td>15.10±2.71</td>
</tr>
</tbody>
</table>
4.2.2 Optical and infrared

The optical and infrared (IR) measurements were obtained using the SMARTS (Small and Medium Aperture Research Telescope System) / ANDICAM camera on the 1.3m CTIO telescope. The optical $V$ and $I$ filters are Johnson-Kron-Cousins, while the IR $J$ and $H$ filters are standards CIT/CTIO. The observed magnitudes were converted into flux density assuming an optical extinction $A_V = 3.7 \pm 0.2$ and using the extinction law by Cardelli et al. (1989). The error on the observed quantities is dominated by the uncertainty in the optical extinction.

4.2.3 Radio

For both datasets, the radio observations were done employing the Australia Telescope Compact Array (ATCA). The east-west oriented array consists of six 22-m antennae with a maximum baseline of 6 km. Most measurements are done at 4.80 and 8.64 GHz, but sometimes additional bands around 1.34, 2.37 and 6.21 GHz are available (see Table 4.1 and 4.2, decay data and rise data respectively).

All the X-ray, IR and radio data mentioned in Table 4.1 and 4.2, have been made available to us by S. Corbel. The details on the radio analysis will be made available in a forthcoming paper (Corbel et al. 2010, in prep.).

4.3 Input physical parameters

In this Section we will explain our choices for the physical parameters in our model that we were able to take from literature. The values are listed in Table 4.3.

As mentioned in Section 4.1.2 the companion in GX 339–4 has eluded direct spectroscopic observation. The relative obscurity of the companion makes it hard to determine the system parameters. So in spite of the fact that GX 339–4 was one of the first proposed BHBs (Samimi et al. 1979) even a basic dynamical parameter, such as the orbital period, remains under dispute. Constraining the orbital period is inevitable when estimating the mass function (Equation 3.1).

Several early measurements of GX 339–4 in quiescence suggested an orbital period of less then a day. Based on modulations in optical photometry, Callanan et al. (1992) first suggested a period of 14.8 h. Using simultaneous photometric and spectroscopic data, Cowley et al. (2002) later found evidence for a similar period of $\sim$0.7 days. Hynes et al. (2003) however use a relatively new approach to determine the orbital parameters. This indirect spectroscopic method, involving the appearance of sharp $N\text{iii}$ emission, was also used to determine the binary parameters for Her X-1 and Sco X-1 (Still et al. 1997 and Steeghs & Casares 2002 respectively). The $N\text{iii}$ lines are thought to develop when the companion is irradiated and thus allow for a detection of the radial velocity curve. This curve is compared to the velocities obtained from He II wing measurements and an orbital period $P_{\text{orb}}$ of 1.7557$\pm$0.0004 days is consistent with both results. Already having obtained the secondary semi-amplitude $K_2 = 317 \pm 10$ km s$^{-1}$ from the same curves, a corresponding mass function can be derived. Using the best-fit period, $f(M) = 5.8 \pm 0.5$ M$_\odot$ (seemingly confirming the fact that GX 339–4 is a dynamical BHB, as this exceeds the maximum neutron star mass). However there is still quite some uncertainty in this value, as using the shortest acceptable alias for $P_{\text{orb}}$ gives a 95% confidence lower limit of 2.0 M$_\odot$. In addition, Hynes et al. (2003) find it likely that the value for $K_2$ has been underestimated, as the N III and C III lines emerge from the inner hemisphere of the secondary (McCintock et al. 1975). Muñoz-Darias et al. (2008) compute the so-called $K$-correction for GX 339–4: modeling the deviation between the reprocessed light center and the center of mass of a Roche-lobe filling star. If they assume a lower limit for the companion consistent with the "stripped giant" model ($\sim$0.17 M$_\odot$), the K-correction yields a BH mass of at least 6 M$_\odot$. Letting the same model agree with typical X-ray flux values and implied

---

4 The orbital period derived with this new method is however clearly inconsistent with the shorter orbital periods from Callanan et al. 1992 and Cowley et al. 2002
Table 4.3: Fixed physical parameters used in the fitting process, obtained from the literature. We omit the error bars, as they are not used in the fits.

<table>
<thead>
<tr>
<th>parameter</th>
<th>value</th>
<th>units</th>
<th>Reference</th>
</tr>
</thead>
<tbody>
<tr>
<td>column density</td>
<td>6</td>
<td>10^{21} cm^{-2}</td>
<td>a</td>
</tr>
<tr>
<td>mass</td>
<td>7</td>
<td>M_⊙</td>
<td>b</td>
</tr>
<tr>
<td>inclination</td>
<td>47</td>
<td>°</td>
<td>c</td>
</tr>
<tr>
<td>distance</td>
<td>6</td>
<td>kpc</td>
<td>d</td>
</tr>
</tbody>
</table>


mass transfer rate for this system suggests \( M_2 \gtrsim 0.3 \, M_⊙ \) and \( M_{BH} \gtrsim 7 \, M_⊙ \). We will adopt this lower limit as a fiducial value for the mass.

For GRS 1915+105 an estimate for the inclination could be made because super-luminal flares were observed, allowing for a determination using the difference in proper motions from the approaching and receding blobs (see Equation 3.2 and Section 3.3.2). Although (in addition to a collimated jet) similar flares have been observed in GX 339–4 (Gallo et al. 2004), this data is of insufficient quality to do the same. Hence there is still a lot of discussion about the inclination and only upper and lower limits are suggested in the literature. As no optical or X-ray eclipses are seen in this system, Cowley et al. (2002) argue for an orbital inclination with the line of sight of less than 60°. On the other hand, the moderate modulation in the V-band when the flux in this band is high and the relatively large secondary mass function are incompatible with a low value for \( i \) (Zdziarski et al. 2004). Therefore it seems reasonable to assume \( i \gtrsim 45° \). Using MNW05, Maitra et al. (2009) obtained inclinations in the range of 45-50 degrees, but the value could not be properly constrained. They chose to fix the value to 47° and we will do the same, for consistency.

As may be suspected, the distance to GX 339–4 has also been a subject of ample debate. Early estimates were exceedingly low, ranging from 1-4 kpc. Using the upper limit on the companion star’s optical luminosity Shahbaz et al. (2001) put a lower limit on the distance of 5.6 kpc. Zdziarski et al. (2004) derive an inclination dependent mass range of \( 6.7 \lesssim d_{\text{min}}(i) \lesssim 9.4 \) kpc, using the orbital parameters and a mass ratio \( q \lesssim 0.08 \) from Hynes et al. (2003), together with the fact that GX 339–4 is a Roche-lobe filling LMXB, so that the Eggleton approximation Eggleton (1983)

\[
\frac{r_L}{a} = \frac{0.49q^{2/3}}{0.6q^{2/3} + \ln(1 + q^{1/3})},
\]

(4.1)

where \( r_L \) is the size of the Roche lobe, \( a \) is the separation, applies. Maccarone (2003) finds a lower limit on the distance, in agreement with these values, of \( \gtrsim 7.1 \) kpc, using the empirical evidence that in most canonical BHBs the transition between soft state and hard states seems to occur at a fixed fraction of their Eddington luminosity (\( L \sim 0.02 \, L_{\text{Edd}} \)). On the basis of a kinematic study, Hynes et al. (2004) argue for an even higher distance of 15 kpc, after observing a +30 km s^{-1} redshifted component. Together with the slightly positive systemic velocity measured for the system this would suggest that GX 339–4 is on the far side of the Galaxy. It is stated however that this is rather uncertain and they deduce a lower limit of \( \sim 6 \) kpc interpreting the same measurements differently. If the dynamical binary parameters and the companion star class had been known, it would have been possible to estimate the radius and the temperature (viz. surface brightness), and, comparing with the apparent brightness, deduce the distance with increased accuracy (see e.g. Gelino et al. 2001). However the true distance to GX 339–4 clearly remains an open question for now. We will therefore adopt the rather conservative lower limit of 6 kpc from Hynes et al. (2004) as the distance in our models, to remain consistent with other works using the MNW05 model (Markoff et al. (2005) and Maitra et al. (2009)).

Lastly we will fix the value for the hydrogen column density to \( 6 \times 10^{21} \) cm^{2}, in agreement with limits given by Hynes et al. (2003) and, again, consistent with Maitra et al. (2009).
4.4 Model and spectral fitting

For spectral modeling we use ISIS (Houck & Denicola 2000), compiled with XSPEC version 12.3.1x libraries (Arnaud 1996). The model is forward folded through the detector response matrix. To account for the additional uncertainties in the PCA response matrix, a 0.6% systematic error has been added in quadrature to all PCA data. As the relative calibration of the PCA and HEXTE instruments is not certain, the normalization factor for the PCA data is set to unity and tied to the radio and IR normalizations during the fits, while the HEXTE data normalization is left free to vary.

All fits for these datasets are done using the following components: (1) The MNW05 steady-state outflow-dominated model that includes a multi-color blackbody accretion disk and a single blackbody for the companion star; (2) an additive Gaussian line profile, with a line energy left free to vary between 6 and 7 keV, and line width $\sigma$ free to vary between 0 and 2 keV; Models (1)+(2) are convolved with Compton reflection from a neutral medium ($\text{reflect}$; Magdziarz & Zdziarski 1995), as expected when if an accretion disk is present, and multiplied with a photoelectric absorption model ($\text{phabs}$) to account for the interstellar medium. For the $\text{reflect}$ model we assume the viewing angle to correspond to the jet inclination.

Consistent with Maitra et al. (2009) we will only consider fits with electron temperatures $T_e$ below $\lesssim 5 \times 10^{10}$ as to avoid pair production processes (see also Section 3.5). In addition, we will assume GX 339–4 to be rotating, placing the ISCO at a minimum of 1 $r_g$ (see Section 1.3), suggesting the inner accretion disk radius should be at least $r_{\text{in}} > 1r_g$. In addition we fix the outer disk radius ($r_{\text{out}} = 10^7 r_g$), as we do not know the exact size of the disk and we want it to be large enough to avoid fitting the IR data with disk features we are not sure about (i.e. the Rayleigh-Jeans regime of the outer annulus). The disk is however not only a weak component in comparison to the complete model: it can be considered a second order effect, as we are exploring the jet physics as prime instigator of the hysteresis and are ignoring the detailed disk-physics. Moreover the disk is not constrained by the X-ray data in most fits, but seems only required to get a correct normalization in the IR. In line with these arguments it is of no importance that the disk is not present in all the rise fits (compare Figure 4.5, bottom two fits and top two fits respectively).

In contrary to GRS 1915+105 no evidence for a significant flux contribution from the stellar companion has been observed in the GX 339–4 hard state, hence we will put the normalization of the single blackbody to zero.

While for GRS 1915+105 we wanted to be sure the jet was long enough to account for the slope in the radio observations, the radio spectrum obtained for GX 339–4 shows distinctly more features. To account for these features we will stop the jet at a length of $10^{14.2}$ cm, consistent with Maitra et al. (2009).

4.4.1 Approach

To explore the possibility of a single physical jet-related mechanism behind the hysteresis, we will start out fitting the rise dataset. Once we locate possible solutions to the rise data we will fit the decay data starting from those solutions, preferably changing the least possible number of model parameters (changing all the parameters would mean the resulting fits are completely unrelated). The choice to start with the rise data set is not only motivated by the temporal sequence of these events, but also by the practical fact that the decay phase shows much less features in the X-ray spectrum: a nearly a perfect power-law with an iron line complex at 6.4 keV and perhaps some reflection. This data set also comprises less radio data points. As a consequence it will be much easier to obtain a fit for the decay data while changing few rise data settings, than the other way around.

To achieve the above we have a limited number of parameters at our disposal, distinguishable in two categories: Geometry-related (location of shock acceleration front and jet-base radius) and energy-related (jet normalization, magnetic domination, lepton temperature, particle distribution index and acceleration rate).
4.4.2 Spectral fitting of the decay phase

As mentioned in the previous Section, we want to keep the highest number of parameters fixed. However we can consider varying some of the parameters which are not directly related to the jet. As explained above, in the HS the SED is dominated by the corona/ base of the jet and hence the disk and line properties are essentially secondary effects. But we do know the disk contribution is likely vary from the rise to the decay phase, as the source will have gone through a soft state, meaning probably at least the size of the disk will not be the same anymore. Varying the normalization of the Gaussian could be justified for the same reason, as the iron line complex is thought to stem from irradiation of disk material by the jet. The bigger the disk, the higher the line contribution. Going from the rise to the decay episode the broadband model flux is also reduced and hence the absolute normalization of the iron line will have to be reduced. Changing the line contribution is further motivated by the fact that the resolution of the PCA is very low (of order 1 keV) (Rau & Greiner 2003) and hence no significant physical statements about the iron line complex amplitude can be made anyway; the Gaussian is mainly included so that the MNW05 model continuum flux does not have to account for features usually located near $\sim 6.4$ keV. However

After these considerations we decided to vary the Gaussian normalization, but to keep the results for the energy and the line width from the rise phase. However we chose not to vary other disk related parameters, such as the inner disk radius $r_{\text{in}}$ and the contribution of the reflection component $\Omega/2\pi$, for the reasons we mentioned in Section 4.4: the disk is a minor component in our paradigm and we want to avoid looking into all the detailed changes in the physics of the disk – from one phase to the other – and focus on the jet. Varying $r_{\text{in}}$ and $\Omega/2\pi$ as well would undoubtedly lead to statistically better fits, but will lead to less insight concerning the jet compelling the hysteresis. Hence for now we will refrain from exploring and quantifying the disk changes in more detail, which may be interesting for a future work.

The broadband luminosity (but in particular the IR flux, see Table 4.2) of the decay data is much lower than that of the rise data. This restricts the possible fits to methods that reduce the overall flux, with emphasis to reducing the IR flux. The main jet-parameters available to reduce the broadband continuum are the jet normalization $N_j$, the electron temperature $T_e$ and the magnetic domination $k$. The index of the particle distribution is also available to this effect, however in addition to changing the normalization of the optically thick radiation, the slope of the optically thin part of the jet will also vary, leading to a change in the X-ray spectrum. For the enhanced reduction of the IR we really only have one parameter at our disposal, i.e. the location of the acceleration front in the jet $z_{\text{acc}}$, after which a percentage of the leptons are accelerated into a power-law, determining the energy of the turnover point from optically thick to optically thin. Increasing the distance will shift the turnover point to lower energies and reduce the IR flux. Lastly me may try to change the acceleration rate $\epsilon_{\text{sc}}$ to change the energy where of the exponential cut-off of the optically thin synchrotron emission occurs.

4.5 Results

4.5.1 Rise phase

Using the MNW05 model, we were able to obtain 4 ostensibly different fits to the rise dataset (See Table 4.4 and Figure 4.5; see Figure 2.5 for explanation of the components in the Figures): The fits differ mainly in terms of the energy budget, namely, the magnetic dominance $k$, the electron temperature $T_e$ and the total energy going in to the jet $N_j$, but there is a natural degeneracy between these parameters, as to some extent they are able to compensate for one another. Looking at the contributions from the individual components in Figure 4.5 the fits are really quite similar and can thus be considered representative of a single class. The naming of the fits is done according

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5The parameters $p$ and $r_0$ are also involved in this degeneracy and can be considered less effective variants of the changes.
Table 4.4: Best-fit parameters obtained for the canonical black hole GX 339–4 rise phase hard state. For definitions of the parameters, see Table 2.1. The error bars have been resolved at 90 percent confidence level. We failed to resolve the error bars for parameters listed in italics.

<table>
<thead>
<tr>
<th>variable</th>
<th>units</th>
<th>LOEQUIP</th>
<th>LOMIDEQUIP</th>
<th>HIMIDEQUIP</th>
<th>HIEQUIP</th>
</tr>
</thead>
<tbody>
<tr>
<td>$N_{ij}$</td>
<td>$10^{-3}$ $L_{Edd}$</td>
<td>69.8 $^{+0.1}_{-0.1}$</td>
<td>52.1 $^{+0.3}_{-0.3}$</td>
<td>36.1 $^{+0.4}_{-0.4}$</td>
<td>29.1 $^{+0.4}_{-0.4}$</td>
</tr>
<tr>
<td>$r_0$</td>
<td>$GM/c^2$</td>
<td>25.8 $^{+0.1}_{-0.1}$</td>
<td>20.8 $^{+0.6}_{-1.7}$</td>
<td>14.2 $^{+0.4}_{-0.4}$</td>
<td>13.1 $^{+0.1}_{-0.1}$</td>
</tr>
<tr>
<td>$T_e$</td>
<td>$10^{10}$ K</td>
<td>4.92 $^{+0.01}_{-0.01}$</td>
<td>4.11 $^{+0.01}_{-0.10}$</td>
<td>3.29 $^{+0.07}_{-0.04}$</td>
<td>3.03 $^{+0.03}_{-0.06}$</td>
</tr>
<tr>
<td>$p$</td>
<td></td>
<td>2.23 $^{+0.00}_{-0.01}$</td>
<td>2.31 $^{+0.04}_{-0.01}$</td>
<td>2.33 $^{+0.03}_{-0.03}$</td>
<td>2.33 $^{+0.01}_{-0.01}$</td>
</tr>
<tr>
<td>$k$</td>
<td></td>
<td>3.29 $^{+0.09}_{-0.09}$</td>
<td>13.5 $^{+0.1}_{-0.1}$</td>
<td>65</td>
<td>142 $^{+1.1}_{-1.1}$</td>
</tr>
<tr>
<td>$z_{acc}$</td>
<td>$GM/c^2$</td>
<td>233 $^{+1}_{-40}$</td>
<td>163</td>
<td>232</td>
<td>206 $^{+15}_{-17}$</td>
</tr>
<tr>
<td>$\epsilon_{sc}$</td>
<td>$10^{-5}$</td>
<td>5.32 $^{+0.04}_{-0.32}$</td>
<td>4.49 $^{+0.89}_{-0.03}$</td>
<td>4.75 $^{+0.29}_{-0.08}$</td>
<td>4.49 $^{+0.36}_{-0.17}$</td>
</tr>
<tr>
<td>$\tau_{in}$</td>
<td>$GM/c^2$</td>
<td>--</td>
<td>--</td>
<td>26</td>
<td>61</td>
</tr>
<tr>
<td>$L_{disk}$</td>
<td>$10^{-3}$ $L_{Edd}$</td>
<td>--</td>
<td>--</td>
<td>39</td>
<td>22</td>
</tr>
<tr>
<td>$T_{disk}$</td>
<td>keV</td>
<td>--</td>
<td>--</td>
<td>0.25</td>
<td>0.14</td>
</tr>
<tr>
<td>$A_{HXT}$</td>
<td></td>
<td>0.89 $^{+0.01}_{-0.01}$</td>
<td>0.89 $^{+0.02}_{-0.01}$</td>
<td>0.89 $^{+0.01}_{-0.01}$</td>
<td>0.89 $^{+0.01}_{-0.01}$</td>
</tr>
<tr>
<td>$A_{line}$</td>
<td>$10^{-3}$</td>
<td>3.6 $^{+0.2}_{-0.2}$</td>
<td>3.9 $^{+0.1}_{-0.4}$</td>
<td>4.1 $^{+0.7}_{-0.8}$</td>
<td>4.0 $^{+1.9}_{-0.5}$</td>
</tr>
<tr>
<td>$E_{line}$</td>
<td>keV</td>
<td>6.22 $^{+0.08}_{-0.09}$</td>
<td>6.19 $^{+0.10}_{-0.09}$</td>
<td>6.12</td>
<td>6.15 $^{+0.11}_{-0.09}$</td>
</tr>
<tr>
<td>$\sigma$</td>
<td>keV</td>
<td>0.84 $^{+0.08}_{-0.08}$</td>
<td>0.87 $^{+0.08}_{-0.08}$</td>
<td>0.92 $^{+0.13}_{-0.08}$</td>
<td>0.96 $^{+0.06}_{-0.08}$</td>
</tr>
<tr>
<td>$\Omega/2\pi$</td>
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<td>5 $^{+1}_{-2}$</td>
<td>3 $^{+2}_{-2}$</td>
<td>6 $^{+2}_{-2}$</td>
<td>5 $^{+2}_{-2}$</td>
</tr>
<tr>
<td>$\chi^2/DoF$</td>
<td></td>
<td>165/81</td>
<td>162/81</td>
<td>152/81</td>
<td>148/81</td>
</tr>
<tr>
<td></td>
<td></td>
<td>(=2.03)</td>
<td>(=2.00)</td>
<td>(=1.88)</td>
<td>(=1.83)</td>
</tr>
</tbody>
</table>

$b$ Derived from model values for $r_{in}$ and $T_{disk}$.

The GX 339–4 hard state and HID hysteresis

to the amount of magnetic domination but there is no real physical motivation for this, other than that for GX 339–4 this is a convenient parameter to order the fits by: In GRS 1915+105 we did not know the order of magnitude of the magnetic domination to expect, however GX 339–4 has been explored many times before with the current paradigm (see Chapter 2) and those results suggest we want to avoid too high values for $k$ for consistency with the canonical black holes, of which GX 339–4 is one. Using the results of the rise fitting process as a basis we were able to obtain a number of statistically convincing fits to the decay phase data set, which we divided into three categories (see Section 4.5.2).

Looking at the reduced $\chi^2$ of the best-fit results for the rise phase we spot a minor trend: The higher the magnetic domination, the better the statistics. Although in this case the effect is minor, the same was seen for the GRS 1915+105 fits (see Table 3.3, SMEGDE and SMEGDE2 columns), where the effect is more pronounced. The two least magnetically dominated fits do not require an accretion disk, but this is related to the location of the acceleration front. If the acceleration front is closer to the base, the jet flux in the IR increases and less disk contribution is required. We can however not keep decreasing the acceleration front distance without changing the slope of the optically thin synchrotron spectrum, which in turn is constrained by the radio data. In theory we would be able to find fits with a disk contribution in all the regimes of the magnetic domination, but for the reasons mentioned in Section 4.4 this was not pursued.

It should be noted that we would normally expect the disk luminosity to be of the same order of magnitude as that of the jet. However the disk in GX 339–4 could very well radiate most of its flux in a regime not covered by the data (between optical to just below the lowest energy channel of RXTE), explaining why we are not able to constrain the disk properties (as shown in Table 4.4).

Interestingly the results are essentially compatible with reflection being absent ($\Omega/2\pi$ is only a few percent), compared to the 20–30% reflection we normally observe (see references in Chapter 2 and Section 3.5). The weak iron line complex (parametrized by $A_{line}$) is consistent with such a low amount of reflection and both parameters indicate a weak disk contribution.
Figure 4.5: Broad band continuum (left) and X-ray band only (right) best-fit results from the rise phase (from top to bottom, cf. Table 4.4: LOEQUIP, LOMIDEQUIP, HIMIDEQUIP and HIEQUIP columns respectively). See the caption of Figure 3.4 for the color-coding of the individual components.
2004: A possible state transition?

Looking at Figure 4.5 the larger residual of the highest-energy radio point is quite salient. If the decrease in flux at this frequency is real, a possible explanation is that GX 339–4 was already in a state transition during this observation. In such an event the flow may no longer be continuous near the base, in which case the jet would transition from optically thick to optically thin and possibly become “bloppy” rather than “steady”. This would eventually result in the synchrotron slope steepening at higher energy, first in the IR and then moving down to the radio. We see this happening explicitly in outbursts and hence we should consider this option here as well. An entirely different model may even be more suitable to fit the data, e.g. a synchrotron blob, i.e. a singular blob of plasma ejected from the accreting binary system. To explore the above possibilities we will explore two methods that are commonly used to make a state determination, analyzing the HID and the PDS, to ascertain if the rise observation is truly hard state.

The HID (Figure 4.3) shows behavior that could be interpreted as a failed state transition: Looking at the (magnified) blue panel in this Figure we see the X-ray luminosity decreasing, while the hardness increases again, shortly after the observation in question. Comparing this outburst to the extensively studied 2002-2003 outburst (see Figure 4.6) one would expect a much smoother track in this part of the diagram for a successful outburst. In addition, the flux level in 2004-2005 was a factor 4 lower the 2002-2003, perhaps explaining why the transition failed the first time. This failed transition may have happened close enough to the observation to influence the spectrum we are trying to fit. Hence we will have to look at the PDS of the rise observation to see if this indicates hard state.

We include the PDSs on the left hand side of Figure 4.7 to see if the rise (and the decay) phases are truly hard state. For our purpose a detailed timing analysis is unnecessary and beyond the scope of this thesis, so we refer you to the caption of that Figure 4.7 for a simple, qualitative analysis, that indicates both observations are indeed hard state. Hence we decided to explore the possibility that the rise data was perhaps obtained when GX 339–4 was already in transition to

**Figure 4.6:** HID (corrected to Crab-units) comparing the evolutionary tracks of the GX 339–4 2002-2003 (black) and 2004-2005 (purple; compare with Figure 4.3) outbursts and GRS 1915+105 class $\chi$ observations from Belloni et al. (2000) (red circles are plateau state and are the same as the stars in Figure 3.2, open circles are class $\chi_2$). The GX 339–4 observations have been shifted to put the source at the same distance as GRS 1915+105. Clearly the plateau state is too soft to lie on the hard state branch. The hard-to-soft state transition for the two GX 339–4 outbursts differs a factor three in count rate. Adapted from Belloni (2009).
Figure 4.7: Power Density Spectra of the RXTE X-ray data obtained during the rise (top left) and decay (bottom left) phases of the 2004-2005 GX 339–4 outburst, and of data obtained a few days before (top right, ObsID 90118-01-05-00, MJD 53080.7218) and after (bottom right, ObsID 80132-01-15-02, MJD 53086.5629) the rise phase observation, created using PCA channels 0-35, a stretch length of 256 s and a top (=Nyquist) frequency of 32 Hz. Left panel: Although the bottom left plot is of lower luminosity and has therefore larger errorbars then the top left one, both spectra clearly show all the classical hard state hallmarks: strong BLN at lower frequencies and a bending power-law at around ∼10 Hz, approximately 1.5-2 decades higher then the bend away from the α ∼ 0 slope. Detailed fitting of these PSDs with Lorentzians is beyond the scope of this work. Right panel: The two plots made from RXTE and on the right are almost indistinguishable from each other and from the top left panel, indicating it is unlikely that the rise observation was done during a state transition.

4.5.2 Decay phase

As explained in Section 4.4.2 we are only able to change the jet normalization $N_j$, the electron temperature $T_e$ and the magnetic domination $k$ – possibly in conjunction with the location of the

a softer state, by making a PDS for one observation shortly before and one shortly after the rise observation. These PDSs are on the right hand side of the same Figure. Since these latter two PDSs are both indicative of a hard state and look very similar to each other and the rise PDS, it is reasonable to assume the rise data itself is indeed a “classical” hard state.

In summary, a (failed) state transition does not provide an explanation for the deviation from a truly flat spectrum noted above. However we may find an explanation considering we are trying to explain the GX 339–4 rise data with a fully steady jet, while in reality there may very well be small fluctuations in the jet. The reduced flux in the highest-energy radio point may then have been the cause of such a “wave of variability”, consistent with adiabatic expansion, traveling outwards along the jet. Markoff et al. (2008) found evidence for such waves in the AGN M81*, where the timescales involved are much better suited to see such evolution in a jet, compared to XRBs (see Section 1.3), clearly confirming the strengths of being able to relate long-timescale AGN physics to short-timescale XRB physics.
not statistically convincing, or, comparing to previous works, the fits are not favorable because I required the change of four parameters, while the other two categories required only three

of solutions listed below.

The categories are simply numbered I, II and III in the order that we discovered them. Category I required the change of four parameters, while the other two categories required only three parameters to be modified. It is not possible to obtain a convincing decay fit in each of the three categories starting from each of the (four) individual rise fits: Sometimes those decay fits are either not statistically convincing, or, comparing to previous works, the fits are not favorable because the values for the parameters $k$ or $z_{\text{acc}}$ are much higher than normally required for the HS (viz. $k \gtrsim 200$ and $z_{\text{acc}} \gtrsim 10000$ $r_g$) and thus not convincing. However as discussed in Section 4.5.1 it really should not be necessary to get all these fits that differ only in the detailed distribution of the energy budget, as all the rise solutions can already be considered representing a single class. As explained, we are mainly interested in the qualitative underlying physical principles driving the hysteresis and behind the redistribution of energy when going from the rise to the decay phase, not in the detailed redistribution of energy from one rise phase energy budget to all possible others.

In line with the previous arguments we are free to choose any rise result as a starting point for the decay fitting. Hence we choose to use the rise result with the lowest value for the magnetic dominance (LOEQUIP; $k \sim 3.3$) as the main starting point for the decay fits, so that we are deviating the least from the values obtained for this parameter in previous works fitting GX 339–4 ($k \sim 2$; see references in Chapter 2). However for category III using LOEQUIP as a starting point

<table>
<thead>
<tr>
<th>variable</th>
<th>units</th>
<th>LOEQUIP</th>
<th>CAT I</th>
<th>CAT II</th>
<th>HIEQUIP</th>
<th>CAT III</th>
</tr>
</thead>
<tbody>
<tr>
<td>$N_j$</td>
<td>$10^{-5} \times L_{\text{Edd}}$</td>
<td>$69.8^{+1.7}_{-0.1}$</td>
<td>$74.3^{+0.1}_{-0.0}$</td>
<td>$19.9^{+0.1}_{-0.6}$</td>
<td>$29.1^{+0.3}_{-0.1}$</td>
<td>$14.1^{+0.1}_{-0.1}$</td>
</tr>
<tr>
<td>$r_0$</td>
<td>$G M/c^2$</td>
<td>$25.8^{+0.1}_{-0.0}$</td>
<td>$3.24^{+0.02}_{-0.01}$</td>
<td>$3.9^{+0.03}_{-0.06}$</td>
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<td>$1.87^{+0.01}_{-0.00}$</td>
</tr>
<tr>
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<td>$10^{10}$ K</td>
<td>$4.92^{+0.01}_{-0.00}$</td>
<td>$3.29^{+0.00}_{-0.09}$</td>
<td>$19.8^{+1.8}_{-1.6}$</td>
<td>$142^{+11}_{12}$</td>
<td>$7819$</td>
</tr>
<tr>
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<td></td>
<td></td>
<td>$206^{+37}_{14}$</td>
<td></td>
</tr>
<tr>
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<td></td>
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<td></td>
<td></td>
<td>$7819$</td>
<td></td>
</tr>
<tr>
<td>$z_{\text{acc}}$</td>
<td>$G M/c^2$</td>
<td>$233^{+10}_{-0}$</td>
<td>$2998^{+5}_{-3}$</td>
<td>$800$</td>
<td>$148^{+11}_{0.1}$</td>
<td>$0.5^{+0.1}$</td>
</tr>
<tr>
<td>$\epsilon_{\text{sc}}$</td>
<td>$10^{-5}$</td>
<td>$5.32^{+0.04}_{-0.32}$</td>
<td>$3.5^{+0.1}_{-0.5} \times 10^4$</td>
<td>$0.5^{+0.1}_{-0.0}$</td>
<td>$4.0^{+1.0}_{0.5}$</td>
<td>$0.5^{+0.1}$</td>
</tr>
<tr>
<td>$A_{\text{line}}$</td>
<td>$10^{-3}$</td>
<td>$3.6^{+0.2}_{-0.3}$</td>
<td>$0.3^{+0.1}_{-0.1}$</td>
<td></td>
<td></td>
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<tr>
<td>$\chi^2/\text{DoF}$</td>
<td></td>
<td></td>
<td>$165/81$</td>
<td>$162/81$</td>
<td>$152/81$</td>
<td>$148/81$</td>
</tr>
</tbody>
</table>

Table 4.5: Best-fit parameters obtained for the canonical black hole GX 339–4 decay phase hard state. For definitions of the parameters, see Table 2.1. The error bars have been resolved at 90 percent confidence level. We failed to resolve the error bars for parameters listed in italics. The values for the LOEQUIP and HIEQUIP columns are the same as in Table 4.4. If an entry in one of the CAT columns this Table is missing, the value is that of the rise result column to its left hand side.

<table>
<thead>
<tr>
<th>fit</th>
<th>$n$ ($10^{14} \text{ cm}^{-3}$)</th>
<th>$B$ ($10^8 \text{ Gauss}$)</th>
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<td>LOEQUIP</td>
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<td>1.6</td>
</tr>
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<td>HIMIDEQUIP</td>
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<td>3.9</td>
</tr>
<tr>
<td>HIEQUIP</td>
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<td>5.3</td>
</tr>
<tr>
<td>CAT I</td>
<td>0.5</td>
<td>0.2</td>
</tr>
<tr>
<td>CAT II</td>
<td>1.1</td>
<td>1.1</td>
</tr>
<tr>
<td>CAT III</td>
<td>2.9</td>
<td>3.8</td>
</tr>
</tbody>
</table>

Table 4.6: Jet base lepton density $n$ and magnetic field strength $B$ for all the best-fit results. The values are consistent with those predicted in Markoff et al. (2001)

acceleration region $z_{\text{acc}}$ and the acceleration rate $\epsilon_{\text{sc}}$ – to obtain the solutions in this Section. As put forward in Section 4.4.1 the objective is to change as few of the above parameters as possible in order to see if we can spot an underlying trend governing the changes in physics when going from the rise to decay phase. With these boundary conditions we are able to find the three categories of solutions listed below.
requires the acceleration zone distance to increase to $z_{\text{acc}} \sim 15,000 \ r_g$ (meaning the location of the acceleration region is even much further out than in the already extreme GRS 1915+105) and results in a bad infrared IR fit (too low normalization). Hence for category III a different starting point is more appropriate. We try the other three rise results and use the one with the best statistics (HIEQUIP).

The final best-fit results for the decay phase HS are in Figures 4.8–4.10 (see Figure 2.5 for explanation of the components) and Table 4.5. All the fits show an increasing, systemic deviation in the highest energy band (above $\sim 20 \text{ keV}$). In Table 4.5 we list only the parameters left free to vary and we refer you to Table 4.4 for the other (fixed) values.

I. Electron temperature, location of acceleration region, jet-base radius and acceleration rate

This fit has gone from a mix of synchrotron and SSC during the rise phase to almost entirely synchrotron post-shock. The reduced IR flux required the post shock jet to be pushed away from the base. To compensate for the reduction in SSC peak-energy the acceleration rate has increased, moving the post-shock cut off out of the spectrum and creating a power-law spectrum all the way through the X-rays. Meanwhile the lower energy X-ray flux also increases and so the base-radius must increase to reduce the SSC normalization.

Compared to the rise, the biggest changes are in the acceleration parameter $\epsilon_{\text{sc}}$, the acceleration front distance $z_{\text{acc}}$ and the jet-base radius $r_0$. The variation in electron temperature $T_e$ is far less significant. The decrease in the latter can be understood from the fact that $r_0$ has increased. We expect the jet base radius to be commensurate the inner edge of the thin accretion disk where the cool outer flow transitions into the hot inner flow. If $r_0$ increases, we expect the disk to be truncated at a larger radius. The electron temperature should correspond to the disk plasma temperature in the region where the jet is launched, so if the disk is further out, the electron temperature should go down. The larger $r_0$/truncation radius at lower luminosity is consistent with the “classical” picture suggested by Esin et al. (1997) (see Figure 1.7).

The jet radius increases, while the power fed into the jet stays the same. Hence the particle and magnetic energy densities decrease at the base (see Table 4.6). Along the jet the magnetic field strength evolves as $B \propto z$ ($z$ is the coordinate along the jet) and the acceleration region $z_{\text{acc}}$ has also moved further out along the jet. These two effects conspire so that the leptons radiate less along the jet and in the acceleration front than during the rise, so the cooling throughout the jet due to synchrotron losses becomes less efficient (see Equation 2.2). We expect a lower cooling rate to be consistent with an increase in synchrotron cut off energy. This increase should however also depend on changes in the acceleration efficiency. For Fermi acceleration, the rate (Equation 2.5) is dependent on the magnetic field $B$ and the plasma parameter $\epsilon_{\text{sc}}$, but as explained in Section 2.2,
\( \epsilon_{sc} \) is a “fudge” factor, used to get an idea of the acceleration efficiency given our lack of knowledge of the true acceleration process. The fits show the acceleration parameter \( \epsilon_{sc} \) has increased by 4 orders of magnitude, while the magnetic field has only gone down by one order. If we assume the acceleration process to be Fermi, the acceleration rate will increase, allowing for the increased synchrotron cut off we observe. However the high value for \( \epsilon_{sc} \) would indicate that the acceleration process is not strictly Fermi acceleration (as do the low values of \( \sim 10^{-5} \) for this parameter in all the other fits in this Chapter). In addition, the fact that the magnetic field strength does appear to influence the cut off frequency (contrary to Equation 2.7) also suggests that synchrotron losses do not dominate the cooling. However these issues will definitely need further investigation before we can make any firm statements about this. As the premise of this Chapter was only to explore if we could find interesting trends, we will leave it as is and leave these issues for a possible future work.

II. Jet normalization, magnetic domination and location of acceleration region

![Figure 4.9: Results from fitting the decay phase, category II.](image)

This fit is actually mainly an overall decrease in broadband flux when compared to the rise phase. The jet power decreases and the energy budget is shifted toward magnetic domination, in order to get the correct amount of SSC contribution in the X-rays, but the peak energy of the SSC stays the same. Again the acceleration region is pushed away from the base, but a factor three less as in the previous category. This decay fit resembles the rise phase fit the most in terms of component contributions.

The lower jet normalization reflects the fact that the system is now accreting at a reduced rate compared to the rise phase. The magnetic field has increased slightly in strength (see Table 4.6), however the change does not appear to be enough to change the electron temperature at the base to first order. Consistent with the unchanged electron temperature the jet base radius also has not changed (in contrast to both values changing in cat 1). The acceleration front has moved out but we can not be sure of the physical mechanism driving this. There appear to be no contradictions in this fit but all we can really say physically is that the energy is redistributed towards magnetic domination, with a value quite high for GX 339–4 or the other canonical BHBs \( (k_{\text{max}} \sim 7, \text{see Table 3.3}) \).

III. Jet normalization, particle distribution index and location of acceleration region

Of all the decay fits, this fit requires the biggest increase (almost by a factor of 40) in \( z_{acc} \). This is because the reduction in jet power required in the X-rays must be compensated in the radio by making the synchrotron spectrum shallower. This in turn pushes the acceleration region even further out so the synchrotron X-ray does not become too high. The acceleration region is now so far removed that the IR flux is reduced too much and must be compensated by another component,
Discussion

Figure 4.10: Results from fitting the decay phase, category III.

i.e. the disk. Although the disk contribution is now quite high, the power radiated by the disk is still of the same order as the power going into the jet (within a factor of 3). But considering the low reflection and iron line contributions such a luminous disk is not expected.

The base radius shows little change and so, as in the previous scenario, we expect the electron temperature to remain unchanged to first order as well. The flatter particle distribution index would indicate that either the shock is more relativistic but can perhaps also be understood from the picture of lepton cooling by synchrotron losses which will be discussed in the next Chapter. If the cooling is less efficient we may not be seeing past a cooling break (see below). The obtained reduction for the index is however \( \sim 0.5 \), while the expected reduction for a cooling break is unity. The case of steepening by unity is however only expected for “perfect” acceleration and there are many intermediate cases, depending on e.g. if particles can escape from the shock through various diffusive processes or turbulent/stochastic effects. At any rate, the new value of \( p = 1.87 \) is fully consistent with the value expected for relativistic diffusive shock acceleration (1.5-2; Heavens & Drury 1988), even more so if we consider the fact that the acceleration front has moved so far out (to a value comparable to those obtained for GRS 1915+105) that cooling is no longer important in the shocked segments and we are looking directly at the uncooled particle distributions (also see Section 3.5 for similar arguments concerning GRS 1915+105).

Additional categories

The work on GRS 1915+105 implies it may be worthwhile to check if a category of solutions can be obtained changing the electron temperature, magnetic domination and location of acceleration region, because the former two appeared to be inversely proportional in this system (see Section 3.5). However, when fitting the spectrum, decreasing the electron temperature also requires a decrease in jet-base radius, as the peak energy of the Compton hump also decreases as a consequence, which would in turn raise the electron temperature again. Some initial fits were attempted, but unsuccessful, perhaps due to this mechanism, and because the objective of this Section was to find a solution changing the least number of parameters, this approach was not further explored.

4.6 Discussion

Although we confirmed both the rise and decay phase observation as hard state (the rise phase more extensively than the decay phase), we see quite some difference in the spectra. Not only is the decay phase less luminous, it also shows a simpler X-ray spectrum that can even be fit using only (post-acceleration front) synchrotron, but is possible to fit using a mix of synchrotron and SSC. We have a limited number of physical parameters available in our jet to account for the spectral differences and consequently explain the hysteresis in the HID using our paradigm.
The HID hysteresis was first observed and described by Miyamoto et al. (1995). Since the launch of RXTE this characteristic evolutionary track is found to be similar for most black hole transients. A number of authors has tried to find the root for this phenomenon, explaining it as the result of varying physical causes. As for explaining the X-ray states, historically people have again looked mainly at the accretion disk/inflow to explain the hysteresis. However recently Corbel et al. (in prep.) have found a similar effect in the radio/X-ray correlation\(^6\), suggesting that the jet is of considerable importance.

In Section 4.5.1 we confirmed the rise observation as hard state. Looking again at the HID we could explain the “failed transition” behavior around the rise observation as a small outburst within that particular hard state, preceding the transition to the soft state. This interpretation would suggest that the flux level at which the hard-to-soft transition occurs depends on the accretion history (evident from Figure 4.6 because the source is leaving the hard state at different luminosity levels; Homan & Belloni 2005). So although the instantaneous accretion rate is no longer considered representative for a certain state (see Section 1.4), the history of the accretion rate may be responsible for the hysteresis. For example, (Smith et al. 2002) suggest a fast secondary accretion flow (in the form of a hot corona) and a slower (thin) disk flow would suffice for explaining the observations. In this picture it is easy to understand why hysteresis is not observed in all the canonical black holes. In an LMXB the disk is expected to be large because Roche-lobe overflow is connected with high angular momentum in the accreting material. A large disk suggests a delay before a change in the accretion rate at the outer edge of the disc has propagated all the way to the hot inner flow (Zdziarski et al. 2004). This would explain why the hysteresis is only observed in LMXB transients, and not in e.g. Cyg X-1 and Cyg X-3 (Maccarone & Coppi 2003). The disks in wind-accreting systems are expected to be much smaller and disturbances would be able to travel through an entire disk with little viscous delay. We suggest that in this physical mechanism the jet may again subsume the role of the (faster) Comptonizing corona, just as it has been shown in the original MNW05 paper to be able to reproduce the observed spectra usually attributed to hot inflow. The fits in cat ii and cat iii would be most suitable to investigate this matter, as the spectral composition is the same as that expected for the corona model, i.e. a synchrotron domination to \(\sim 10\) keV transcending into IC at higher energies. However going into the details of this suggestion is beyond the scope of this Chapter, which is meant to be exploratory.

In Smith et al. (2002) the arguments depend on the dominant cooling process in the corona being Comptonization of soft thin disk seed photons. Meyer-Hofmeister et al. (2005) attribute the hysteresis to changes in the Compton cooling (or heating) in the corona. Although the dominant cooling in our jet need not be due to Compton processes, this physical mechanism seems to be related to our cat iii fit (and to a lesser extent the cat i fit), which also appears to portray a change in the dominant cooling rate or process.

Zdziarski & Gierliński (2004) and Zdziarski et al. (2004) again look at the inflow for an explanation, attributing the hysteresis to two different regimes in radiative efficiency. The hard-to-soft transition starts with an optically thin inflow, while before the soft-to-hard transition the disk is optically thick, so the luminosity is reduced. The soft-to-hard transition seems to be instigated by a lower limit to the luminosity in the soft state, probably related to evaporation of the optically thick disk (Meyer et al. 2000). The limited range of possible luminosities for his transition (see Section 4.3) is also characteristic for other sources (Maccarone 2003). What constitutes the history determining the hard-to-soft state transition (to an optically thick disk) is however far less understood, as in GX 339–4 the transition is found to vary at least a factor three (see Figure 4.6) in luminosity. The maximum possible luminosity is probably limited by cooling of the hot plasma, but the transition is thought to not only depend on the existence of a hot flow, but also on the extent of the truncation of the surrounding cold disk. In all three fit categories of the decay phase, our jet has also changed in radiative efficiency from the rise to the decay phase, due to the increase in acceleration front distance, resulting in the jet being optically thin over a broader spectral band. This band is however dominated in intensity by the optically thick region

\(^6\)Interestingly, Coriat et al. (2009) found no evidence for hysteresis in the IR/X-ray correlation, suggesting the radio and the IR are connected to the X-ray differently.
before the acceleration front, which has grown in accordance. Hence a larger part of the jet at higher energies has become radiatively inefficient, consistent with the above works. Interestingly, Heinz & Sunyaev (2003) also found a similar scaling-relation for radiatively inefficient disk X-rays and uncooled synchrotron jet X-rays, which, assuming a relativistic shock is responsible for the value of $p$ is qualitatively consistent with what we see in cat III. The increase in optical depth may also provide a hint to why cat III required a higher magnetic domination to get a good fit: An increased optical depth is generally consistent with an increase in the Compton y-parameter and enhanced Comptonization. To reduce the normalization of the Compton “bump”, without changing its peak energy one normally increases the magnetic domination at the expense of the jet luminosity, which appeared required here.

The results in this work are consistent with those obtained in Markoff et al. (2003a), who were able to fit a number of GX 339–4 broadband data sets, changing only the acceleration zone distance $z_{\text{acc}}$ and the power input into the jet $N_j$. We varied a number of parameters, mainly related to the energy distribution$^7$, but are not able to fit any of the decay data without moving this zone further out in the jet, compared to the rise phase. Our values for $z_{\text{acc}}$ for the rise phase $\sim 200 \ r_g$ are consistent with those in the other GX 339–4 /MNW05 related works (see Section 2), but quite low compared to the values obtained in Markoff et al. (2003a) ($z_{\text{acc}} \sim 1.8 \times 10^3\ r_g$). The latter values are quite consistent to those we found for the decay phase.

4.7 Conclusions

In this Chapter we found three possible mechanisms driving the hysteresis, represented by three different categories of solutions to the decay phase data. For two of these categories the hysteresis would be caused by the change of three physical parameters, while the third category requires one additional parameter. All three categories show some interesting trends, that can be associated qualitatively with physical mechanisms already described in the literature. A quantitative analysis of the implications is however beyond the scope of this thesis, but the results in this Chapter provide a basis for future work in the form of a paper.

Clearly there is quite some degeneracy in the GX 339–4 rise fits that are used as starting point for the research in this Chapter. The degeneracy is partly due to the model, but also stems from the fact that some of the basic physical parameters in GX 339–4 are still quite uncertain. This problem could be alleviated by a high SNR detection of relativistic ejecta, allowing for tighter constraints on the inclination (largely determining the extent of inversion of the radio spectrum) and the mass (function), which are both still quite uncertain at this point. A direct spectroscopic detection of the companion star during an extended quiescent phase would also help constraining the mass, however this is much less likely, as attempts at this have been done, but were as of yet unsuccessful. For now, the system does not seem able to quiet down to a level where the contamination of the optical band caused by irradiation of the cold accretion disk reduces enough for the (putative) star becomes visible, but in due course GX 339–4 will undoubtedly go into outburst again and perhaps even create some blobs of plasma suitable for a proper detection.

Looking at the results for the decay phase, the main result is the required increase in acceleration front distance in all the fits. Also it appears we must consider cooling mechanisms to understand the evolution of GX 339–4 through the canonical states. For cat I we can not avoid discussing cooling to understand the results and for cat III changes in cooling appear to have had quite an effect on GX 339–4’s spectrum as it went from one hard state to the next. Although changes in the inflow and its cooling processes are just as likely to influence the spectrum and drive the hysteresis, in this work we will remain focussed on the outflow and motivated by the above results we will investigate the cooling processes in the jet in more detail in the next Chapter, expanding our paradigm to include cooling effects from synchrotron losses. As we shall see in that Chapter the value of $z_{\text{acc}}$ and the efficiency of the cooling in the jet due to these losses appear to be quite related.

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$^7$Basically, if the power input $N_j$ remained the same, we changed the use of the available power.
The synchrotron cooling break

5.1 Introduction

As hinted at in Section 2.2, energy losses for relativistic particles due to radiation can be quite significant and we expect the spectrum emitted by a relativistic jet to vary significantly depending on the exact cooling mechanisms at work. The MNW05 paradigm incorporates cooling due to adiabatic expansion and synchrotron and IC losses under the assumption of a constant injection of fresh, high energy power-law leptons in each segment after the acceleration front (see Chapter 2). The objective of this Chapter is to investigate what happens to the emitted synchrotron spectrum if the cooling is fast and the injection rate is not high enough to compensate for this cooling.

In other works with the MNW05 model the particle distribution indices is generally found to be $p \sim 2.5$, which is higher than expected if a relativistic shock is responsible for the acceleration ($p \sim 1.5 - 2$; Heavens & Drury 1988). This can be understood if we take into account synchrotron losses. If significant, these losses steepen the index of a steady power-law electron distribution by unity, from $p$ to $p + 1$. Equivalently, relating the particle distribution index $p$ to the spectral index $\alpha$

\[ \alpha = \frac{p + 1}{2} \]

Figure 5.1: Example expected accelerated synchrotron spectrum with cooling break implemented. The photon index $\Gamma$ increases by 1/2 (in this case above $\sim 0.02$ keV). At higher energy the photon index increases exponentially (in concordance with the acceleration rate, parametrized by $\epsilon_{\text{sc}}$).
the observed power-law index in the optically thin synchrotron spectrum should increase with \( \Delta \alpha = 0.5 \). The turnover point where the indices steepen is referred to as a \textit{cooling break} (Kardashev 1962). So in previous works using MNW05 where we observed \( p > 2 \) it was implicitly assumed that we were looking beyond a cooling break, i.e. the break lies below the X-rays (near the peak of the thermal electron distribution; see below) and we only observe the steepened power-law. It is however possible that the break occurs at a slightly higher energy and be visible in the optically thin synchrotron spectrum (see Figure 5.1).

Although MNW05 already incorporated synchrotron cooling to get an estimate for the maximum energy to which a lepton could be accelerated (see Section 2.2) we will now implement the above described break into the paradigm. We may find two new distinct classes of solutions: If this spectral break occurs at energies below the X-ray data, a new set of jet solutions should be possible, starting with a different range of values for \( p \). This solution set would still resemble the MNW05 solutions. However if the break occurs in the X-ray band, we may find an entirely different class of solutions.

5.1.1 Goal

Comparing the dynamical timescale of a system (\( t_{\text{dyn}} \); the time particles have to cool) to the synchrotron lifetime of the particles (\( t_{\text{syn}} \); i.e. the time such particle needs to radiate away most of its energy), we can get an estimate of the energy where the cooling break should occur. We are satisfied with an estimate as the cooling rate is time-dependent parameter, which we are essentially introducing in a steady-state code. We will incorporate the new cooling routine into the model as a switch, so that it can be turned on and of at will and we can compare the results.

We refer you to the appendix for the most important modifications done to the code. In this Chapter, we first show you how we calculate the break energy. We then show some additional calculations required to implement the cooling break as well as some other modifications made to the original paradigm in Section 5.2.2. We apply the final version of the modified paradigm to GX 339–4 hard state data (from Section 4.5.1) in Section 5.3 and discuss the model and its results in Section 5.4, before ending with our conclusions.

5.2 The new model

5.2.1 Cooling break energy

As stated above, in MNW05 injection of fresh leptons (with a distribution \( CE^{-p} \)) is assumed in each segment after the acceleration front. The time leptons have to cool in each of these segments, or the dynamical timescale of such a segment, is given by the time needed to traverse that segment. This \textit{residence} time can be calculated from the bulk speed, which is given by the jet velocity profile (see Figure 2.4; Falcke 1996; Falcke et al. 2009). The model has a subroutine \textit{jetpars} that calculates the various variables in each segment of the jet as a function of the geometry, magnetic field strength, particle density and lepton Lorentz factor in the nozzle. These variables include the bulk velocity in the jet expressed as a combination of the usual relativistic speed factors, \( \beta \gamma \). In the code this parameter is called \( rvel \). From \( rvel \) we can easily calculate the velocity of an average lepton in the flow.

\[
\nu_{\text{lepton}} = \beta c = \frac{rvel}{\sqrt{1 + rvel^2}} c.
\]  

The residence time is then simply the length of a segment divided by \( \nu_{\text{lepton}} \). We can qualitatively compare the residence time to the synchrotron lifetime of the electrons (see Figure 5.2).
Figure 5.2: Comparing the time an electron needs to emit most of its energy when subjected to a certain magnetic field strength (the synchrotron lifetime $t_{\text{syn}}$) to the timescale they reside in that field and hence have to cool (the dynamical timescale, or residence time $t_{\text{dyn}}$), we can estimate the change in the particle distributions and, consequently, possible changes in the spectrum that results from the distribution. Here we show examples of the lifetime of typical minimum, maximum and average energy power-law electrons and the increase of the residence time as each subsequent jet-segment decreases in magnetic field strength $B$ and increases in length $z$. The left hand side of these graphs corresponds to the acceleration front $z_{\text{acc}}$. When the lifetime of a maximum energy power-law electron exceeds the residence time in a certain segment, electron cooling is no longer relevant. Clearly, if $t_{\text{dyn}} \ll t_{\text{syn}}$, cooling has no significant effect. Only when the timescales are roughly of the same magnitude, cooling starts to become important and we see a break emerge, where the power-law steepens by unity. If $t_{\text{syn}} \ll t_{\text{dyn}}$, there will not be a break in the power-law anymore, but the entire power-law will have steepened with said factor.

We can estimate the break energy in the case of synchrotron dominated cooling from Equation 2.3, which gives the lepton energy as a function of time, using Equation 2.4, which gives the timescale of the energy loss. If the lepton remains relativistic we can assume $\gamma \gg 1$ and $\beta \approx 1$. We then see that because $t_{\text{syn}}$ is inversely proportional to the initial energy $E_0$ even leptons that have an infinite amount of energy at their disposal will end up with a maximum energy of

$$E_{\text{br}} \approx \frac{6\pi m_e^2 c^3}{\sigma_T B^2 T},$$

within a finite amount of time. The quantity $E_{\text{br}}$ gives an estimate for the energy where the synchrotron cooling break occurs. For the timescale $t$ we can use the residence time $t_{\text{dyn}}$. To calculate the break energy we assume the break energy in equation 5.3 to be exact.

In Section 2.2 where we calculated the energy equation, we assumed the synchrotron cooling was dominant in the acceleration region, consistent with Markoff et al. (2001). However as put forward above, there are other cooling processes at work in our paradigm. We can neglect losses due to Compton processes, as these are only considered important at the base of the jet, not in the acceleration region. Including adiabatic expansion and continuous injection of ”fresh” leptons with a spectrum $C E^{-p}$ the relevant energy equation becomes

$$\frac{\partial N}{\partial t} = -\frac{\partial}{\partial E} \left[ \dot{E} N \right] + C E^{-p}.$$  

From this equation we can derive the correct break energy for our jet. We will omit the complete derivation (see Kardashev 1962) here, and only state the end result, which increases from
The synchrotron cooling break

Figure 5.3: The relativistic Maxwellian, or Maxwell-Juttner distribution, for several temperatures. Below $T \sim 6 \times 10^9$, an individual electron would become sub-relativistic. When integrated over all $\gamma$, the distribution should approach unity. For the lowest temperature, the fraction of electrons exceeding $\gamma = 1$ becomes negligible.

the previous result, Equation 5.3 by a factor of four. The change in cooling break energy basically depends on the strength of two competing processes: Injection shifts the cooling break to higher energy with time, while adiabatic losses shift it to lower energy, in comparison with pure synchrotron losses. Including the correction factor our final break energy becomes

$$ E_{\text{br}} = \frac{24 \pi m e^2 c^3}{\sigma_T B t^2}. \quad (5.5) $$

We will now show the necessary calculations to implement the cooling break into the MNW05 paradigm.

5.2.2 Modifications to MNW05

As explained in Section 2.1 the total radiating lepton distribution in MNW05 not only comprises power-law distributions. The leptons in the jet start out in a quasi-thermal (relativistic Maxwellian, or Maxwell-Juttner; see Figure 5.3) distribution, described by a probability distribution for the speed

$$ F(E) = \frac{\beta t E}{mc^2} \sqrt{\frac{E}{mc^2}^2 - 1} \exp \left(-\beta t \frac{E}{mc^2}\right), \quad (5.6) $$

where $\beta t = mc^2/kT_e$ and $K2$ is the modified Bessel function of the second kind. Starting from the acceleration front a percentage of the thermal particles is continuously accelerated into a (single) power-law and hence, from this segment on, the total lepton distribution consists of an addition of the quasi-thermal distribution and a single power-law distribution (see Figure 5.4, top panel; ignore the discontinuity in the total lepton distribution for now, we will get back to this later). The normalization of both parts of the total distribution is determined by the fraction of particles that exist in each of these parts. After the acceleration front the total distribution gets shifted down, as in each subsequent jet segment the number density $n$ decreases according to adiabatic expansion (see Figure 5.5, bottom panel). The total distribution in each segment will radiate one of the blue-dashed post-shock components as shown in Figure 5.5, top panel; the most energetic component is radiated by the first acceleration zone and as we move out along the jet the spectral components decrease in energy, as explained in Chapter 2. Adding all the contributions from the distributions in each accelerated segment will result in the the green-dashed total synchrotron component.
Switch to broken power-law

For a single power-law distribution with a spectrum \( N_E dE \propto N_0 E^{-p} dE \) (where \( N_0 \) is the normalization constant), the proper normalization can be derived analytically from the lepton number density \( n \), using the fact that for a power-law with slope \(-p\) the total number density \( n_{\text{tot}} \) between \( E_{\text{min}} \) and \( E_{\text{max}} \) is given by the integral

\[
n_{\text{tot}} = \int_{E_{\text{min}}}^{E_{\text{max}}} N_0 E^{-p} dE, \tag{5.7}
\]

where \( N_0 \) is the normalization constant. Now that the power-law is broken, a normalization constant needs to be calculated for both the pre-break \(-p\) and post-break \(-(p+1)\) regimes. This can again be done analytically as long as we know the total number density, so that \( n_{\text{tot}} \) is given by

\[
n_{\text{tot}} = \int_{E_{\text{min}}}^{E_{\text{br}}} N_0 E^{-p} dE + \int_{E_{\text{br}}}^{E_{\text{max}}} N'_0 E^{-(p+1)} dE, \tag{5.8}
\]

where \( E_{\text{min}} \) is the minimum energy of the leptons in the power-law particle distribution, \( E_{\text{br}} \) is the energy where the cooling break occurs and \( E_{\text{max}} \) is the maximum energy of the leptons. The two constants \( N_0 \) and \( N'_0 \) are the normalization constants of the \(-p\) and \(-(p+1)\) distributions respectively. In order to have a “continuous” broken distribution. We require the particle density of both parts of the distribution to be the same at the cooling break energy:

\[
N_0 E_{\text{br}}^{-p} = N'_0 E_{\text{br}}^{-(p+1)}, \tag{5.9}
\]

so that

\[
N'_0 = N_0 \frac{E_{\text{br}}^{-p}}{E_{\text{br}}^{-(p+1)}} = N_0 E_{\text{br}}^{-p} (p+1) = N_0 E_{\text{br}}. \tag{5.10}
\]

Putting this into Equation 5.8 and rearranging, we calculate \( N_0 \)

\[
N_0 = \text{plfrac} \times n_{\text{tot}} \left( \frac{E_{\text{min}}^{1-p}}{p-1} + \frac{E_{\text{br}}^{1-p}}{p-1} + \frac{E_{\text{max}}^{1-p}}{p} - \frac{E_{\text{br}}^{1-p}}{p} \right)^{-1}, \tag{5.11}
\]

where \( \text{plfrac} \) is the fraction of leptons that is accelerated into the broken power-law. This expression is easily calculated in the code, and using Equation 5.10 also \( N'_0 \) is known. Using \( N_0 \) and \( N'_0 \) will result in a total acceleration front segment distribution similar to Figure 5.4, center panel.

Before we move on we note that although the normalization factors of our broken power-law are now determined, we can distinguish two regimes where we only need a single normalization constant: If the break energy \( E_{\text{br}} \) is above the maximum lepton energy \( E_{\text{max}} \) we see only a \(-p\) slope and if the break energy \( E_{\text{br}} \) is below the minimum lepton energy \( E_{\text{min}} \) we see only a \(-(p+1)\) slope. We again calculate the normalization constants using Equation 5.7. In the first case \( N_0 \) reduces to the same constant that is calculated in MNW05

\[
N_0 = \text{plfrac} \times n_{\text{tot}} \frac{1-p}{E_{\text{max}}^{1-p} - E_{\text{min}}^{1-p}}, \tag{5.12}
\]

while in the latter case the normalization constant will become

\[
N_0 = \text{plfrac} \times n_{\text{tot}} \frac{-p}{E_{\text{max}}^{1-p} - E_{\text{min}}^{1-p}}, \tag{5.13}
\]

Now that the normalizations of all possible (broken and un-broken) power-laws are set up properly we show how the total lepton distributions evolve along the jet and the resulting spectrum,
Figure 5.4: Lepton distributions at the first acceleration segment, as a function of Lorentz factor $\gamma$, showing changes made to the MNW05 model. For the break-less (top) distribution the shape of the combined distribution stays the same as the leptons travel along the jet, only the minimum power-law energy $E_{\text{min}}$ shifts down according to adiabatic cooling (Figure 5.5, bottom panel, demonstrates this for model parameters representing GX 339–4). The distributions get down-shifted in density and energy, as the jet evolves outward and subsequent segments cool adiabatically. The same happens for the (middle and bottom) distributions that include a cooling break, however for these, the break energy, where the power-law slope increases by unity, is calculated for each subsequent segment (cf. Figure 5.6), until the break energy exceeds the energy $E_{\text{max}}$ of the most energetic particles in the power-law. Note that for the calculation of all these distributions, an injection of fresh high energy particles is assumed in every segment (see Section 2.1). Also note the “bumps” in the power-laws due to the Maxwellian contribution.
Figure 5.5: GX 339–4 jet model spectrum (top panel), showing synchrotron radiation components in dashed light-blue, and the post-shock lepton distributions (bottom panel), obtained using the original MNW05 model. Only the post-shock lepton distributions are shown, as the pre-shock segment distributions are not yet accelerated into a power-law and would therefore not exhibit a cooling break. The top distribution represents the one in the first acceleration segment, or acceleration front. The subsequent distributions evolve according to adiabatic cooling as described in the caption of Figure 5.4. Clearly around $\gamma \sim$ the Maxwellian still contributes, resulting in a slight curve in the power-law slopes.
The synchrotron cooling break

Figure 5.6: Intermediate jet model spectrum (top panel; for this plot the GX 339–4 data has been omitted, as no fit was attempted using this intermediate result/lepton distribution) and post-acceleration zone lepton distributions (bottom panel). To obtain this result, the particle distribution index $p$ was already reduced by unity with respect to Figure 5.5 for comparison, so that the cooled power law regime of the post-acceleration synchrotron components are equally shallow. Hence in the cooled, highest energy synchrotron contributions shown in the top panel, the resulting spectral photon index $\Gamma$ is the same as in the aforementioned Figure. However the contributions from the segments that are not affected by cooling definitely show this diminished photon index, which emerges most saliently as the asymptote feature near $10^{16}$ Hz.
The new model

for typical GX 339–4 parameter values, in Figure 5.6, bottom panel and top panel respectively. Cooling definitely has a significant effect in this BHB. In the bottom panel we also see the break energy increasing as we get further removed from the acceleration front. This behavior can be understood from the following: Because the magnetic field $B$ and the residence time $t_r$ change in each jet segment, the break energy (Equation 5.5) also varies. Were the magnetic field to stay the same all along the jet, the cooling break would move to lower energy in the particle distribution because of the increase in residence time. However in our model we assume Poynting conservation of the magnetic field along the jet, i.e. it diminishes outward, as $B \propto r$. As the break energy depends on the magnetic field squared, cooling becomes less and less significant, until at a certain point the distribution looks again like a fresh, unaffected injection of electrons. This is what we see in Figures 5.6 and 5.7.

The spectrum in Figure 5.6 is only an intermediate result and has been included for clarity/comparison. We will however make some more modifications to MNW05 before using the expanded paradigm on actual data, because there remains another issue – shortly mentioned above – related to the relative normalizations of the quasi-thermal and power-law distributions: Looking at the top two plots in Figure 5.4 we see an intriguing feature, namely that of a jump in the total lepton MNW05 distribution at the minimum energy of the power-law. The issue was never considered of great importance, as the imprint it potentially has on the spectral results is tempered by the fact that it is only relevant in the optically thick part of the synchrotron spectrum, where, due to large self-absorption the effect is diminished and largely integrated out. However it remains unlikely that "step-functions" occur in nature. Now that we are modifying the original MNW05 distributions anyway we will explore if we can find a solution and try to either remove or reduce this minor flaw and investigate the influence it has on the emitted model spectrum.

Further normalization issues

As mentioned above, the normalizations of both the Maxwellian and the power-law are largely determined by the fraction of total particles in each respective distribution. In MNW05 this fraction is fixed to a value determined empirically from some initial BH SED fits. Through some kind of undetermined process, 75% is accelerated into a power-law\(^1\), while 25% remains in a Maxwell-Juttner.

The energy where the thermal distribution is transcending into the power-law is expected to occur at the peak of the Maxwell-Juttner distribution (Equation 5.6), at an energy of $E_{\text{min}} \sim 2.23kT$ (Maitra et al. 2009). Evidence for this comes from MHD simulations, that indicate the particles at the peak of the relativistic Maxwellian are most susceptible for acceleration. We refer to this energy as $E_{\text{min}}$ as it is the minimum energy of the particles in the power-law. After the peak, the thermal distribution falls of exponentially, but a significant fraction (depending on the exact temperature of the distribution; see Figure 5.3) of the thermal particles can exist in the Maxwellian at energies this peak. To this remainder the power-law is added is added, the normalization of which is mainly determined by the value of the peak energy $E_{\text{min}}$ (see e.g. Equation 5.12). Because the power-law cuts in instantaneously at $E_{\text{min}}$ this addition makes the final distribution discontinuous.

There are several methods to try and resolve this “inconsistency”, each with their own regimes of validity and other shortcomings. We first attempt to find the solutions without the power-law break. If successful we can always add the break later.

1. As put forward above we usually assume that the power-law starts at the peak of the Maxwellian, consistent with various MHD works. However this need not be true. The simplest solution would be to commence the power-law at $\gamma = 1$, in which case the two distributions are just added at all energies, so we are ensured of a smooth final distribution, and we will just see a minor “bump” at energies where the Maxwellian contributes significantly. However because of the MHD inconsistency, combined with the fact that the

\(^{1}\)Indeed, as far as astrophysical processes go, this is considered very efficient and this value may perhaps be on the high side. (Maitra et al. 2009)
Figure 5.7: GX 339-4 jet model spectrum (top panel) and final post-acceleration zone lepton distributions (bottom panel) including all modifications to MNW05. In the top panel note the increase in steepness of the lower-energy optically thin part of each post-acceleration synchrotron component with respect to those in Figure 5.6, due to the increased fraction of quasi-thermal particles. The exponential decrease above the peak of this Maxwell-Juttner distribution causes the total post-acceleration synchrotron components to gain an extra feature: Before, the intensity of the optically thick part decreased monotonically with a single value for the slope, related to $p$. Now the spectrum is steeper just after the transition from optically thick to optically thin while it becomes shallower at even higher energies.
resulting spectrum is completely different from those obtained with MNW05 we choose not to explore this option.

2. Another possibility depends on the assumption that most of the particles in the thermal distribution are below the peak energy (an assumption that becomes more accurate as the electron temperature increases), as above this energy the distribution decreases exponentially. We can now solve for \( E_{\text{min}} \), as the point where the Maxwellian cuts off immediately and the power-law starts. We get three equations to solve for three constants; one normalization constant for the Maxwellian, one for the power-law and the value of \( E_{\text{min}} \). The power-law normalization factor can be calculated semi-analytically from the integral Equation 5.12, leaving \( E_{\text{min}} \) as a variable. The Maxwellian normalization factor is calculated with an integral over the Maxwellian (Equation 5.6) from \( \gamma = 1 \) up to \( E_{\text{min}} \), again leaving \( E_{\text{min}} \). The third equation is obtained setting the values of the Maxwellian and the power-law equal at \( E_{\text{min}} \). The biggest drawback of this solution is that it only works for a limited parameter space. Because \( \text{plfrac} \) is fixed it is often the case that the Maxwellian and the power-law do not intersect and we will not be able to solve the above set of equations.

A similar method is used by Pe’er & Casella (2009), who ascertain they have a smooth particle distribution by choosing a Maxwellian with a dimensionless temperature of \( E_{\text{min}} / 2 \) such that the power-law connects up smoothly at \( E_{\text{min}} \). However, as above, at energies higher than \( E_{\text{min}} \) the Maxwellian does not contribute to the total distribution anymore and so the thermal component itself is again discontinuous (private communication with P. Casella). For this reason (and the limited range of validity explained above) we choose not to use these methods.

3. A third possible solution is to let the fraction of particles accelerated into a power-law (\( \text{plfrac} \)) vary, and to solve for that in stead of for \( E_{\text{min}} \). Initial exploration using the Mathematica software package indicate we may expect a reduction the accelerated fraction, to about a half. Considering an acceleration efficiency of 75 % is quite high this method seems promising. Indeed, the method has the drawback that the final solution is still discontinuous, albeit just by a factor of two (because the two component distributions have the same value at \( E_{\text{max}} \); see Figure 5.4, bottom panel), which we should be able to completely ignore considering the size of the discontinuity before was already ignored because of the optical depth argument explained above. The other drawback is that we have an extra free variable. Although the variable is worked out by the model itself (and depends greatly on the values for \( z_{\text{acc}} \) and \( p \), see Section 5.4), as we shall see below we get an entirely different model with distinct behavior. However considering the drawbacks of the other solutions we decide to go with this one, as in this case the model should at least be able to provide us with a fit under all circumstances.

To solve the equations in method 2 and 3 above, we implement a root finding routine into the code. The method of choice is the (multi-dimensional) Newton-Raphson, described in Press et al. (1992), section 9.7, as it does not require analytic derivatives for the Jacobian matrix.

Using method 3 we obtain an acceleration front distribution similar to Figure 5.4, bottom panel. The evolution of the lepton distribution shows the behavior depicted in Figure 5.7, bottom panel and the resulting spectrum is shown in the same Figure, top panel.

**Switching from analytic to numerical approach**

As explained above we choose to make the fraction of accelerated leptons variable, as this solution can be applied over the broadest range of parameter space, while making the least changes to the original model. Using the Equations in the above sections and the root finder routines we can analytically calculate how the parameters of the particle distributions, such as \( E_{\text{min}}, E_{\text{br}}, E_{\text{max}} \) and \( n \) change in each segment of the jet, according to the cooling processes and changes in physical parameters along the jet. However to do this analytically for each segment requires an enormous amount of CPU time on a computer. To prevent this MNW05 only calculates the lepton
distributions in the nozzle and at the acceleration front. After this front the general shape of the lepton distribution is assumed to remain the same, safe for the changes caused by the adiabatic cooling. Because of this assumption an analytic approach to calculating the distributions is not required for the segments after the acceleration front, but the change in the total distribution can be calculated quickly from the stored distributions in the previous segment, greatly reducing the required CPU time. Even with this approach MNW05 can take many days to finish a complete run and produce a usable fit. So in order for the new model to be usable we have to use a similar approach, which is done using some simple geometry, explained in Figure 5.8.

Making use of the fact that in each subsequent segment the break moves up in energy (for reasons explained below), we can write a simple routine (see Appendix A) that essentially only has to calculate the break energy in each segment and modifies the stored distribution (comprising an un-broken power-law) when three conditions are met:

- The power-law must have become fully dominant over quasi-thermal distribution (this only happens at an energy higher than the quasi-thermal peak).
- We are in an energy bin with an energy exceeding the break energy.
- The break energy $E_{br}$ is less than the maximum lepton energy $E_{max}$.

We want to be sure the power-law has become dominant over the quasi-thermal distribution to avoid an issue when $E_{br}$ is of the same order as the top quasi-thermal particles. As stated above the routine will start modifying the distributions when it has reached an energy bin exceeding the break energy. Because of the slight “bump” in the power-law regime – caused by the Maxwellian contribution to the total distribution, see e.g. Figure 5.4 – the newly calculated power-law will have too high a normalization if the routine starts at an energy where the quasi-thermal distribution is still relevant, because it will connect up to the distribution where ever it is activated. When an energy bin has an energy higher than $E_{br}$ the power-law is steepened by unity. As soon as $E_{br}$ exceeds $E_{max}$ we can stop modifying the single power-law distributions already calculated in the model.

If the energy of the break were to decrease along the jet we would encounter the problem that we can not account for the aforementioned “bump” in the lepton distribution caused by the Maxwellian contribution using this routine (basically the routine only calculates straight lines in a log-log plot).
5.3 The model applied to GX 339–4

Now that we have a working model, that is fast enough to be used, we can apply it to some data and see if we get some credible results. Considering the BHB discussed in Chapter 3, GRS 1915+105 is rather unpredictable, we apply the model to the GX 339–4 hard state (rise phase; see Chapter 4), for some preliminary results. Detailed fitting with this new paradigm is beyond the scope of this work.

The fit very much resembles the results for the same data set in Section 4.5.1, in terms of which components fit what. The statistics have slightly improved, by $\Delta \chi^2_{\text{red}} \sim 0.1$, which seems promising. We note the same feature in the optically thin synchrotron spectrum already spotted in Figure 5.7, namely that of a steepening of the power-law just above the turnover from optically thick to thin, followed by a second regime where the power-law becomes shallower. The steepening can be attributed to the larger fraction of thermal particles ($74\%$ thermal leptons, where it was $25\%$ before) making the exponential decay of the Maxwellian more dominant.

Most parameters are in fact similar to those listed in Table 4.4, taking again into consideration the fact that there is a natural degeneracy between the energy related parameters $N_j$, $T_e$ and $k$. However the values for the particle distribution index $p$ and the acceleration front $z_{\text{acc}}$ are notably different.

The acceleration front has moved inwards by a factor of 3. This reduction can be understood...
from the slightly better fit in the infrared: The front was able to move inwards without worsening the fit to the X-ray data, because of the new “bend” in the synchrotron power-law.

The change in $p$ was to be expected up to a certain extent. Considering the reduction in $z_{\text{acc}}$ we are likely looking beyond the cooling break, at a cooled spectrum, while the code returns particle distribution index as if it were uncooled. However the index has become shallower by a value greater than unity, suggesting an extremely relativistic shock. In fact the value for $p$ is even lower than the steepest value obtained for GRS 1915+105, ($p \sim 1.38$, see Table 3.3, SYNCH fit), for which we suggested we must be looking at the uncooled injected spectrum. It is however promising to see that $p$ has decreased by $\sim 1$, suggesting the new model does in fact work. We would need to explore the new parameter space quite a bit more before making any definite statements on the validity and applicability of the model (over the whole of parameter space) and about how the parameters (especially $z_{\text{sh}}$ and $p$) are related in this new paradigm, so although these preliminary results are very interesting, we will stop the discussion here and investigate a bit more how the model behaves.

Lastly, as mentioned in Section 2.2, the efficiency of the cooling is expected to change the maximum energy to which the leptons can be accelerated and hence the relative importance of the synchrotron and IC processes. Now that we adjusted the cooling, has this ratio changed? Comparing the value of $\epsilon_{\text{sc}}$ in the fit to the results in Table 4.4, the acceleration rate has indeed decreased, moving the cut off to lower energy and making the IC contribution slightly more dominant at higher energies. The effect is marginal though.

5.4 Discussion

Now that the model appears to be working, we should explore the extent of its applicability due to possible limitations in the allowed parameter space. The main difference with MNW05 is that we made the fraction of leptons accelerated into a power-law a variable to be solved by the code. To see if this change limits the model we plot how the accelerated fraction varies as a function of the acceleration front distance $z_{\text{acc}}$ and particle distribution index $p$ (see Figure 5.10), fixing the values of the other parameters to the ones in the Table of Figure 5.9. Of all the model parameters, $z_{\text{acc}}$ and $p$ are expected to influence the accelerated fraction the most, when keeping the other parameters fixed: $z_{\text{acc}}$ determines the magnetic field strength $B$ in the acceleration segments and hence the significance of the synchrotron cooling losses, while $p$ determines the slope of the power-law and hence determines how many leptons can “fit” under the power-law normalization of this distribution becomes to high to line up at the Maxwellian peak (within a factor of two, as described in Section 5.2.2).

In Figure 5.10, top panel, we see the fraction going to zero for $z_{\text{acc}} \sim 15 \, r_g$. At this point the cooling at the acceleration front has become so efficient that all the accelerated leptons have definitely ended up in a steepened distribution. Because the normalization of the Maxwellian is already determined by the number density at the base, it is impossible to accelerate particles into the steep power-law, because the power-law normalization would immediately become too high for it to line up with the Maxwellian. Hence all the particles remain in a thermal distribution. This could be a problem, as values for $z_{\text{acc}} \sim 7 \, r_g$ have been needed before. However the bottom limit of $z_{\text{acc}}$ has about the same value as the jet base radius $r_0$ (see Table in Figure 5.9). Of course the distance to the first acceleration zone is not expected to be smaller than the jet base radius. The fact that the minimum value for $z_{\text{acc}}$ seems to indicate the first acceleration zone is the first jet segment, provides us with a natural explanation as to why we get no power-law leptons and only a strictly thermal distribution, as all the leptons are required for the thermal distribution that is injected into the jet. The other end of the $z_{\text{acc}}$ range seems to give no problem at all, which is consistent with the cooling becoming insignificant when the magnetic field is too low and so the original MNW05 behavior returns. For the parameter values used, cooling becomes inefficient at $z_{\text{acc}} \sim 10^3 \, r_g$, after which the fraction does not change much anymore. This magnitude of $z_{\text{acc}}$ seems to support our conclusion that in GRS 1915+105 we may be looking at uncooled spectra, but to be sure of this we would have to verify this result with best-fit parameter values for this
Figure 5.10: Dependence of the fraction of accelerated particles on the location of the acceleration front $z_{acc}$ (top) and the particle distribution index $p$ (bottom), created with values typical for GX 339–4. **Top panel:** For the parameter values used, $z_{acc}$ develops a clear minimum at a distance $\sim 15$ $r_g$. The fraction accelerated makes a jump every time it increases, because the jet is divided in segments and $z_{acc}$ needs to be in a subsequent segment before a change is observed. To sample every jump over multiple orders of magnitude we had to change the step size $\Delta z$ a number of times, which is why not every jump is sampled an equal number of times. **Bottom panel:** The expected range in $p$ for a relativistic diffusive shock is indicated by the purple lines and suggests we can expect $\sim 20\%$ of the leptons gets accelerated into the broken power-law (for the used parameter range).

BHB, that, as shown in Chapter 3, can deviate quite a lot from the canonical GX 339–4. For now we will have to limit ourselves to the only fit values we have, those for GX 339–4 in Section 5.3.

As stated, we also looked the behavior of the particle distribution index $p$ (Figure 5.10, bottom panel). The same argument as above applies here: The steeper the index becomes, the “harder” it is to accelerate particles into the power-law. The “possible” range for $p$ is however far less than that deemed possible for $z_{acc}$, so the effect on the fraction is far less pronounced and the accelerated fraction never goes to zero. Over the limited range allowed for a relativistic shock the accelerated fraction hardly varies (with a value between 18-22%), so this parameter should not limit the model.

Comparing the expected total synchrotron spectra (Figure 5.1) to those in Figures 5.7 and 5.9, the results differ from what was expected beforehand. In stead of a steepening of the spectral index, we see the index becoming more shallow. This effect can be attributed to the fact that
The synchrotron cooling break has become more dominant with respect to the power-law lepton distribution, due to method used. The thermal domination hides the shallow power-law regime, because the exponential decay of the Maxwellian determines the shape of the spectrum as soon as the spectrum becomes optically thin, steepening the spectral index significantly. In the next regime of the spectrum, where it bends up again, we are already looking at the cooled spectrum. We find that the results of the model have become slightly less predictable using our method of varying the accelerated fraction of leptons. In Figure 5.10 we saw that depending on the value of $z_{\text{acc}}$ and $p$ the accelerated fraction can vary significantly. Hence the jet segments in the modified paradigm can potentially radiate spectral components anywhere between the purely thermal shape (e.g. Figure 5.7, the pre-acceleration zone blue-dashed components) and the old shape (e.g. Figure 5.5, blue-dashed post-acceleration front components), depending on the exact accelerated fraction. Although we will always see the inverted spectrum, in some cases (e.g. where $z_{\text{acc}}$ is very small) there may be no optically thin component at all anymore. This means the post-acceleration synchrotron components change in shape continuously because the involved parameters tend to vary significantly during the fitting process, making it much harder to fit the data by hand, which remains the most important step in the entire process.

5.5 Conclusions

The synchrotron cooling break is definitely of importance in the canonical BHB GX 339–4. At this stage it is too early to say if the same is true for other BHBs. The influence the cooling break has on the spectrum is much related to the distance to the acceleration region, because of the dependency of the break energy on the magnetic field and residence time. The results suggest that if the acceleration front is too far out in a jet, synchrotron cooling becomes too inefficient to modify the particle distributions. However if the cooling break occurs at an energy between the minimum and maximum energy a power-law lepton can have, based on acceleration and cooling rates and arguments, the effect on the spectrum can be dramatic.

We were only able to find one of the new categories of solutions mentioned in the introduction of this Chapter. Because of the alterations to the model the results are hard to compare with one another and with findings from the previous version of the model, due to the continuously changing shape of the post acceleration synchrotron spectrum. In the modified paradigm the shape of the synchrotron components is found to vary continuously between strictly quasi-thermal and the absorbed power-law spectrum we saw in MNW05, depending on the other jet parameters, most importantly $z_{\text{acc}}$ and $p$. The continuous variation makes the new model hard to use when fitting data.

With regards to the modifications we made to the distributions and their relative normalization in our paradigm we can conclude that completely removing the Maxwellian/power law jump appears impossible without introducing other physical objections. The best solution we found is to allow for an extra a free parameter, which makes sure the jump is reduced to a maximum of a factor two. This solution appears the best solution for now, as it is applicable to the widest range of parameter space, but has the drawback that the resulting synchrotron spectrum becomes more unpredictable. Considering the many orders of magnitude this parameter space can subsume (emerging already when only comparing the two BHBs featured in this work: GRS 1915+105 and GX 339–4) this is crucial.

As in the previous Chapter on GX 339–4, where changing the acceleration front distance was necessary in every mechanism behind the hysteresis, $z_{\text{acc}}$ again emerges as a most important parameter. In this Chapter the parameter reasserts its importance because synchrotron cooling and the energy of the cooling break depend on the strength of the magnetic field, which is largely determined by the distance of the acceleration front to the jet base. Although the physical parameters in a jet, like the distribution of the energy budget or the magnetic field strength, can vary over a large range, the jet structure seems to be the deciding factor in the phenomena we encounter. Should we be able to pin down, or at least gain more insight into the location of $z_{\text{acc}}$ using other arguments and/or simulations (e.g. MHD), jet-related research is sure to progress significantly.
Conclusions

As explained in Chapter 1, one of the “hot topics” at the current frontier of black hole related research is the extent to which the accretion physics driving the processes we see scales with mass. All the work in this thesis is to a higher or lesser degree related to this important issue.

In Chapter 3 we saw a positive correlation between the magnetic domination and the overall power input into the jet in the BHB GRS 1915+105. Evidence for a such a correlation was also found in the AGN M81+. In addition we found qualitatively similar results between GRS 1915+105 and the canonical BHBs on one hand, and the blazar sequence on the other, as they both show a decrease in energy range where the synchrotron power peaks with increasing jet power. To quantify the latter results we would need to investigate time-dependent effects, such as cooling breaks (see below), but both results suggest that the physics in black hole related phenomena is indeed scaling with mass.

In Chapter 4 we have been looking numerous times at the canonical states in the HID and how the transition between these states depends, with a varying degree, on luminosity. As explained in Chapter 1 an important part of current research is to find if the AGN spectral states map on to black hole states. The hardest problem to tackle concerning this topic is to correct for selection effects due to differences in luminosity and hence the observability of different classes of AGN. The most important result from this Chapter is the required increase in distance to the first acceleration zone $z_{\text{acc}}$ for all the decay phase fits. Also in Chapter 5 we found this parameter to be most influential on the results. Finding the location of $z_{\text{acc}}$, and its relation to other physical parameters, is a major topic in current MHD research (e.g. Polko et al., in prep.). The problem relates to understanding the field configuration before and after the soft state and hence understanding the soft-to-hard state transition. As noted in that Chapter, contrary to the hard-to-soft state transition, the soft-to-hard state transition occurs within a very limited luminosity range and is therefore better understood and seems therefore a logical place to start research away from steady, hard-state jets. As explained, AGN evolution progresses too slow for human observation and so the “intermediate” states in AGN can look like distinct classes of active galaxies. Hence trying to understand the transitional behavior is expected to shed light on the relation between AGN and BHB states and the extent to which the mass-scaling principle/assumption (see Section 1.3) is valid.

Cooling is important on every scale. (Heinz 2004) find evidence for radiative cooling in most of the AGN contributing to the “fundament plane” and emphasize the need to include synchrotron jet cooling to properly model the FP correlation. Above we mentioned that also the relation between the broadness of the synchrotron spectrum and input power in the jet required modeling
of cooling breaks. Although in Chapter 5 we made an estimate for the synchrotron cooling break to implement it in our steady state jet, in reality cooling processes are always time dependent and so in the end we will have to move towards using non-steady state models. Studying the detailed evolution of the lepton distribution throughout the jet would imply solving the Fokker-Planck equation for every segment and work on this is already in progress (e.g. Maitra et al. in prep.). Studying time-dependent cooling is also unavoidable if we want to understand the soft-to-hard and hard-to-soft state transitions and efforts along these lines will again contribute to our insight into AGN states and the relation to their XRB counterparts. However judging from the results in this thesis, it is not necessary to use hugely CPU intensive code to discover interesting trends in jet related physics. To fully understand how jets are formed and sustained, more detailed modeling of the processes involved will be necessary. However the trends found using our paradigm should at least provide some basic constraints these future efforts.
Delving deeper into the physics of jets we have encountered a number of issues that can not be resolved in a timescale reasonable for a Master thesis. Finding solutions for and understanding some of these issues may even prove a basis for a PhD, but the results obtained are interesting, to say the least, by themselves. In summary, the next tasks will obviously qualify for further investment of time:

- Chapter 3 will be submitted for publication to Monthly Notices of the Royal Astronomical Society (MNRAS).
- We need to investigate in detail the extent to which the jet can replace the hot inflow with regards to driving and explaining the HID hysteresis from Chapter 4. In order to do this we need to better understand the cooling and acceleration processes at work. The results of this investigation should be enough to base another scientific paper on.
- Concerning Chapter 5 we will most certainly have to use the developed paradigm to fit other hard state and GRS 1915+105 plateau state data, to check its validity and applicability.

For now we will leave these tasks for the future.
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Appendix A

Code Modifications

In this section the most important modifications to the code (excluding the root finder routines, for those see Press et al. 1992) are included for reference.

Calculate the break energy and the residence time in the acceleration segment

\[
\text{if}(\text{breaksw.eq.}1.0d0) \text{ then}
\begin{align*}
\text{brcst} &= 6.0d0 * \pi * \text{emgm} * \text{emgm} * \text{clite}**3 / \text{sigtom} \\
\text{rspeed} &= \text{rvel} / \sqrt{1 + \text{rvel} * \text{rvel}} * \text{clite} \\
\text{t_br} &= \text{delz} / \text{rspeed} \\
\text{e_br} &= 4.0d0 * \text{brcst} / (\text{t_br} * \text{bfield} * \text{bfield}) \\
\le_br &= \log_{10}(\text{e_br})
\end{align*}
\]

\text{c Calculate the plfrac and pl norms using root finder routines}

\[
\begin{align*}
\text{x(1)} &= 1.0d5 * 10.d0**((11.0d0 - 5.0d0)*\alpha) \\
\text{x(2)} &= 1.0d0 \\
\text{x(3)} &= 1.0d1 * 10.d0**((11.0d0 - 5.0d0)*\alpha) \\
\text{call newt(x,N,checkpoint)} \\
\text{call funcv(N,x,f)} \\
\text{if (checkpoint) then} \\
\text{ write(*,*) 'Convergence problems.' } \\
\text{endif} \\
\text{cnormbef} &= \text{x(1)} \\
\text{plfrac} &= \text{x(2)} \\
\text{cnormaft} &= \text{x(3)} \\
\text{thmfrac} &= 1.0d0 - \text{plfrac}
\end{align*}
\]

\text{c Calculate normalizations factor if break is lower than e_min or higher than e_max}

\[
\begin{align*}
\text{cnorm} &= \text{plfrac} * \text{ntot} * (1.0d0 - \alpha) / (\text{emax}**(1.0d0 - \alpha) - \text{emin}**(1.0d0 - \alpha)) \\
\text{cnormcool} &= \text{plfrac} * \text{ntot} * (-\alpha) / (\text{emax}**(-\alpha) - \text{emin}**(-\alpha)) \\
\text{end if}
\end{align*}
\]
At the acceleration front, make energy array between lelec(1) and log10(emax), add power-law and particle distribution together in appropriate ratios.

\[
ebot = \text{lelec}(1)
\]
\[
einc = (\log10(\text{emax}) - \text{ebot}) / \text{nelec}
\]
\[
e_{\text{mxw}}_{\text{m}} = \text{lelec(\text{nelec})}
\]
\[
do \ i = 1, \text{nelec}
\]
\[
etemp(i) = \text{ebot} + (i - 0.5) * \text{einc}
\]
\[
\text{if (etemp(i).le.lelec(nelec)) then}
\]
\[
edtrm = \text{thmfrac} * \text{mxwl}(10.0**etemp(i))
\]
\[
\text{mxwlcomp(i) = log10(edtrm)}
\]
\[
\text{else}
\]
\[
edtrm = 0.0
\]
\[
\text{mxwlcomp(i) = -200.0}
\]
\[
end if
\]
\[
\text{if (breaksw.eq.1.00 .and. e_br .lt. emax) then}
\]
\[
\text{if (le_br .lt. log10(emin)) then}
\]
\[
\text{pltrm = cnorm} * (10.0**etemp(i))**(alpha - 1.0)
\]
\[
\text{plcomp(i) = log10(pltrm)}
\]
\[
\text{else}
\]
\[
pltrm = 0.0
\]
\[
plcomp(i) = -200.0
\]
\[
end if
\]
\[
\text{else}
\]
\[
\text{if (etemp(i).ge.log10(emin).and.}
\]
\[
etemp(i).lt.log10(emin)) then
\]
\[
pltrm = cnormbef * (10.0**etemp(i))**(alpha)
\]
\[
plcomp(i) = log10(pltrm)
\]
\[
else
\]
\[
pltrm = cnormaft * (10.0**etemp(i))**(alpha - 1.0)
\]
\[
plcomp(i) = log10(pltrm)
\]
\[
end if
\]
\[
\text{else}
\]
\[
pltrm = 0.0
\]
\[
plcomp(i) = -200.0
\]
\[
end if
\]
\[
\text{if (etemp(i).ge.log10(emin)) then}
\]
\[
pltrm = cnorm * (10.0**etemp(i))**(alpha)
\]
\[
plcomp(i) = log10(pltrm)
\]
\[
else
\]
\[
pltrm = 0.0
\]
\[
plcomp(i) = -200.0
\]
\[
end if
\]
\[
dtemp(i) = \text{log10}(edtrm + pltrm)
\]
\[
\text{lelec(i) = etemp(i)}
\]
\[
\text{elen(i) = 10.0**lelec(i)}
\]
\[enddo\]
At this point a loop is started to calculate the lepton distributions in all the segments after the acceleration front: For every segment, because in each subsequent segment the break moves up in energy, if below the break energy use the stored particle distribution, when above steepen the power law by unity, but only use make modifications to stored distribution if the number density in the quasi-thermal dist has decreased (after the peak) to four magnitudes below that in the power-law.

```plaintext
if(breaksw.eq.1.d0)then
  t_br=delz/rspeed
  e_br=4.d0*brcst/(t_br*bfield*bfield)
  le_br=log10(e_br)
  do i=1,nelec
    if(mxwlcomp(i).lt.plcomp(i)-4.d0
      * .and.etemp(i).lt.le_br) then
      dtemp(i+1)=dtemp(i)-(etemp(i+1)-etemp(i))*alpha
    else if (mxwlcomp(i).lt.plcomp(i)-4.d0
      * .and.etemp(i).ge.le_br) then
      dtemp(i+1)=dtemp(i)-(etemp(i+1)-etemp(i))*(alpha+1.d0)
    else
      dtemp(i)=dtemp(i)
    end if
  end do
end if
```

Of course the distribution is renormalized after this so that the particle number density in each segment stays consistent. However this piece of the code is not modified.


BIBLIOGRAPHY


BIBLIOGRAPHY


