Black Hole X-ray Binaries: Radiation and High-Redshift Feedback

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Gillian Knevitt

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Theoretical Astrophysics Group
Department of Physics and Astronomy
University of Leicester
Abstract

The accretion of matter onto black holes results in their characteristic spectrum through which we can identify them and study their properties. Furthermore, this radiation can couple to their surroundings, resulting in complex interactions between black holes and their environments. In this thesis, I study the accreting properties of stellar mass black holes, and examine the effect that such interactions may have had on the early universe. I also consider the observational characteristics of the lowest luminosity stellar mass black hole binary systems in our own galaxy.

Approximately one billion years after the Big Bang, the universe underwent a huge baryonic phase change, in which neutral hydrogen became ionized by the first sources of radiation. Massive stars are thought to drive this process, but their ionizing lifetimes could have been extended by a later phase in their evolution: black hole X-ray binary formation. However, the extent of this enhancement is not known, and has been highly debated in recent literature. In this thesis, I show that X-ray binaries were unlikely to be present in sufficient numbers to exert a significant effect on the intergalactic medium. Using a stellar population synthesis model of a single starburst event, I show that radiation from X-ray binaries dominates the ionizing power of a cluster after the most massive stars have ended their lives. However, their high energy spectra and short lifetimes mean their ionizing timescales are too long for them to affect the progress of reionization. Even so, the high escape fraction of X-rays from galaxies still provides scope for low level heating and ionization of the distant intergalactic medium under different circumstances, such as in the context of continuous star formation.

I also assess the detectability of the dimmest black hole binary systems in the Milky Way. Using a catalogue of black hole binaries in our galaxy, I find that there is a statistically significant lack of short orbital period systems, when compared to the neutron star binary population. I show that these sources may be hidden from view, rather than being truly absent, due to radiatively inefficient accretion, in which energy is lost to the black hole. However, this conclusion requires that the switch to inefficient accretion occurs sharply at a threshold mass accretion rate. In the case of a smoother switch, alternative observational or evolutionary arguments must be put forward to explain this dearth.
Declaration

The work described in this thesis was undertaken between October 2010 and March 2014 while the author was a research student under the supervision of Dr Graham Wynn and Prof Mike Watson in the Department of Physics and Astronomy at the University of Leicester. The author acknowledges the support of an STFC studentship.

Please note that the author has changed her surname during this PhD from her maiden name of Gillian James, to Gillian Knevitt.

Two of the chapters in this thesis have been published:


Another chapter has been submitted for peer review:

For Oliver, and my family
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Introduction

The chances of finding out what’s really going on in the universe are so remote, the only thing to do is hang the sense of it and keep yourself occupied.

Douglas Adams, The Hitchikers Guide to the Galaxy
CHAPTER 1. INTRODUCTION

In 1784, John Michell, a natural philosopher and geologist, wrote a letter to Henry Cavendish postulating the existence of an object so massive that even light was prevented from escaping. The mathematician Pierre-Simon Laplace investigated the same concept in his book *Exposition du Systeme du Monde* in 1796 and, over a century later, Albert Einstein gave these massive objects mathematical grounding in General Relativity. Nevertheless, these “dark stars”, or black holes, as they later became known, were merely considered as mathematical curiosities.

However, the advent of X-ray astronomy led to the actual discovery of a black hole in our own Galaxy, Cygnus X-1 (Webster & Murdin, 1972). Since then, there have been over twenty confirmed black hole detections in the Milky Way (Ritter & Kolb, 2003). Furthermore, there is now compelling evidence for the presence of super-massive black holes at the centre of most galaxies (Lynden-Bell, 1969; Kormendy & Richstone, 1995; Magorrian et al., 1998; Merritt & Ferrarese, 2001)

Despite this success, we have only found a small fraction of the hundred million black holes that theorists predict to be present in the Milky Way alone (Shapiro & Teukolsky, 1983; van den Heuvel, 1992). This is because we cannot yet directly detect black holes. Instead, we can only infer their presence by the behaviour of material in their vicinity. Consequently, our detections of black holes are based on their gravitational influence on nearby stars (Harms et al., 1994; Kormendy & Richstone, 1995), and emission from interacting objects (Shakura & Sunyaev, 1973).

Arguably the most powerful tool for observing black holes is studying the properties of the material that orbits them. Black holes in binary systems often draw material from their companion star, which forms an accretion disc around them. Radiation from these accretion discs give black holes a distinct spectrum, from which we can gain significant insight into the properties of the black hole itself.

Furthermore, the effects of accretion onto black holes are likely to be far reaching. Radiation can heat and ionize the surroundings, and drive huge kinetic outflows and jets (see e.g. Done et al., 2007; Fabian, 2012). Strong correlations between the properties of supermassive black holes and their host galaxies suggest a physical connection between their evolutionary paths (Kormendy & Richstone, 1995; Magorrian et al., 1998). Thus, black holes may have had a vital role in shaping the universe we see today.

It is these wide ranging effects of accretion that this thesis will seek to understand. Firstly, I will examine the role of the brightest stellar mass black hole binaries in the early universe, and their potential to drive important feedback processes. In the later part of my thesis, I will then investigate the other extreme, studying the faintest black hole binary populations in our Galaxy. In particular, I will address the possibility that the nature of their accretion may hinder their detection.

In this introductory chapter, I provide an overview of the main physics behind accreting black holes in binary systems. This will include a discussion of the mass transfer
processes between the binary components and a mathematical grounding for the behaviour of accretion discs themselves. Following this, I describe the main observational features of black hole binaries, and the physical mechanisms through which their radiation couples to their surroundings. Finally, I provide a short review of the relevant recent research into the influence of X-rays on the early universe, with particular emphasis on the potential impact that radiation from black hole binaries may have had on their primordial surroundings.

1.1 WHAT IS AN X-RAY BINARY?

In 1962, American Science and Engineering (AS&E) launched an Aerobee 150 sounding rocket carrying an X-ray detector, with the aim of observing fluorescence from solar X-rays on the moon. While this search failed, they made an unexpected discovery which gave birth to a whole new field of astronomy: the first extra-solar X-ray source, Scorpius X-1 (Giacconi et al., 1962). In the next few years, scientists postulated that these X-rays were coming from the accretion of material onto a compact, massive object known as a neutron star (Hayakawa & Matsuoka, 1964; Shklovsky, 1967; Burbidge, 1972). The material donor was then found to be a companion star, as inferred from the spectra of optical signals emanating from the same location (Sandage et al., 1966). Over the next decade, many similar sources were discovered, using rockets, balloon borne missions and finally, in 1970, the first X-ray detecting satellite, UHURU. With this came more evidence in favour of mass exchanging binaries with a compact component. In particular, observations of the pulsating X-ray source, Cen X-3 (Giacconi et al., 1971; Schreier et al., 1972), showed that the delays of the pulse arrival times were in phase with the periodic eclipses. This was a clear indication that X-rays were being generated by accretion onto the eclipsing object. Finally, in 1972 a mass measurement of the compact X-ray source Cygnus X-1 was derived from its orbital parameters (Bolton, 1972; Webster & Murdin, 1972), showing it was too heavy to be a neutron star. It was, instead, the first detected black hole. The global picture arising from these studies was of a galactic population of X-ray binaries, defined as systems comprising a compact black hole or neutron star orbiting a normal star from which they accrete material to produce an X-ray dominated spectrum. An artist’s impression is shown in Figure 1.1.

It is useful to separate X-ray binaries into two types of system, with different characteristics. While the nature of the primary (the compact neutron star or black hole accretor) has some bearing on the source’s properties, more important is the nature of the donor star (the secondary) since its size controls the method and rate at which mass is transferred onto the primary. X-ray binaries are therefore historically split into two types as follows:

(i) A **Low Mass X-ray Binary** (LMXB) comprises a black hole or neutron star primary orbited by a $\lesssim 1.5M_\odot$ star. This star is usually type A or later.
(ii) A **High Mass X-ray Binary** (HMXB) comprises a black hole or neutron star primary orbiting a massive secondary with \( M \geq 3M_\odot \), typically a type O or B star. ¹

### 1.1.1 Origins

“Main sequence” stars are in hydrostatic equilibrium because the inwards pull of their own gravity balances the outwards pressure due to nuclear fusion. Consequently, when their fuel is depleted, they collapse under their own weight. Collapse continues until atoms are so tightly packed that electron degeneracy pressure begins to counteract the pull of gravity. This form of pressure is a consequence of the Pauli Exclusion Principle whereby two half-integer spin fermions cannot occupy the same quantum state. A star whose structure is maintained by this pressure is known as a White Dwarf. However, above the Chandrasekhar mass limit of \( \sim 1.4M_\odot \) (Chandrasekhar, 1931) electron degeneracy pressure is not strong enough to withstand gravity, so the star continues to contract until a further degeneracy pressure becomes significant, this time due to short-range neutron-neutron repulsion. If neutron degeneracy pressure can balance gravity, a compact object called a neutron star is formed. If not, the star undergoes utter collapse to form a black hole. This happens when its mass exceeds the Tolman-Oppenheimer-Volkoff limit (Tolman, 1939; Oppenheimer & Volkoff, 1939). Unlike the Chandrasekhar limit, this constraint is less certain, as it depends on an unknown equation of state. However, it is understood to be approximately \( 3M_\odot \) (e.g. Heiselberg & Pandharipande, 2000). Although Schwarzschild found a solution of general relativity that characterised a black hole in 1916, these objects remained hypothetical until mass measurements of X-ray binary components confirmed their existence (Webster & Murdin, 1972; Press & Teukolsky, 1977).

In a binary system, the most massive star runs out of fuel first, because it fuses hydro-

¹Note that secondaries with \( 3M_\odot \leq M \leq 10M_\odot \) are sometimes instead referred to as Intermediate Mass X-ray Binaries, a subclass that will not be independently addressed in this thesis.
gen at a higher rate in order to remain in hydrostatic equilibrium (main sequence lifetimes \( \propto M^{-2.5} \)). When this happens, it expands as hydrogen fusion is ignited in its outer layers, which are then lost to the secondary by accretion. High mass secondaries accrete all of this material, but lower mass companions spiral inwards, becoming engulfed in a common envelope. Finally, the primary goes supernova, and collapses to form a neutron star or black hole. For massive secondaries, a wide HMXB is formed, while for low mass secondaries the supernova ejects the common envelope and a close LMXB remains.

Alternative formation scenarios have been suggested, particularly for globular clusters, which contain more LMXBs than theories of binary evolution predict (Clark, 1975; Katz, 1975). In these populations, some of the X-ray binaries may have formed from the tidal capture of high velocity neutron stars or black holes by main sequence stars in the cluster (Fabian et al., 1975; Press & Teukolsky, 1977).

1.1.2 Mass transfer

Mass can be transferred from a secondary onto a compact primary in two different ways. For sufficiently large secondaries, or close enough orbits, Roche Lobe Overflow can occur, in which the primary pulls material away from its companion. Alternatively, the secondary may eject its mass as a wind, which is captured by the primary in a process known as Stellar Wind Accretion. The latter is more common for HMXBs, whereas LMXBs, which are typically in closer orbits, accrete via the former mechanism.

1.1.2.1. Roche lobe overflow

If we consider two stars of mass \( M_1 \) and \( M_2 \) in orbit about each other, then Newton’s generalisation of Kepler’s third law gives

\[
P_{\text{orb}}^2 = \frac{4\pi^2a^3}{G[M_1 + M_2]} \tag{1.1}
\]

where \( a \) is the separation between the centres of mass of the binary components. By introducing \( q = M_2/M_1 \) as the mass ratio of the two stars, this can be rearranged to express \( a \) as

\[
a = 3.5 \times 10^{10}\left(\frac{M_1}{M_\odot}\right)^{1/3}(1 + q)^{1/3}P_{\text{orb}}^{2/3}\text{ (hr) cm}. \tag{1.2}
\]

The potential of the binary system is given by the sum of the gravitational potentials of the two stars and the effective potential of the centrifugal force. In Cartesian coordinates, where the \( x \)-axis lies along the line of the centres, the \( y \)-axis is in the direction of orbital motion of the primary, and the \( z \)-axis is perpendicular to the orbital plane, the potential can be written as
Figure 1.2  **Roche geometry of a binary system.** This diagram shows the Roche lobes surrounding the primary (M$_1$) and secondary (M$_2$) in a binary system with separation $a$. $R_1$ and $R_2$ are the radii of spheres having the same volume as the lobes themselves, and $L_1$ is the inner Lagrange point, connecting the two lobes at a saddle point in the Roche potential.

\[
\Phi_R = \frac{GM_1}{(x^2+y^2+z^2)^{1/2}} - \frac{GM_2}{[(x-a)^2+y^2+z^2]^{1/2}} - \frac{1}{2} \frac{\Omega_{\text{orb}}^2}{\Omega^2} [(x-\mu a)^2 + y^2], \tag{1.3}
\]

(Pringle & Wade, 1985; Frank et al., 2002), where $G$ is the gravitational constant, $\mu = M_2/(M_1 + M_2)$ and $\Omega_{\text{orb}} = 2\pi/P_{\text{orb}}$. This is known as the Roche potential, after the French mathematician, Edouard Roche, who studied it in connection with the destruction or survival of planetary satellites.

Combining (1.1) and (1.3) shows that

\[
\Phi_R = \frac{GM_1}{a} F \left( \frac{x}{a}, \frac{y}{a}, q \right) \tag{1.4}
\]

i.e. that the lines of equipotential ($\Phi_R =$ constant) are functions of $q$ only and that their scale is then determined by the binary separation, $a$. The most important feature of these equipotentials is the figure of eight area shown in Figure 1.2. The inner Lagrange point, $L_1$, is a saddle point in $\Phi_R$ which connects the potential wells surrounding each star, known as Roche Lobes. This is the position where the effective attractions of the two stars exactly balance.

In Figure 1.2, both stars lie within their Roche Lobe, but if the secondary expands beyond $L_1$, mass will move from the secondary into the Roche lobe of the primary, since
it is easier to pass through this point than out of the system all together. A binary in which both stars lie within their Roche lobes is called a detached system, whereas a system where a secondary fills its Roche lobe is semi-detached.

The distance, $b_1$, of the $L_1$ point from the centre of the primary is labelled in Figure 1.2 and can be estimated to good accuracy by (Plavec & Kratochvil, 1964)

$$\frac{b_1}{a} = 0.500 - 0.277 \log q.$$ (1.5)

for $0.1 \leq q \leq 10$. The radii of spheres having the same volume as the lobes themselves, $R_1$ and $R_2$, are well approximated as (Eggleton, 1983)

$$\frac{R_2}{a} = \frac{0.49q^{2/3}}{0.6q^{2/3} + \ln(1 + q^{1/3})},$$ (1.6)

which is accurate to better than 1%. Here, $R_1$ can be found by replacing $q$ by $q^{-1}$.

To initiate mass transfer, there must be a physical mechanism which causes the secondary to fill its Roche lobe. Either its radius must increase or the binary orbit must shrink. There are phases in stellar evolution where rapid expansion can lead to a semi-detached binary. However, long term mass transfer in LMXB systems is driven by orbital decay. This can be dominated by either of two mechanisms:

(i) **Gravitational Radiation**: According to general relativity, the changing quadrupole moment of the orbiting bodies produces gravitational waves that carry energy and angular momentum from the orbit. This is believed to drive orbital decay for binaries with $P_{\text{orb}} < 2$ hours (King, 1988).

(ii) **Magnetic Braking**: When $P_{\text{orb}} \geq 2$ hours, a magnetically coupled stellar wind emanating from the secondary is the dominant angular momentum loss mechanism.

The mass transfer rate resulting from these two processes is given by (King, 1988):

$$-\dot{M}_2 = \begin{cases} 
10^{-10} \left( \frac{P(h)}{2} \right)^{-2/3} M_\odot \text{yr}^{-1} & \text{if } P(h) < 2, \\
6 \times 10^{-10} \left( \frac{P(h)}{3} \right)^{5/3} M_\odot \text{yr}^{-1} & \text{otherwise.}
\end{cases}$$ (1.7)

### 1.1.2.2. Wind accretion

For donor stars that reside well within their Roche lobes, mass transfer is possible via a strong stellar wind (e.g. Livio, 1988; Matsuda et al., 1991). This is particularly relevant for HMXBs, which often contain an early-type star in a close orbit. As mass is ejected, wind particles passing close enough to the compact object such that their kinetic energy is less than the gravitational potential energy, will be captured and ultimately accreted. This is an exceedingly inefficient process compared to Roche lobe overflow, because only
a fraction of the material ejected from the secondary can be accreted onto the primary. However, mass loss rates from early type O and B stars are often extremely high \((10^{-6} - 10^{-5} M_\odot \text{yr}^{-1})\), which is why sources powered in this way are observable.

### 1.1.3 Accretion energy

As material is accreted onto the primary, gravitational energy can be extracted. For a body of mass \(M\) and radius \(R_*\), the gravitational potential energy released by the accretion of mass \(m\) onto its surface is

\[
\Delta E_{\text{acc}} = \frac{GMm}{R_*}.
\]

Therefore, in the case of steady accretion at a rate of \(\dot{M}\), the accretion luminosity is

\[
L_{\text{acc}} = \frac{GM\dot{M}}{R_*}.
\]

For a neutron star, this can be expressed, using typical orders of magnitude, as

\[
L_{\text{acc}} = 1.3 \times 10^{36} \dot{M}_{16} (M/M_\odot)(10\text{km}/R_*) \text{erg s}^{-1}
\]

where \(\dot{M} = 10^{16} \dot{M}_{16} \text{ g s}^{-1}\). However, the energy released from accretion onto black holes is less clear because their “radius” is not a hard surface. Instead, it is the distance from the singularity inside which matter cannot escape, known as the Schwarzchild Radius\(^2\),

\[
R_* = \frac{2GM}{c^2}.
\]

In this case, it is possible for accretion energy to disappear into the hole rather than being radiated. We parametrize this uncertainty with the dimensionless quantity, \(\eta\); the efficiency at which the rest mass energy is converted into radiation as it accretes onto the hole. Therefore

\[
L_{\text{acc}} = 2\eta \frac{GMM\dot{M}}{R_*}.
\]

Then, substituting (1.11)

\[
L_{\text{acc}} = \eta \dot{M} c^2
\]

If material is accreted infinitesimally slowly, then \(\eta = 1\). However, realistic values of \(\eta\) are typically around 0.1 (Frank et al., 2002).

\(^2\)The surface at the Schwarzchild radius represents the event horizon in a non-rotating body. However, for a rotating black hole, the event horizon is inside the Schwarzchild Radius.
1.1.3.1. Eddington limit

A term that will be used regularly in this thesis is the Eddington limit. This is the maximum luminosity a body can achieve when there is balance between the force of radiation acting outward and the gravitational force acting inward, i.e.

\[
\frac{GMm_p}{r^2} = \frac{L\sigma_T}{4\pi cr^2},
\]  

(1.14)

where the left hand side of the equation represents the gravitational force a proton (mass \(m_p\)) and electron (mass \(m_e\)) would feel from an accreting body of mass \(M\), at radial distance \(r\). The electron’s mass is negligible compared to the proton mass, so has been neglected. The right hand side shows the force radiation exerts on free electrons through Thomson scattering. The outwards radial force on each electron is the rate at which it absorbs momentum, \(\sigma_T S/c\), where \(S = L/4\pi r^2\) is the radiation flux and \(\sigma_T = 6.7 \times 10^{-25}\) cm\(^2\) is the Thomson cross section. The attractive electrostatic Coulomb force between electrons and protons causes them to move out together, so this again represents a net force on an electron proton pair. Rearranging for luminosity, the Eddington limit is

\[
L_{\text{Edd}} = \frac{4\pi GMm_pc}{\sigma_T} \sim 1.3 \times 10^{38} (M/M_\odot) \text{ erg s}^{-1}.
\]  

(1.15)

If the luminosity exceeds the Eddington limit, then the outward pressure of radiation would exceed the inward gravitational attraction. Therefore accretion would stop, and the source would no longer be luminous, unless it radiated by other means. If other radiation was produced, by nuclear burning, for example, then the outer layers of material would be blown off and the source would become unstable.

1.2 A CLOSER LOOK AT ACCRETION

When gas flows through the \(L_1\) point it does not fall directly onto the primary because the binary system is rotating. On leaving \(L_1\), the stream of material tends towards a circular orbit, since this is the lowest energy orbit for which it can retain its angular momentum. The circularization radius \(R_{\text{circ}}\) is such that the specific angular momentum is the same as the material had when passing through \(L_1\). The circular velocity is therefore

\[
v_\phi(R_{\text{circ}}) = \left(\frac{GM_1}{R_{\text{circ}}}\right)^{1/2}.
\]  

(1.16)

Typically, \(R_{\text{circ}}\) is a factor of 2-3 times smaller than the lobe radius \(R_{L1}\) of the primary. However, unless \(R_{\text{circ}} < R_s\), the gas must lose angular momentum before it can be accreted. Conservation of angular momentum causes angular momentum to diffuse outwards as some material falls inwards. Subsequently, a disc forms around the compact object. The structure of this accretion disc is defined by a set of differential equations that
can be solved analytically under certain conditions. These solutions include a *steady thin disc* approximation that is suitable for many physical accretion flows, and an *Advection Dominated Accretion Flow* that may be present in systems with very low accretion rates. These will be described in the text below, as will the physical processes which lead to angular momentum transfer and disc instabilities.

### 1.2.1 Accretion disc structure

It is often the case that disc flow in a binary is confined so closely to its orbital plane that we can consider it as a two dimensional structure. Within these bounds, the motion of gas in the disc can be quantified by calculating the mass and momentum transfer rates through the disc.

#### 1.2.1.1. Mass transfer

The geometry of a thin disc is shown in Figure 1.3. An annulus of disc material lies between $R$ and $R + \Delta R$, of which an infinitesimally small element subtending $d\phi$ is highlighted in yellow, containing mass $dm$. The surface density of this element is $\Sigma = dm/(R\Delta R d\phi)$ and the mass, $m$, of the annulus can be found by integrating $dm$ around the disc:

$$m = \int_0^{2\pi} dm = \int_0^{2\pi} R\Delta R \Sigma d\phi = 2\pi R\Delta R \Sigma.$$  

(1.17)

The rate of change of the mass of the annulus is given by the net flow to and from neighbouring annuli, so can be written as

$$\frac{\partial}{\partial t}(2\pi R\Delta R \Sigma) = v_R(R,t)2\pi R \Sigma(R,t) - v_R(R + \Delta R, t)2\pi (R + \Delta R) \Sigma(R + \Delta R, t) \approx -2\pi \Delta R \frac{\partial}{\partial R} (R \Sigma v_R).$$  

(1.18)

In the limit $\Delta R \to 0$, this gives the *Mass Conservation Equation*:

$$R \frac{\partial \Sigma}{\partial t} + \frac{\partial}{\partial R} (R \Sigma v_R) = 0.$$  

(1.19)
1.2.1.2. Angular momentum transport

Since $v_\phi = R\Omega$, the angular momentum of the element $dm$ is $dL = Rdv_\phi = dm R^2 \Omega$. The total angular momentum of the annulus is therefore

$$L = \int_0^{2\pi} dL = \int_0^{2\pi} R^2 \Omega dm = \int_0^{2\pi} R^2 \Omega R \Delta R \Sigma d\phi$$

$$= 2\pi R \Delta R \Sigma R^2 \Omega. \quad (1.20)$$

In addition to $v_\phi$, the gas possesses a small radial "drift velocity" $v_R$, which is negative near the central compact object, so that mass is being accreted. $v_R$ is a function of both time and distance from the compact object, and leads to a change in the mass and angular momentum of the annulus over time. The rate of change of angular momentum can be written

$$\frac{\partial}{\partial t} (2\pi R \Delta R \Sigma R^2 \Omega) = v_R(R,t) 2\pi R \Sigma(R,t) R^2 \Omega(R) - v_R(R + \Delta R,t) 2\pi (R + \Delta R) \Sigma(R + \Delta R,t)(R + \Delta R)^2 \Omega(R + \Delta R) + \frac{\partial G}{\partial R} \Delta R$$

$$\simeq -2\pi \Delta R \frac{\partial}{\partial R} (R \Sigma v_R R^2 \Omega) + \frac{\partial G}{\partial R} \Delta R, \quad (1.21)$$
where the additional term involving $\partial G/\partial R$ takes account of transport due to the net effects of viscous torques. The total torque acting on an annulus at radius $R$ is given by:

\[ G(R,t) = 2\pi R\nu \Sigma R^2 \Omega' \]  

(1.22)

where a kinematical viscosity, $\nu$, has been introduced, the origins of which will be discussed in Section 1.2.2.

In the limit of $\Delta R \to 0$, (1.21) becomes

\[ R \frac{\partial}{\partial t} (\Sigma R^2 \Omega) + \frac{\partial}{\partial R} (R\Sigma v_R R^2 \Omega) = \frac{1}{2\pi} \frac{\partial G}{\partial R}. \]  

(1.23)

If we now assume $\partial \Omega/\partial t = 0$, which holds for orbits in a fixed gravitational potential, then, using (1.22), this can be written as

\[ R\Sigma v_R (R^2 \Omega)' = \frac{1}{2\pi} \frac{\partial G}{\partial R}. \]  

(1.24)

Combining (1.19) and (1.24) to eliminate $v_R$ yields

\[ R \frac{\partial \Sigma}{\partial t} = - \frac{\partial}{\partial R} (R\Sigma v_R) = - \frac{\partial}{\partial R} \left[ \frac{1}{2\pi(R^2 \Omega)'} \frac{\partial G}{\partial R} \right]. \]  

(1.25)

and finally, in the case of a Keplerian orbit, where $\Omega_K = (GM/R^3)^{1/2}$ then, substituting for $G$ from (1.22) gives the equation for angular momentum transport

\[ \frac{\partial \Sigma}{\partial t} = \frac{3}{R} \frac{\partial}{\partial R} \left\{ R^{1/2} \frac{\partial}{\partial R} [\nu \Sigma R^{1/2}] \right\}. \]  

(1.26)

This equation governs the evolution of the surface density $\Sigma$, over time. From (1.24), the radial velocity is shown to follow

\[ v_R = - \frac{3}{\Sigma R^{1/2}} \frac{\partial}{\partial R} [\nu \Sigma R^{1/2}]. \]  

(1.27)

1.2.1.3. The spreading disc solution

For further analysis, it would now be necessary to know the dependence of the viscosity parameter, $\nu$, which could be both radius and time varying. However, to gain some insight into the formation of an accretion disc, it is possible to continue this line of inquiry with the assumption that $\nu = \text{constant}$. In this case (1.26) can be written as

\[ \frac{\partial}{\partial t} (R^{1/2} \Sigma) = \frac{3\nu}{R} \left( R^{1/2} \frac{\partial}{\partial R} \right)^2 (R^{1/2} \Sigma) = \frac{12\nu}{s^2} \frac{\partial^2}{\partial s^2} (R^{1/2} \Sigma). \]  

(1.28)
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Figure 1.4  Viscously spreading disc. This figure shows a ring of matter of mass \( m \) placed in a Keplerian orbit at \( R = R_0 \) as it spreads out under the action of viscous torques. Different times are shown in terms of the dimensionless time variable, \( \tau = 12 \nu t R_0^{-2} \), where \( \nu \) is the constant kinematic viscosity. This plot is reproduced from the thesis of Fergus Wilson, University of Leicester, 2010, which in turn is based on Figure 5.1 in Frank et al. (2002).

where \( s = 2R^{1/2} \). Hence, if we write \( R^{1/2} \Sigma \) as the separable function: \( T(t) S(s) \), we find

\[
\frac{T'}{T} = \frac{12 \nu S''}{s^2 S} = \text{constant} = -\lambda^2, \tag{1.29}
\]

which we can solve to show that \( T \) and \( S \) are exponential and Bessel functions respectively.

We can use this equation to predict what would happen to material first passing through the \( L_1 \) point to form an initial matter distribution in a ring around \( R = R_0 \), with a surface density of \( \Sigma(R, t = 0) = (m/2\pi R_0) \delta(R - R_0) \), where \( \delta(R - R_0) \) is the Dirac delta function. Standard methods give

\[
\Sigma(x, \tau) = \frac{m}{\pi R_0^2} \tau^{-1} x^{-1/4} \exp\left\{ -\frac{(1+x^2)}{\tau} \right\} I_{1/4}(2x/\tau) \tag{1.30}
\]

where \( I_{1/4}(z) \) is a modified Bessel function and \( x = R/R_0, \tau = 12 \nu t R_0^{-2} \) are dimensionless radius and time variables respectively. Figure 1.4 shows \( \Sigma(x, \tau) \) as a function of \( x \) for various values of \( \tau \). Over time, the gas spreads out in both directions from its initial annulus, eventually forming a smooth disc with higher surface density nearer to the central object. Solving for the radial velocity gives two solutions that account for the asymptotic
behaviour of $I_{1/4}$:

$$v_R \sim \frac{3v}{R_0} \left\{ \frac{1}{4\chi} + \frac{2\chi - \frac{1}{2}}{\tau} \right\} > 0 \text{ for } 2\chi \gg \tau$$

$$v_R \sim -\frac{3v}{R_0} \left\{ \frac{1}{2\chi} - \frac{2\chi}{\tau} \right\} < 0 \text{ for } 2\chi \ll \tau.$$  \hspace{1cm} (1.31)

Matter at $2\chi \gg \tau$ moves outwards, taking away the angular momentum of the inner parts that move inwards towards the accreting object. At later times, $v_R$ changes sign at further distances, so material at even greater radii begins to drift inwards, eventually leading to the majority of mass being accreted while all of the original angular momentum is carried to large distances by a small fraction of the mass.

1.2.1.4. Steady thin discs

The most famous self-consistent solution of the time-dependent disc equations is that of a steady thin disc, which was developed by Shakura & Sunyaev (1973); Novikov & Thorne (1973) and Lynden-Bell & Pringle (1974) among others. Provided changes in disc structure only take place over long timescales, discs can settle into a steady state structure which can be examined by setting $\partial / \partial t = 0$ in the conservation equations (1.19) and (1.24). From (1.19) this gives

$$R \Sigma v_R = \text{constant},$$  \hspace{1cm} (1.32)

which represents the constant flow of mass through all points in the disc. As $v_R < 0$, this can be written as a mass accretion rate, $\dot{M}$ so that

$$\dot{M} = 2\pi R \Sigma (-v_R).$$  \hspace{1cm} (1.33)

Next, with $\partial / \partial t = 0$, (1.24) becomes

$$R \Sigma v_R R^2 \Omega = \frac{G}{2\pi} + \frac{C}{2\pi},$$  \hspace{1cm} (1.34)

where $C$ is the constant of integration. Substituting (1.22) for $G$ gives:

$$-v \Sigma \Omega' = \Sigma(-v_R) \Omega + \frac{C}{2\pi R^3}.$$  \hspace{1cm} (1.35)

The constant, $C$, relates to the rate at which angular momentum flows onto the central object. For the accretor to remain intact, its rotation rate must be less than Keplerian. Therefore, if the disc extends all the way to the stellar surface, there must be a point at which it begins a transition from angular momentum $\Omega_K$ to $\Omega_*$. If this occurs in a
boundary layer of thickness $b$, we can write

$$\Omega(R_\ast + b) = \left(\frac{GM}{R^3_\ast}\right)^{1/2} [1 + \mathcal{O}(b/R_\ast)].$$

(1.36)

Rearranging (1.35) for $C$ at $R = R_\ast + b$ gives

$$C = 2\pi R^2 \Sigma v_R \Omega(R_\ast + b)|_{R_\ast + b},$$

(1.37)

which, using (1.36) and (1.33) means

$$C = -\dot{M}(GMR_\ast)^{1/2}$$

(1.38)

to terms of order $b/R_\ast$. Consequently, if we substitute this back in for $C$ in (1.35), while setting $\Omega = \Omega_K$, we find

$$\nu \Sigma = \frac{M}{3\pi} \left[ 1 - \left(\frac{R_\ast}{R}\right)^{1/2} \right].$$

(1.39)

### 1.2.2 Viscosity

The evolution of the surface density of a disc is driven by the transport of angular momentum through viscous processes. However, although the kinematic viscosity term, $\nu$, is a central feature of the disc equations, both its magnitude and nature remain highly uncertain. The viscous timescale is the typical timescale for matter to diffuse through the disc via the effects of viscous torques, and is given by:

$$t_{\text{visc}} = \frac{r^2}{\nu}.$$

(1.40)

It is reasonable to expect that viscosity arises due to turbulent flow in the disc. However, the obvious choice of molecular viscosity as the driver of disc evolution can be ruled out because the timescales are too long (see Frank et al., 2002).

One of the most promising proposed mechanisms is the magneto-rotational instability (Balbus & Hawley, 1991; Brandenburg et al., 1995; Hawley et al., 1996). Here, elements of ionized disc material on the same frozen-in field lines behave as if attached by springs. As the innermost element orbits faster, it stretches the magnetic field line, which slows the high velocity (outer) element and speeds up the slower (inner) element, hence transferring angular momentum that forces the elements to larger and smaller radii respectively. This further stretches the field line, so leads to a runaway instability.

Alternative processes include the hydrodynamical instability (Godon & Livio, 2000; Klahr & Bodenheimer, 2003; Umurhan & Regev, 2004; Johnson & Gammie, 2005), whereby vortices in discs drive angular momentum transport. However, these vortices are
often short lived and do not lead to long term viscosities. Additionally, gravitational instabilities may become important when the masses of the disc and star are similar (Lodato & Rice, 2004, 2005). Here, thermal pressure and rotation are countered by self-gravity, which acts to destabilise the disc. Self gravity leads to clumping of disc material, and spiral density waves are formed that typically carry both energy and angular momentum outwards.

1.2.2.1. The Shakura-Sunyaev solution

In 1973, Shakura & Sunyaev proposed a simple parameterisation for a disc viscosity driven by turbulent motion. By reasoning that the largest turbulent scales are set by the geometry of the flow, they argued that the vertical scale height of the disc, \( H \), was a suitable representative scale. Furthermore, they assumed that turbulence must be subsonic, since if it were supersonic it would rapidly thermalise by shocks. Hence, they decided that the speed of sound, \( c_s \), must be the characteristic speed of turbulent motion. This reasoning led to the following formula for the kinematic viscosity term \( \nu \):

\[
\nu = \alpha c_s H,
\]

where \( \alpha \) is a dimensionless parameter, expected to be \( \lesssim 1 \). Although our uncertainty has essentially been shifted to a new parameter, \( \alpha \), this formalism is extremely powerful for predicting general disc behaviour. Typically quoted values of \( \alpha \) range from \( \sim 0.01 \) for protostellar discs (Hartmann et al., 1998) to \( \sim 0.1 \) for X-ray binaries (e.g. Lasota, 2001). A more detailed discussion of possible values for \( \alpha \) can be found in King et al. (2007).

1.2.3 The thermal-viscous instability

Although this steady thin disc scenario is a highly powerful and applicable approach to a variety of astrophysical phenomena, it is important to note that various stability criteria must be met for a steady state solution to exist. Of particular importance to this thesis is the thermal-viscous instability, under which a disc can evolve through hot and cold states as it tries to reach a stable configuration. This will be described, after a brief diversion to discuss some relevant timescales.

1.2.3.1. Important timescales

Alongside the viscous timescale defined in (1.40), there are two further timescales which are important for disc evolution. The first of these is the dynamical timescale, \( t_{\text{dyn}} \), which is the period of a Keplerian revolution and also the vertical sound crossing time of the disc. There is also the thermal timescale, \( t_{\text{th}} \), which is the ratio of the thermal content to the local dissipation rate. The three timescales are related to each other as follows:
where $D(R)$ is the viscous dissipation per unit area. This means that for $\alpha \ll 1$, there is a hierarchy of timescales where $t_{\text{dyn}} \ll t_{\text{th}} \ll t_{\text{visc}}$. In fact, using reasonable parameter choices (and further formulae for steady thin discs; see e.g. (5.49) of Frank et al., 2002),

\[
t_{\text{dyn}} \sim \alpha t_{\text{th}} \sim 100 m_1^{-1/2} R_{10}^{3/2} \text{s} \quad (1.45)
\]
\[
t_{\text{visc}} \sim 3 \times 10^5 \alpha^{-4/5} \frac{\dot{M}_{16}}{\Sigma}^{-3/10} m_1^{1/4} R_{10}^{5/4} \text{s}. \quad (1.46)
\]

where $R \text{ cm} = 10^{10} R_{10} \text{ cm}$ and $\dot{M} \text{ g s}^{-1} = 10^{16} \dot{M}_{16} \text{ g s}^{-1}$. For typical parameters, the dynamical and thermal timescales are of the order of minutes and the viscous timescale is of the order of days to weeks. These differences mean that we can distinguish between disc processes by measuring the timescale over which they occur. In particular they are associated with different types of instability. For example, disruption of the energy balance of a disc will grow over $t_{\text{th}}$. Since $\Sigma$ changes over the longer $t_{\text{visc}}$, $\Sigma$ will remain constant over this timescale while in contrast, the vertical structure, which changes over a dynamical timescale, would respond quickly and remain in hydrostatic equilibrium. This type of instability would be called a thermal instability.

**1.2.3.2. Conditions for a local thermal-viscous instability**

Linear perturbation analysis of (1.26) shows that a disc is locally viscously stable at distance $R$ provided:

\[
\frac{\partial [\nu\Sigma]}{\partial \Sigma} > 0, \quad (1.47)
\]

or equivalently,

\[
\frac{\partial \dot{M}}{\partial \Sigma} > 0, \quad (1.48)
\]

or

\[
\frac{\partial T}{\partial \Sigma} > 0. \quad (1.49)
\]
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The second form of the relation (1.48) is intuitively justifiable. If an annulus at radius $R_0$ is in the unstable regime, $\partial M/\partial \Sigma < 0$, then increasing its surface density would lead to a lower local mass accretion rate. However, the mass supply rates from neighbouring annuli would still have their original higher values. Consequently, mass would accumulate locally, increasing the surface density even more, again lowering the mass accretion rate; this is an indication of instability.

Since this is a viscous instability, one might expect that the annulus is in both dynamic and thermal equilibrium, as the timescales are significantly shorter. However, it is worth considering the consequences of these relations on the thermal stability.

In a steady annulus, the heating rate, $Q_+$ and the cooling rate, $Q_-$ must balance. The heating rate given by the viscous dissipation, $G\Omega'dR$ for a ring of width $dR$. Since this energy is radiated over the upper and lower faces of the disc, an area of $4\pi R dR$, the energy emitted per unit area in the annulus is

$$D(R) = \frac{G\Omega'}{4\pi R} = \frac{1}{2} \nu\Sigma(R\Omega')^2$$

where $G$ has been substituted from (1.22). $D(R)$ is therefore proportional to $\nu\Sigma$. With the Shakura-Sunyaev solution, $\nu = \alpha c_s H$, where $c_s \simeq 10\sqrt{(T/10^4K)}$ km s$^{-1}$, the dependence can alternatively be written as $D(R) \propto T^{1/2}\Sigma$.

The rate of energy lost from a disc annulus, via radiation, per unit disc area is given by

$$F(H) \approx \frac{4\sigma T^4}{3\kappa \rho H},$$

where the Kramer’s opacity, $\kappa = \rho T_c^{-3.5}$. Substituting this formalism in, we find that $F(H) \propto T^{8}\Sigma^{-2}$.

In equilibrium, $D = F$ (heating and cooling balance). If a viscous perturbation increases $\Sigma$ while the temperature remains constant, then the heating rate increases and the cooling rate drops, so that heating dominates over cooling and the temperature rises. However, if the disc is viscously unstable, then $\partial(\nu\Sigma)/\partial\Sigma < 0$. Therefore, if a perturbation $\partial\Sigma$ is greater than 0, $\Delta\nu < 0$, so a new steady solution would require a lower temperature because $\nu \propto T^{1/2}$. Therefore, the thermal equilibrium, where the annulus needs to heat up, is unstable.

The consequence of this discussion is that a viscous instability of this nature implies a thermal instability, causing the disc to evolve on a thermal timescale.

1.2.3.3. The “S” curve

For a ring of material at radius $R_0$, there are multiple solutions to thermal equilibrium. When plotted in the log $T$ (or equivalently, log $\dot{M}$), log $\Sigma$ plane, they form an S shaped curve (see, e.g. Meyer & Meyer-Hofmeister, 1981; Cannizzo et al., 1982; Faulkner et al.,.
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Figure 1.5 The S-curve and limit-cycle behaviour. In a plot of temperature against surface density, the S-shaped curve shows the equilibrium of heating and cooling of an annulus in an accretion disc. The negative gradient is unstable. Consequently, an annulus whose equilibrium position lies on this part of the curve will instead follow the cycle shown from A to B, to C, to D in the diagram. The switch from the cold (lower) positive branch to the hot (upper) positive branch is associated with the ionization of hydrogen at temperature \( T_H \sim 6500\, \text{K} \).

1983), as shown in Figure 1.5. Away from the curve, the temperature of the annulus changes on a thermal timescale, moving directly up or down on the diagram, to reach the nearest point of equilibrium. After this, it will evolve on a viscous timescale, moving along the S-curve until it reaches a stable accretion rate, which is set externally by the mass flowing in from neighbouring annuli. However, from (1.49), we know that the positive gradients on this plot are stable, whereas the negative gradient is unstable. If the equilibrium position for the annulus naturally lies somewhere along the negative gradient, then stability is unobtainable, because the slightest random perturbation sends the annulus away from equilibrium again. In this situation, limit-cycle behaviour begins, whereby the surface density and temperature evolve from point A, to B, to C, to D in Figure 1.5. Evolution along the positive gradients of the S-curve occur on a viscous timescale, as the annulus is in thermal equilibrium. However when \( \Sigma_{\text{max}} \) is reached, at point B, the disc becomes thermally unstable, and increases its temperature on a thermal timescale to reach point C on the higher level. The same happens in the opposite direction from D to A.

Physically, for Keplerian steady-state accretion discs, the unstable state is associated with the partial ionization of hydrogen, at \( T_H \sim 6500\, \text{K} \). (e.g. Mineshige & Osaki, 1983;
Pojmanski, 1986; Cannizzo & Mattei, 1992). Between $T \sim 6000 - 10000K$, there is an extremely steep dependence of opacity on the temperature, so that a small increase in temperature leads to a correspondingly large rise in opacity. Radiation from material in a thin disc is thermal (see Section 1.3.1), and photons in the Wien tail of the distribution can ionize some of the hydrogen in the disc. The reabsorption of photons into the disc raises its temperature instead of cooling it. Thus, more photons are able to ionize hydrogen in the disc, the temperature rises further, and a runaway rise in temperature begins. This ends only when the disc is fully ionized, at which point the opacity no longer rises so rapidly. On the upper branch, the gas is said to be in a "hot" or "high" state. Conversely, on the lower branch, hydrogen is predominantly neutral, and the disc is in a "cold" or "low" state. These expressions will be used interchangeably in the rest of this thesis.

$\Sigma_{\min}$ and $\Sigma_{\max}$ can be estimated using the full vertical structure disc equations, assuming an $\alpha$-prescription viscosity (Frank et al., 2002) as:

$$\Sigma_{\min} = 8.25 R_{10}^{1.05} M_1^{-0.35} \alpha_h^{-0.8} \text{gcm}^{-2}, \quad (1.52)$$
$$\Sigma_{\max} = 11.4 R_{10}^{1.05} M_1^{-0.35} \alpha_c^{-0.86} \text{gcm}^{-2}. \quad (1.53)$$

In these equations, $\alpha_h$ and $\alpha_c$ are the characteristic values for the hot and cold states respectively. Observations show that the $\alpha$ parameter is different depending whether the annulus is in a high or low state, so $\alpha_h \sim 0.1$ and $\alpha_c \sim 0.01$. Note that these different values of $\alpha$ result in a longer viscous time for a ring of material on the lower branch than on the upper branch. Therefore, the annulus will spend more time in the cold state.

1.2.3.4. Global thermal-viscous instabilities

So far this discussion has been limited to a single annulus. However, the onset of the thermal-viscous instability in one part of the disc can alter the mass inflow/outflow rate for other parts of the disc, triggering further instabilities. Analysis by e.g. Meyer (1984); Papaloizou & Pringle (1985); Lin et al. (1985); Cannizzo et al. (1995); Vishniac & Wheeler (1996); Vishniac (1997); Menou et al. (1999) shows that heating and cooling waves propagate through the disc, leading to a global instability. Since mass outflow rate increases significantly in the high state, the moving front means mass is accreted rapidly onto the central object, emptying the disc until the $\Sigma_{\min}$ is reached and the disc drops into a cool state again. The disc then cycles between two states: a hot and mostly ionised state with a large local accretion rate, and a cold, neutral state with a low accretion rate.

While this instability was originally proposed to account for large-amplitude luminosity variations observed in cataclysmic variables (Meyer & Meyer-Hofmeister, 1981; Smak, 1982), it is also believed that the same mechanism is responsible for high luminosity phases in soft X-ray transients (Cannizzo et al., 1995; Dubus et al., 2001; Lasota,
2001). However, X-ray binaries have extremely luminous inner discs that can irradiate the rest of the disc, keeping it ionized for extended periods. These outbursts will be discussed in further detail in Section 1.3.

1.2.4 Advection dominated accretion flows

Beyond the steady thin disc model, there are three further self-consistent solutions that can be derived from the hydrodynamic disc equations. In 1976, Shapiro, Lightman & Eardley developed a second, hotter solution, in which the accreting gas forms a two temperature plasma with the ion temperature greater than the electron temperature. However this was found to be unstable and unlikely to exist for real flows (Piran, 1978).

The other two solutions to the disc equations involve a process known as an Advection Dominated Accretion Flow (ADAF hereafter). Here, energy is trapped in the gas, and is carried inwards rather than being radiated. This solution is found first at super-Eddington accretion rates, where the large optical depth of the inflowing gas traps radiation (Katz, 1977; Abramowicz et al., 1988). The second case originates in the opposite limit of low, sub-Eddington accretion rates, for which the main formulism is derived in Narayan & Yi (1994, 1995a,b); Abramowicz et al. (1995); Chen et al. (1995), although some ideas were introduced previously by (Ichimaru, 1977). In this optically thin solution, the accreting gas has a very low density, so cannot cool effectively. A two temperature structure forms, where hot ions hold most of the energy but transfer it inefficiently to the electrons that produce most of the radiation. Radiation is therefore largely prevented from escaping from the gas, and the energy is instead advected inwards towards the central accretor. These low accretion rate ADAFs are thought to be present in many accreting systems, from X-ray binaries, to low luminosity active galactic nuclei (see e.g. Narayan & McClintock, 2008, for a review), and are particularly relevant to this thesis.

A synopsis of the mathematical treatment behind Advection Dominated Accretion Flows can be found in e.g. Narayan et al. (1998). However, in this work, only a broad understanding of the observational consequences of inefficient accretion is required.

An ADAF forms when the mass accretion flow drops below a critical value of

\[ \dot{M}_{\text{crit}} \sim f \dot{M}_{\text{Edd}} \]  

where \( f \) is expected to be between 0.01 and 0.1 (Narayan & Yi, 1994; Narayan, 1996; Esin et al., 1997). The defining property of an ADAF is to reduce the efficiency \( \eta \) at which matter is converted to radiation (see (1.13)). For efficient flows, \( \eta \approx 0.1 \). However, for an ADAF,

\[ \eta = \frac{L}{\dot{M}c^2} \ll 0.1. \]  

(1.55)
At $\dot{M}_{\text{crit}}$, continuity requires that $\eta \sim 0.1$. Below this value, Narayan & McClintock (2008) roughly estimate that

$$\eta \sim 0.1 \frac{\dot{M}}{fM_{\text{Edd}}}.$$  (1.56)

### 1.3 OBSERVING X-RAY BINARIES

Accretion discs around stellar mass black holes emit characteristic but widely variable spectra. While HMXBs are generally persistent sources, exhibiting spectral changes without large variations in magnitude, LMXBs undergo outbursts in which their X-ray luminosity dramatically rises and before tailing off in an exponential fashion. In this section, I will explain the main features of the spectra and light curves of these systems, and the theoretical background that underlies them. Following this I will provide a brief overview of the black holes we see in our own galaxy.

#### 1.3.1 Spectral shape

The X-ray spectra of black hole binaries are typically composed of thermal emission from the optically thick accretion disc, and some additional higher energy, non-thermal emission. However, the relative power of these two components varies considerably, leading to a wide diversity in observed spectra. Generally, spectra alternate between two very different “states”, over a range of timescales. These are shown for the HMXB Cygnus X-1 in Figure 1.6 (Gierliński et al., 1997). The spectra are plotted in $\nu F_\nu$, so that peaks represent the characteristic photon energy of the source output. Plotted in red is the “high” or “soft” state, which is dominated by thermal emission below $\sim 10\text{keV}$, with a non-thermal tail extending to $> 500\text{keV}$. The “low” or “hard” state, shown in blue, peaks at $\sim 100\text{keV}$ with the non-thermal component overpowering the thermal one.

For transient sources, “low” and “high” state spectra correspond to significantly reduced, and raised, X-ray luminosities respectively. However, in the case of HMXBs, this definition is largely historical because they are persistent sources. Variations in the luminosity of Cygnus X-1 were originally thought to be greater than a factor of 5, when observations only covered the 2-10keV range. However, over the full spectral range, the change is less dramatic (e.g. Nowak, 1995; Zhang et al., 1997; Gierliński et al., 1999), and Done & Gierliński (2003) have shown that the long-term luminosity of Cygnus X-1 only varies by a factor of 3.

I will describe the broad features of the two spectral states, and their origins, below. However, this will not be a comprehensive review of all the emission processes which contribute to the X-ray spectra, or the components which are included in detailed spectral
modelling. For a thorough evaluation of all the radiative processes involved in black hole X-ray binary emission, see e.g. Zdziarski & Gierliński (2004).

In many cases, spectral components are modelled as power-law spectra, where the flux density follows

\[ F(E) \propto E^{-\alpha} \]  

and \( \alpha \) is the spectral index. This can also be written as a photon flux density,

\[ N(E) \propto E^{-\Gamma}, \]  

where \( \Gamma = \alpha + 1 \) is the photon index.

### 1.3.1.1. Soft state spectra

The soft state spectrum is dominated by emission from the thin accretion disc, which is optically thick. Radiation is therefore emitted as a blackbody spectrum described by the Planck function,

\[ B_\nu(T) = \frac{2h\nu^3}{c^2} \frac{1}{\exp\left(\frac{\nu}{kT}\right) - 1} \text{ erg cm}^{-2} \text{ Hz}^{-1} \text{ ster}^{-1}, \]  

where \( h \) is the Planck constant and \( k \) is the Boltzmann constant. Because the temperature of the disc decreases with distance from the source, a combination of different blackbody spectra emitted from each disc annulus combine to form what is commonly referred to as a multicolour disc blackbody spectrum (Dotani et al., 1997). The disc luminosity can be
CHAPTER 1. INTRODUCTION

written as \( L_{\text{disc}} = 4\pi r_{\text{in}}^2 \sigma T_{\text{in}}^4 \), where \( r_{\text{in}} \) and \( T_{\text{in}} \) are the inner disc radius and temperature respectively.

The high energy tail of the spectrum is non-thermal, and can be modelled as a power law, with \( \Gamma \sim 2 - 2.2 \), extending beyond 500 keV (Gierliński et al., 1999). This component is well fitted by Compton scattering on a non-thermal electron population.

1.3.1.2. Hard state spectra

In the hard state, the disc component of the spectrum is no longer dominant, and is rarely detected (exceptions include the black hole binary J1118+480; Esin et al., 2001). Instead, the spectrum follows an approximate power law with \( \Gamma \sim 1.5 - 2 \). This state corresponds to lower mass accretion rates where the reduced density of the flow near the source leads to fewer collisions and incomplete thermalisation. Electrons gain only a small amount of energy via Coulomb collisions, which they lose by radiation. The spectrum is broadly fitted by a thermal Comptonization model where, for Cygnus X-1, \( kT_e = 75 - 110 \) keV (Ibragimov et al., 2005).

Geometrically, a thin disc model, though largely consistent with the soft state spectra, is no longer applicable. Instead, accretion in the low hard state is more analogous to a quasi-spherical, ADAF-like flow. This has led to a standard theory arguing that discs in the hard state are truncated far from the hole itself. A hot flow closer to the source is then responsible for main spectral features.

1.3.1.3. Spectral evolution

The spectral evolution of three X-ray binaries (GX 339-4, XTE J1859+226 and XTE J1550-564) is shown in Figure 1.7, reproduced from Figure 3 of Fender et al. (2004). Here, their Eddington-ratioed luminosities have been plotted against X-ray colour, in a hardness-intensity diagram. X-ray binaries follow anti-clockwise tracks, beginning in a low state (LS on diagram), which is associated with a near-vertical rise in luminosity and a hard X-ray spectrum. The spectrum then becomes progressively soft as it moves from right-to-left in the diagram, towards the high state (HS). It then decreases in luminosity, hardens, and returns to the low state.

The transition between soft and hard state spectra implies a dramatic change in the structure of the disc, which is currently understood as follows. In the soft state, the optically thick disc extends close to the black hole, and the disc blackbody dominates the spectrum. However, in periods of low \( L/L_{\text{Edd}} \), the inner disc is replaced by an optically thin hot flow, most likely through evaporation (Meyer & Meyer-Hofmeister, 1994; Różańska & Czerny, 2000; Mayer & Pringle, 2007). The emitted spectrum is then dominated by the thermal Compton scattering of the electron population in the hot inner flow. At later times, as the mass inflow rate increases again, the disc extends progressively
Figure 1.7 Combined X-ray hardness-intensity diagram for GX 339-4, XTE J1859+226 and XTE J1550-564. In this figure, Eddington-ratioed luminosities have been plotted against X-ray colour to chart the spectral evolution of three X-ray binaries. Binaries in their low state (LS) with hard spectra progress vertically up the diagram as their luminosity increases, before crossing horizontally to the high state (HS) passing through an intermediate branch. In this state the X-ray spectrum is much softer, and the luminosity is high. The spectra then progress anti-clockwise around the diagram, eventually returning to the hard state. This figure is reproduced from Fender et al. (2004).

Further underneath the flow, and the blackbody component again becomes dominant.

The transition between the low and hard states is given theoretically as $L \sim 0.01L_{\text{Edd}}$ under ADAF assumptions (Narayan & Yi, 1995b), or alternatively $L \sim 0.1L_{\text{Edd}}$ under the relaxed constraints of Yuan (2001). Observationally, the switch is found to occur between 0.2 and 0.003$L_{\text{Edd}}$ for the best studied case of XTE J1550-564 (Done & Gierliński, 2003). This can be further constrained to $L/L_{\text{Edd}} = 0.02$, for just the soft-to-hard state transition (Maccarone & Coppi, 2003) but there are often deviations from this value (Done & Gierliński, 2003).

This range of transition luminosities is a manifestation of a property known as hysteresis, where the same source can show different spectra at the same luminosity. A consequence of this property is that there is no single relation between spectral state and $L/L_{\text{Edd}}$ for the majority of black hole X-ray binaries (Narayan & Yi, 1995b; Maccarone & Coppi, 2003)

Although widely accepted, the disc truncation, hot flow, model for hard state spectra is still regularly contested. Several alternative mechanisms have been put forward (see e.g. Done et al., 2007, for a review) but these generally require that the X-ray source be
mildly beamed away from the disc, in conflict with observations that show no trend in X-ray binary properties as a function of inclination (Fender & Belloni, 2004; Narayan & McClintock, 2005). Furthermore, various studies have argued that there is observational evidence for an untruncated disc. For example, both Miller et al. (2002) and Miniutti et al. (2004) show that extremely smeared iron lines and reflection features in XTE J1650-500 can be interpreted as evidence that the disc extends inwards towards the hole. However, Done & Gierliński (2006) were able to refit the model to show that it is in line with current theory.

Additionally it should be noted that Cygnus X-1, among other systems, actually exhibits a broader range of states than addressed here. These can all be accounted for within the truncated disc model. For a comprehensive review of these states, see e.g. Tanaka & Lewin (1995) or Remillard & McClintock (2006).

1.3.1.4. Radio emission and jets

Radio emission from X-ray binaries is strongly correlated to their spectral state. The low state is associated with a steady jet, and a non-linear radio and X-ray luminosity correlation of $L_{\text{Radio}} \propto L_X^{0.7}$ (Corbel et al., 2003; Gallo et al., 2003). This close correlation is due to the ADAF-like inner flow in the the disc acting as a base to the jet (Heinz & Sunyaev, 2003). In the high state, however, radio emission is strongly suppressed (Tananbaum et al., 1972; Fender et al., 1999; Corbel et al., 2001). The transition between the low and high states corresponds to a transition from strong to suppressed radio emission, during which transient radio outbursts are often observed (e.g. Mirabel & Rodríguez, 1994; Hjellming & Rupen, 1995). The disc-jet relationship through progressive states is explained in detail for G1915+150 in Fender & Belloni (2004) and as a unified model, in Fender et al. (2004)

1.3.1.5. Neutron star X-ray binary spectra

There is fundamental difference between black hole and neutron star X-ray spectra because neutron stars have a solid surface. In an accretion flow, half of the energy is stored as kinetic energy. This is radiated at the surface in a boundary layer for neutron stars, leading to an additional spectral component. Spectra generally comprise soft and hard components from the disc and Comptonized boundary layer respectively, which both contribute in varying extents to the X-ray luminosity, depending on the spectral state. For a discussion of the main features of neutron star spectra, see for example, Gilfanov et al. (2003), Revnivtsev & Gilfanov (2006) or Section 7 of Done et al. (2007)

1.3.2 Light curves

All black hole, and many neutron star, LMXBs are subject to the thermal-viscous instability. Consequently, they undergo outbursts. These are phases of high mass accretion
Figure 1.8 The visual light curve of SS Cygni. In this figure, the magnitude of SS Cygni, the brightest dwarf nova, is plotted over 200 days. The source exhibits four complete outbursts during this period. This figure was created using the AAVSO light curve generator at http://www.aavso.org/adata/curvegenerator.shtml.

rate, and hence luminosity, where part or all of the disc becomes ionized. Outbursts were originally identified in the dwarf novae subclass of accreting white dwarf binary systems (see Cannizzo, 1993, for a review), and a characteristic light curve for such a system is shown in Figure 1.8. Here, the visual luminosity of SS Cygni, the brightest dwarf nova, is plotted against time for two hundred days during which it experiences four outbursts, with rise times of about a day, which last for around a week. In between outbursts, there are weeks to months of quiescence.

LMXB outbursts are considerably different to those of dwarf novae, having much longer timescales of months to years. A typical outburst light curve is shown on the left panel of Figure 1.9 for the black hole LMXB A 0620-00. Unlike the sharp decline seen in Figure 1.8, the outburst follows a quasi-exponential decay (see Tanaka & Lewin, 1995; Tanaka & Shibazaki, 1996, for reviews). A second X-ray maximum is also visible after which the exponential decline continues. Finally, while just 3-5% of the disc mass is accreted by the white dwarf in a dwarf nova outburst (Cannizzo, 1993), LMXB outbursts result in a much larger fraction of the disk mass being accreted, which may approach 100% in the case of small disks.

The difference between the outbursts of X-ray binaries and dwarf novae is due to irradiation of the accretion disc. Irradiation keeps the disc hot, even at large radii (van Paradijs, 1996), extending the time the disc spends on the hot branch, and greatly altering the progression of the outburst.
1.3.2.1. Disc irradiation

If the temperature of a disc is kept above the hydrogen ionization temperature, $T_H$, then it will always be ionized, and will not undergo the thermal-viscous instability. Hence irradiation, such that $T_{\text{irr}} > T_H$ at all radii, yields a stable disc (Tuchman et al., 1990). For this reason, HMXBs, with more massive secondaries, higher mass transfer rates and, consequently, higher disc temperatures, are usually persistent (van Paradijs, 1996). Conversely, LMXBs, with generally lower accretion rates, may be transient. All known black hole LMXBs are transient (King et al., 1996; Dubus et al., 1999), but neutron star systems may be persistent, because they typically have smaller discs, due to tidal truncation, so the outer disc is often hotter.

The surface temperature of a concave disc face heated by a compact central X-ray source was calculated by van Paradijs (1996) as:

$$T_{\text{irr}}(R) = \frac{\eta M_c c^2 (1 - \beta)}{4 \pi R^2 \sigma} \left( \frac{H}{R} \right)^n \left[ \frac{d \ln H}{d \ln R} - 1 \right],$$  

for a central mass accretion rate of $\dot{M}_c$ with accretion efficiency of $\eta$, c.f. (1.13). $\beta$ is the albedo of the disc faces, and lies between 5/7 and 7/8. $H$ is the disc scale height, and the index $n = 1$ or 2, for neutron star and black hole sources respectively. For $n = 1$ the central accretor is regarded as a point source. However, for black holes, without a solid surface, the source of irradiation is actually the inner region of the disc (Fukue, 1992; King et al., 1997), which leads to an extra factor of $H/R$ in (1.60).

From (1.60), the requirement of $T_{\text{irr}} > T_H$ for disc stability can alternatively be written as $\dot{M}_c > \dot{M}_{\text{crit}}$, i.e. the central mass accretion rate must be greater than some critical value, for the entire disc to be irradiated. In the standard disc instability model, described in Section 1.2.3, outbursts end when a cooling wave moves inwards from the outer edge of the disc. However, if $\dot{M}_c > \dot{M}_{\text{crit}}$ at the outer edge of the disc, such a cooling wave will be prevented. Instead, the disc must remain in the hot state until $\dot{M}_c$ drops below $\dot{M}_{\text{crit}}$. Since the surface density of a disc in the hot state is quasi-steady state, $\dot{M}_c$ only drops as the mass of the disc decreases. For this reason, LMXB outbursts consume most of the disc, and the outburst decays over a much longer timescale.

1.3.2.2. Outbursts of irradiated discs in X-rays

King & Ritter (1998) developed an analytic overview of an LMXB outburst which accounts for disc irradiation. They treat the surface-density profile of the disc in the hot state as quasi-steady state, obeying

$$\Sigma(R) \sim \frac{\dot{M}_c}{3 \pi \nu},$$

which is obtained from (1.39) with $R \gg R_s$, and $\nu \sim$ constant (a significant approxima-
Accretion flows in XRB

Figure 1.9 Light curves for an irradiated disc outburst; observed and theoretical. On the left, the X-ray light curve of BHB A 0620-00 is plotted in outburst, using data from the Ariel V all sky monitor. For comparison the theoretical outburst lightcurve of an irradiated disc from King & Ritter (1998) is plotted on the right. This figure is reproduced from Figure 2 of Done et al. (2007) with kind permission from Springer Science and Business Media.

...tion). If we assume the disc is heated to some radius, $R_h$, then integrating $\Sigma(R)$ gives the mass of the heated zone:

$$M_h = 2\pi \int_0^{R_h} \Sigma R dR \approx M_c \frac{R_h^2}{3\nu}.$$  \hfill (1.62)

$M_h$ can only change through central accretion, so $\dot{M}_c = -\dot{M}_h$. Substituting this relation into (1.62) gives

$$-\dot{M}_h = \frac{3\nu}{R_h^2} M_h,$$  \hfill (1.63)

which can be integrated to show that $M_h$ decays exponentially with time

$$M_h = M_0 \exp(-3\nu t / R_h^2),$$  \hfill (1.64)

where $M_0$ is the initial mass of the heated zone. King & Ritter (1998) integrate the maximum surface density, $\Sigma_{\text{max}} \approx \rho R/2\pi$, out to $R_h$, to estimate $M_0 \approx \rho R_h^3 / 3$. Inserting this into (1.64), gives

$$M_h = \frac{\rho R_h^3}{3} \exp(-3\nu t / R_h^2),$$  \hfill (1.65)

which can be written in terms of $\dot{M}_c$ as

$$\dot{M}_c = -\dot{M}_h = R_h \nu \rho \exp(-3\nu t / R_h^2).$$  \hfill (1.66)
Therefore, the first part of the X-ray light curve follows the observed exponential profile. After the onset of an outburst, the outer disc becomes irradiated, increasing $R_h$ to some radius $R_0$. The normalisation then increases in (1.66) leading to a jump in the X-ray luminosity, which accounts for the second X-ray maximum observed in many XRT lightcurves at approximately 50-75 days after that start of the outburst. Note that $R_0$ will be the disc radius for short orbital period systems, since $R_D$ increases with $P_{\text{orb}}$.

X-rays decay exponentially until the outer edge of the disc cannot be kept in the hot state. Rearranging (1.60), $R_h$ is given by

$$R_h^2 = \frac{\eta \dot{M}_c c^2 (1 - \beta)}{4\pi \sigma T_H^4} \left( \frac{H}{R} \right)^n \left[ \frac{d \ln H}{d \ln R} - 1 \right] = B_n \dot{M}_c,$$  \hspace{1cm} (1.67)

where $B_1 \sim 4 \times 10^5$, and $B_2 \sim 5 \times 10^4$(cgs). The condition that $R_h < R_0$ is therefore, $\dot{M}_c < \dot{M}_{\text{crit}}$. King and Ritter use typical values of $\eta = 0.15$, $d \ln H/d \ln R = 45/38$ and $\eta = 0.10$, $d \ln H/d \ln R = 43/36$ for neutron star and black hole systems respectively to estimate:

$$\dot{M}_{\text{crit}}(\text{NS}) = 4.1 \times 10^{-10} R_{11}^2 M_\odot \text{yr}^{-1},$$  \hspace{1cm} (1.68)

$$\dot{M}_{\text{crit}}(\text{BH}) = 2.9 \times 10^{-9} R_{11}^2 M_\odot \text{yr}^{-1},$$  \hspace{1cm} (1.69)

where $R_0 = R_{11} \times 10^{11}$cm, $\beta = 0.8$ and $H/R = 0.2$ in both cases (de Jong et al., 1996; King et al., 1997). Note that if $\dot{M}_c(t = 0) > \dot{M}_{\text{crit}}$, sources may never experience a phase of exponential decay.

Inserting the condition of (1.67) into (1.62) shows that the exponential decay ends when

$$\dot{M}_c = \left( \frac{3\nu}{B_n} \right)^{1/2} \dot{M}_h^{1/2}.$$  \hspace{1cm} (1.70)

After this condition is met, the heated disc mass changes, not only via central accretion, but also due to the decrease in the outer heated radius $R_h$ with $\dot{M}_c$. Hence,

$$\dot{M}_h = -\dot{M}_c + 2\pi \Sigma R_h \dot{R}_h.$$  \hspace{1cm} (1.71)

Using (1.62) and (1.67), this can be written as

$$\dot{M}_h = -\dot{M}_c + \frac{B_n \dot{M}_c}{3\nu} \dot{M}_c,$$  \hspace{1cm} (1.72)
which can be further simplified by substituting (1.70) to show

\[
\frac{B_n \dot{M}_c}{3\nu} = -1. \quad (1.73)
\]

Substituting this relation back into (1.71) gives

\[
\dot{M}_h = -2\dot{M}_c \quad (1.74)
\]

so, with (1.70),

\[
\dot{M}_h = -2\left(\frac{3\nu}{B_n}\right)^{1/2} M_h^{1/2}. \quad (1.75)
\]

Integrating, gives

\[
M_h = \left[ M_h^{1/2}(T) - \left(\frac{3\nu}{B_n}\right)^{1/2} (t - T) \right]^2, \quad (1.76)
\]

finally yielding

\[
\dot{M}_c = \left(\frac{3\nu}{B_n}\right)^{1/2} \left[ M_h^{1/2}(T) - \left(\frac{3\nu}{B_n}\right)^{1/2} (t - T) \right], \quad (1.77)
\]

i.e. a phase of linear decay, beginning at time T. Here, \(M_h(T)\) is the hot state disc mass at T, with T = 0 for systems with long orbital periods. The outburst ends when \(\dot{M}_c = 0\) after a further time

\[
t_{\text{end}} - T = \left(\frac{B_n}{3\nu}\right)^{1/2} M_h^{1/2}(T). \quad (1.78)
\]

Of course, the source may be very faint in X-rays before the end of outburst, so may no longer be detectable before quiescence is reached.

The outburst profile predicted by this model – an exponential decay, including secondary X-ray maximum, and a linear tail-off – is shown in the right hand panel of Figure 1.9. The observed light curve of A 0620-00 in the left-hand panel is well matched by these theoretical calculations.

On reaching quiescence, almost all of the originally heated disc mass \(M_0\) is accreted. Therefore, the mass needs to build up again before another outburst and, on average, we expect

\[
t_{\text{rec}} = \frac{M_0}{-\dot{M}_2}. \quad (1.79)
\]
However, observations show widely variable recurrence times for different sources.

### 1.3.3 Galactic population

There are over 300 X-ray binaries in our galaxy (McClintock & Remillard, 2006), approximately a third of which are transient. Of these, there are at least 17 confirmed black hole LMXBs (Ritter & Kolb, 2003), all of which are transient, and 3 black hole HMXBs, which are persistent sources. Additionally there are many candidate systems, and new transients are regularly being identified, such as MAXI J1659-152 (Kuulkers et al., 2013).

The difference in lifetimes between HMXBs and LMXBs mean they are typically found in very different environments. HMXBs, with lifetimes of just $10^{-5} - 10^{-7}$ years are generally observed in young stellar populations and star forming regions, while LMXBs, with much longer lifetimes of $10^7 - 10^9$ years are more likely to be found in older stellar populations, including globular clusters.

### 1.4 X-RAYS AND THEIR SURROUNDINGS

Accreting black holes can interact with their surroundings though radiation, jets, and other outflows. These effects are highly relevant to the study of Active Galactic Nuclei (accreting supermassive black holes at the centre of galaxies; AGN hereafter) and their surroundings (see Fabian, 2012, for a review). The tight correlation between the black hole mass and many physical properties of its host galaxy, has led researchers to conclude that a feedback process is occurring (Alexander & Hickox, 2012), a popular model for which is wind feedback (King, 2003; Pounds et al., 2003; King & Pounds, 2003). In low mass galaxies, feedback is instead likely to be dominated by massive stars and supernovae (e.g. Dekel & Silk, 1986; Ceverino & Klypin, 2009).

X-ray binaries can influence their surroundings both radiatively and kinetically. Locally, the energy deposited by ULXs is probably comparable to their observed X-ray luminosities (Pakull & Mirioni, 2003). Furthermore, radiation from X-rays may be an important source of feedback at high redshifts (Cantalupo, 2010; Justham & Schawinski, 2012), as discussed in this thesis. The kinetic influence of X-ray binaries on their surroundings is not the focus of this work, but may be equally significant. For example, ULXs such as SS 433 are known to drive outflows (see e.g. Blundell et al., 2001; Goodall et al., 2011), and the kinetic outflow of Cygnus X-1 is comparable to its X-ray output (Gallo et al., 2005).

The propagation and interaction of radiation with its environment is described by Radiative Transfer theory, which tracks the specific intensity:

$$I_{\nu} = \frac{dE}{dA dt d\nu d\Omega}.$$ (1.80)
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which is defined as the energy dE passing through a surface area dA over time dt and with frequency dv through solid angle dΩ. This is often averaged over a sphere as:

$$J_v = \frac{1}{4\pi} \int I_v d\Omega.$$  (1.81)

The transfer of radiation through its surroundings depends on various physical processes, which alter the temperature and ionization state of the traversed medium.

1.4.1 Physical interaction processes

The main physical interactions between photons and particles are summarised below. In Section 1.4.2, the radiative transfer code used in this thesis (developed by Bolton et al., 2004; Bolton & Haehnelt, 2007) will be introduced. In preparation for this, the relevant formalisms for each process will be included in their description. These will be assessed for both hydrogen and helium, the elements included in the model.

1.4.1.1. Photoionization: $X^0 + \gamma \rightarrow X^+ + e^-$

Photoionization occurs when an energetic photon is absorbed by an atom or ion, which releases an electron and increases its ionization state. The photon must have a minimum threshold ionization energy of 13.6eV, 24.6eV or 54.4eV for the ionization of H$^0$, He$^0$ or He$^+$ respectively. The cross-section to photoionization events is estimated in Osterbrock & Ferland (2006):

$$\sigma_{H0} = 6.30 \times 10^{-18} [1.34(\nu/\nu_{H0})^{-2.99} - 0.34(\nu/\nu_{H0})^{-3.99}] \text{cm}^2, \quad \sigma_{He0} = 7.03 \times 10^{-18} [1.66(\nu/\nu_{He0})^{-2.05} - 0.66(\nu/\nu_{He0})^{-3.05}] \text{cm}^2, \quad \sigma_{He^+} = 1.50 \times 10^{-18} [1.34(\nu/\nu_{He^+})^{-2.99} - 0.34(\nu/\nu_{He^+})^{-3.99}] \text{cm}^2.$$  (1.82)

Bolton & Haehnelt (2007) calculate the photoionization rate per unit volume [s$^{-1}$cm$^{-3}$] of elements of size $\delta R$ in a discretised grid along the line of sight from a source at timestep $\delta t$ as:

$$n_i^l \Gamma_i^l = \left[ \frac{1}{4\pi R^2 \delta R} \sum_k \frac{L_{vk}}{h_{p\nu_k}} P_i e^{-T_{vk}} \right] + n_i^l \Gamma_i^b,$$  (1.83)

for the species $i$ in grid element $l$. Here, $\nu_i$ is the ionization threshold frequency for species $i$, $P_i$ is the absorption probability (see Section 1.4.2) and $T_{\nu}$ is the transmission factor:

$$T_{\nu} = \begin{cases} 0 & (l = 0), \\ \sum_{j=0}^{l-1} (\sigma_{H0} n_{H0}^j + \sigma_{He0} n_{He0}^j + \sigma_{He^+} n_{He^+}^j) \delta R & (l > 0). \end{cases}$$  (1.84)

In the final term of (1.83), $\Gamma_i^b$ is the optically thin photoionization rate due to the diffuse...
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metagalactic UV background and is calculated in Bolton & Haehnelt (2007) as:

\[
\Gamma_{H^0}^b = 1.27 \times 10^{-11} j_{-21}^b (\alpha_b + 3)^{-1},
\]

\[
\Gamma_{He^0}^b = 1.51 \times 10^{-11} (0.553) \alpha_b j_{-21}^b (\alpha_b + 2)^{-1},
\]

\[
\Gamma_{He^+}^b = 3.03 \times 10^{-12} (0.250) \alpha_b j_{-21}^b (\alpha_b + 3)^{-1}
\]

with units \([s^{-1}]\), where the specific intensity of the metagalactic radiation field \([\text{ergs} \cdot \text{s}^{-1} \cdot \text{cm}^{-2} \cdot \text{sr}^{-1} \cdot \text{Hz}^{-1}]\) is assumed to have a power-law spectrum with index \(\alpha_b\) normalized by \(J_{-21}^b\) at the \(H^0\) ionization threshold,

\[
j_{\nu}^b = j_{-21}^b \times 10^{-21} \left( \frac{\nu}{\nu_{H0}} \right)^{-\alpha_b}.
\]  

1.4.1.2. Photoheating

Excess energy after photoionization becomes kinetic energy that heats the gas. The average heating rate per unit volume \([\text{ergs} \cdot \text{s}^{-1} \cdot \text{cm}^{-3}]\) for species \(i\) at grid element \(l\) is given by

\[
n_l^i \epsilon_{i}^l = \left[ \frac{1}{4\pi R^2 \delta R} \sum_k \frac{L_{vk} h_p (v_k - v_i) e^{-T_{vk}}}{\nu_k} \right] + n_l^i \epsilon_{i}^b,
\]

so that the total heating rate per unit volume at \(l\) is \(H_l^{\text{tot}} = n_{H^0} \epsilon_{H^0}^l + n_{He^0} \epsilon_{He^0}^l + n_{He^+} \epsilon_{He^+}^l\).

The background heating rates are given, by Bolton & Haehnelt (2007), in \([\text{ergs} \cdot \text{s}^{-1}]\) as:

\[
\epsilon_{H^0}^b = 2.94 \times 10^{-22} j_{-21}^b [(\alpha_b + 2)(\alpha_b + 3)]^{-1},
\]

\[
\epsilon_{He^0}^b = 6.48 \times 10^{-22} (0.553) \alpha_b j_{-21}^b [(\alpha_b + 1)(\alpha_b + 2)]^{-1},
\]

\[
\epsilon_{He^+}^b = 2.80 \times 10^{-22} (0.250) \alpha_b j_{-21}^b [(\alpha_b + 2)(\alpha_b + 3)]^{-1}.
\]

1.4.1.3. Collisional recombination: \(X^+ + e^- \rightarrow X^0 + \gamma\)

The reverse process of photoionization is collisional recombination, where electrons collide with ions, combining and releasing a photon. Recombination can be separated into two cases. In Case A recombination, the resulting atom can be in any energy state. However when an atom goes straight to its ground state, the released photon must have high enough energy to ionize another atom so that there is no net recombination. Case B recombination excludes transitions to the ground state for this reason. The latter is used in
this thesis, where recombinations are estimated in \([\text{cm}^3 \text{s}^{-1}]\) (Abel et al., 1997):

\[
\alpha_{\text{H}^+} = \exp(-28.6130338 - 0.72411256 \ln(\bar{T}) - 2.02604473 \times 10^{-2} \ln(\bar{T})^2 - 2.38086188 \times 10^{-3} \ln(\bar{T})^3 \\
- 3.21260521 \times 10^{-4} \ln(\bar{T})^4 - 1.42150291 \times 10^{-5} \ln(\bar{T})^5 \\
+ 4.98910892 \times 10^{-6} \ln(\bar{T})^6 + 5.75561414 \times 10^{-7} \ln(\bar{T})^7 \\
- 1.85676704 \times 10^{-8} \ln(\bar{T})^8 - 3.07113524 \times 10^{-9} \ln(\bar{T})^9),
\]

where \(\bar{T}\) is the gas temperature in eV.

1.4.1.4. Recombination cooling

When recombinations occur, the number of gas particles decreases and some of the kinetic energy is transferred to the photon, resulting in net cooling of the gas. The recombination cooling rates are given in \([\text{cm}^3 \text{s}^{-1}]\) in Bolton & Haehnelt (2007) as:

\[
\Lambda_{\text{H}^+}^{\text{rec}} = 1.036 \times 10^{-16} T \alpha_{\text{H}^+} n_e n_{\text{H}^+},
\]

\[
\Lambda_{\text{He}^+}^{\text{rec}} = (1.036 \times 10^{-16} T \alpha_{\text{He}^+} + 6.526 \times 10^{-11} \alpha_{\text{He}^+}^d)n_e n_{\text{He}^+},
\]

\[
\Lambda_{\text{He}^2+}^{\text{rec}} = 1.036 \times 10^{-16} T \alpha_{\text{He}^2+} n_e n_{\text{He}^2+}.
\]

1.4.1.5. Collisional ionization: \(X + e^- \rightarrow X^+ + 2e^-\)

In collisional ionization, the impact of a free electron with an atom results in the ejection of a bound electron, taking energy from the free electron. Collisional ionization rates in \([\text{cm}^3 \text{s}^{-1}]\) from Theuns et al. (1998) are given by:

\[
\Gamma_{\text{e}0} = 1.17 \times 10^{-10} T^{0.5} e^{-157809.4/T[1 + T_5^{1/2}]}^{-1},
\]

\[
\Gamma_{\text{eHe}^0} = 4.76 \times 10^{-11} T^{0.5} e^{-285335.4/T[1 + T_5^{1/2}]}^{-1},
\]

\[
\Gamma_{\text{eHe}^+} = 1.14 \times 10^{-11} T^{0.5} e^{-631515.0/T[1 + T_5^{1/2}]}^{-1},
\]

where \(T_5 = T/10^5\) K.

1.4.1.6. Collisional ionization cooling

The energy for collisional ionization comes from the kinetic energy of the reactants and therefore results in a reduction of the gas temperature. The collisional ionization cooling
rate is calculated in [erg s\(^{-1}\) cm\(^{-3}\)] as (Theuns et al., 1998):

\[
\Lambda_{eH_0} = 2.18 \times 10^{-11} \Gamma_{eH_0} n_e n_{H_0},
\]

\[
\Lambda_{eHe_0} = 3.94 \times 10^{-11} \Gamma_{eHe_0} n_e n_{He_0},
\]

\[
\Lambda_{eHe^+} = 8.72 \times 10^{-11} \Gamma_{eHe^+} n_e n_{He^+}.
\]

### 1.4.1.7. Collisional excitation cooling

In this process, a free electron colliding with an atom excites a bound electron into a higher state. It then decays, emitting a photon which carries energy away, making the gas cooler. Collisional excitation cooling [erg s\(^{-1}\) cm\(^{-3}\)] can be estimated for different species as (Cen, 1992):

\[
\Lambda_{eH_0}^\text{ex} = 7.50 \times 10^{-19} e^{-118348/T} [1 + T_5^{1/2}]^{-1} n_e n_{H_0},
\]

\[
\Lambda_{eHe_0}^\text{ex} = 9.10 \times 10^{-27} T^{-0.1687} e^{-13179.0/T} [1 + T_5^{1/2}]^{-1} n_e n_{He_0},
\]

\[
\Lambda_{eHe^+}^\text{ex} = 5.54 \times 10^{-17} T^{-0.397} e^{-473638/T} [1 + T_5^{1/2}]^{-1} n_e n_{He^+},
\]

### 1.4.1.8. Bremsstrahlung cooling

In Bremsstrahlung cooling, also known as free-free emission, radiation is produced through the deceleration of free electrons by ionized atoms. Kinetic energy is lost from the electrons, resulting in net cooling of the gas. The Bremsstrahlung cooling rate is given in [ergs\(^{-1}\)cm\(^{-3}\)] by Cen (1992):

\[
\Lambda_{ff} = 1.43 \times 10^{-27} T^{1/2} g_{ff} (n_{H^+} + n_{He^+} + n_{He^{2+}}) n_e,
\]

where the Gaunt factor for free-free emission, \(g_{ff}\) is given by

\[
g_{ff} = 1.1 + 0.34 \exp [- (5.5 - \log_{10}(T))]^2.
\]

### 1.4.1.9. Inverse compton cooling

Inverse Compton cooling occurs when low energy photons are scattered to higher energies by fast-moving electrons. This is different to the standard Compton effect because it is electrons that lose energy, rather than photons. The loss of energy in the electrons cools the gas at a rate [erg s\(^{-1}\) cm\(^{-3}\)] of (Peebles, 1993):

\[
\Lambda_c = 5.65 \times 10^{-36} [T - 2.73(1 + z)](1 + z)^4 n_e.
\]
1.4.2 A one-dimensional radiative transfer code

To study the ionization and temperature of the high-redshift IGM, Bolton et al. (2004) developed a one-dimensional radiative transfer model, which I apply to my research in Chapter 3. The model builds on the photon conservation algorithm developed by Abel et al. (1999), and has the advantage that energy conservation is nearly independent of spatial resolution, lowering the computational expense of the model.

The code estimates the propagation of radiation through a medium using equations for the probability of photons being absorbed by H, He and He$^+$:

\[
P_{\text{H}0} = p_{\text{H}0}q_{\text{He}0}q_{\text{He}^+} \left[ \frac{1 - \exp(-\tau^\text{tot}_V)}{D} \right],
\]

(1.97)

\[
P_{\text{He}0} = q_{\text{H}0}p_{\text{He}0}q_{\text{He}^+} \left[ \frac{1 - \exp(-\tau^\text{tot}_V)}{D} \right],
\]

(1.98)

\[
P_{\text{He}^+} = q_{\text{H}0}q_{\text{He}0}p_{\text{He}^+} \left[ \frac{1 - \exp(-\tau^\text{tot}_V)}{D} \right],
\]

(1.99)

where \( p_i = 1 - \exp(-\tau_i^V) \), \( q_i = \exp(-\tau_i^V) \), \( \tau^\text{tot}_V = \tau^\text{H}_V + \tau^{\text{He}0}_V + \tau^{\text{He}^+_V} \), \( \tau_i^V \) is the optical depth for species \( i \), and \( D = p_{\text{H}0}q_{\text{He}0}q_{\text{He}^+} + q_{\text{H}0}p_{\text{He}0}q_{\text{He}^+} + q_{\text{H}0}q_{\text{He}0}p_{\text{He}^+} \). These absorption probabilities are then used to calculate photoionization and photoheating rates in (1.83) and (1.88) respectively.

The model consists of a discretised grid spaced uniformly in logarithm, through which radiation is assumed to propagate instantaneously. It is the choice of \( c = \infty \) that conserves photon numbers. The removal of limitations that prevent light from propagating instantly through a cell means that spatial resolution does not need to be prohibitively high for accuracy.

The ionization state and temperature of atomic hydrogen and helium are computed at each time-step by solving a set of four coupled first-order differential equations. The first three determine the abundances of the ionized species:

\[
\frac{dn_{\text{H}^+}}{dt} = (\Gamma_{\text{H}0} + n_e \Gamma_{\text{eH}0})n_{\text{H}0} - n_{\text{H}^+}n_e \alpha_{\text{H}^+},
\]

(1.98)

\[
\frac{dn_{\text{He}^+}}{dt} = (\Gamma_{\text{He}0} + n_e \Gamma_{\text{eHe}0})n_{\text{He}0} + n_{\text{He}^+}n_e \alpha_{\text{He}^+} - (\Gamma_{\text{He}^+} + n_e \Gamma_{\text{eHe}^+} + n_e \alpha_{\text{He}^+})n_{\text{He}^+},
\]

(1.99)

\[
\frac{dn_{\text{He}^{2+}}}{dt} = (\Gamma_{\text{He}^+} + n_e \Gamma_{\text{eHe}^+})n_{\text{He}^+} - n_{\text{He}^{2+}}n_e \alpha_{\text{He}^{2+}},
\]

(1.100)
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with the closing conditions

\[ n_{H_0} = n_H - n_{H^+}, \quad (1.101) \]
\[ n_{He0} = \frac{Y}{4(1-Y)} n_H - n_{He^+} - n_{He^{2+}}, \quad (1.102) \]
\[ n_e = n_{H^+} + n_{He^+} + n_{He^{2+}}. \quad (1.103) \]

where \( n_i \) is number density \([\text{cm}^{-3}]\) of species \( i \) and \( Y \) is the helium mass fraction, taken to be \( Y = 0.24 \) (e.g. Olive & Skillman, 2004). \( \Gamma_i [\text{s}^{-1}] \) is the photoionization rate, and \( \Gamma_{ei} [\text{cm}^3 \text{s}^{-1}] \) is the collisional ionization rate, which are given in (1.83) and (1.91) respectively. \( \alpha_i [\text{s}^{-1}] \) is the recombination rate calculated in (1.89).

The fourth differential equation determines the temperature of the medium:

\[ \frac{dT}{dt} = \frac{(\gamma - 1) \mu m_H}{k_B \rho} [H_{\text{tot}} - \Lambda(n_i, T)] - 2H(t)T, \quad (1.104) \]

where \( \mu \) is the mean molecular weight of the gas, \( \gamma = 5/3 \), \( \Lambda(n_i, T) [\text{erg s}^{-1} \text{ cm}^{-3}] \) is the radiative cooling function, which is a combination of all of the cooling rates described in equations (1.90) to (1.96).

The final term of (1.104) represents cooling by the adiabatic expansion of the universe, where \( H(t) \) is the Hubble parameter (Peebles, 1993)

\[ H(t) = \Omega_m^{1/2} H_0 \cosh \left[ \frac{3H_0}{2} (1 - \Omega_m)^{1/2} t \right] a^{-3/2}. \quad (1.105) \]

In this equation, \( H_0 = 100 \ h \ \text{km s}^{-1} \ \text{Mpc}^{-1} \) is the present day Hubble parameter, \( \Omega_m \) is the matter density as a fraction of the critical density and \( a = (1 + z)^{-1} \) is the cosmological scale factor.

These equations are solved using a first order implicit integration scheme (Anninos et al., 1997) with a time-step that is typically comparable to the hydrogen ionization time-scale (\( t_{\text{ion}} = 1/\Gamma_H \)). When the rate of change in the electron number density is very low (\(< 10^{-8}\)), equilibrium is assumed for solving the ionic abundance equations, so that only the temperature equation needs to be integrated, allowing larger time-steps and increasing the speed of the calculation. Secondary ionisations by fast electrons are included using the results of Furlanetto & Stoever (2010).

1.4.2.1. Code tests

The limitations of the radiative transfer implementation described above are tested in detail in Appendix C of Bolton & Haehnelt (2007), to which the reader is referred for a full analysis of the code’s behaviour.

However, I have made several tests of my own implementation of the radiative trans-
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Temperature (10^4 K)
Distance (comoving Mpc)
\[ \alpha = 3.0, \ 10\text{Myr} \]
\[ \alpha = 1.5, \ 10\text{Myr} \]
\[ \alpha = 0.5, \ 10\text{Myr} \]
\[ \alpha = 3.0, \ 100\text{Myr} \]
\[ \alpha = 1.5, \ 100\text{Myr} \]
\[ \alpha = 0.5, \ 100\text{Myr} \]

Figure 1.10 Test of the one-dimensional time-dependent radiative transfer code used in this thesis. This plot reproduces the upper panels of Figure 5 of McQuinn (2012). Here, the temperature profile of the IGM around a point source emitting 10^{54} ionizing photons s^{-1} is shown after 10 (dashed lines) and 100 (solid lines) Myr. The source spectrum follows a power-law profile, with spectral indexes of \( \alpha = -0.5, -1.5 \) and \(-3\), plotted in black, red and blue respectively. In each case, \( z = 8 \) and a neutral, homogeneous IGM is initially assumed.

The code before applying it to my model SEDs in Chapter 3. One such test is shown in Figure 1.10. In this model, I calculated the temperature around a single point source with a luminosity of 10^{54} ionizing photons s^{-1}, propagating into an initially neutral medium at \( z = 8 \). The spectrum of the source is distributed between 13.6 – 4000eV, following a power-law spectrum with \( \alpha = 0.5, 1.5 \) and \( 3 \) plotted as black, red and blue lines respectively. The temperature profiles for each spectrum are plotted at 10 Myr (dashed lines) and 100 Myr (solid lines). Figure 1.10 can be compared directly to the upper panels of Figure 5 in McQuinn (2012), who made the same calculation with their own radiative transfer code. For this implementation, I neglected cooling due to adiabatic expansion and Inverse Compton processes, which were not included in the McQuinn (2012) calculation. There is a strong qualitative agreement between the two plots. These results, and the success of further tests in Bolton & Haehnelt (2007), mean that results from this model can be considered reliable, within the simplification constraints imposed by the implementation.
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1.5 X-RAYS IN THE EARLY UNIVERSE

1.5.1 Ionization

Approximately 370 000 years after the Big Bang, the baryonic gas in our Universe cooled sufficiently to recombine, becoming predominantly neutral, and radiation decoupled from matter to form the Cosmic Microwave Background (CMB: Penzias & Wilson, 1965). A wealth of observations show us that the baryons in the Universe are now predominantly ionized (Fan et al., 2006), which suggest the Universe has undergone an Epoch of Reionization (Barkana & Loeb, 2007).

Radiation with energy greater than 13.6eV is required to ionize hydrogen, the most abundant element in the universe. Theoretically, we can therefore link reionization to the epoch at which this radiation was first produced; the formation of the first stars and galaxies at around 100-250 million years after recombination (z ∼ 15-30; Bromm et al., 2009; Loeb, 2010).

1.5.1.1. Observational evidence

Observing the Universe during reionization is extremely challenging because of the huge distances involved. Observers are therefore largely limited to studying the final stages of reionization, at approximately z ∼ 6 (Fan et al., 2006; Robertson et al., 2010).

Gunn & Peterson (1965) predicted that even very low densities of neutral hydrogen could produce deep absorption troughs in the UV parts of distant spectra. The first observational evidence of such a feature was found by Djorgovski et al. (2001) in the quasar spectrum of SDSS J1030+0524 at z=6.28. Similar observations followed (Pentericci et al., 2002; White et al., 2003, 2005; Oh & Furlanetto, 2005), all of which indicated that some neutral hydrogen was still present at this redshift, and that reionization was in its final stages (White et al., 2005).

Quasar spectra can be modelled to predict the extent of ionized material around them, to calculate the fraction of neutral hydrogen in their surroundings (Mesinger et al., 2004). Mesinger & Haiman (2004) infer that \( f_{\text{HI}} \geq 0.2 \) at \( z = 6.28 \) from the spectrum of SDSS J1030+0524, and Wyithe et al. (2005) modelled the spectra of 7 quasars between \( 6 \leq z \leq 6.42 \) to show \( f_{\text{HI}} \geq 0.1 \). Yu & Lu (2005) also estimated that hydrogen is 80 - 90% ionized at \( z = 6.2-6.4 \), by observing the sizes of small regions of transmitted flux in quasar absorption spectra.

Less certain, from these spectra, is the rate at which reionization progressed. Some studies of quasar spectra (Fan et al., 2002; Cen & McDonald, 2002; Fan et al., 2006) suggest that the hydrogen neutral fraction increased by a factor of \( \gtrsim 10 \) from \( z=5.7-6.4 \), implying a brief, fast reionization epoch at \( z \sim 6 \). Alternatively quasar observations by Songaila & Cowie (2002) between \( z=4.42-5.75 \) predict a steadily evolving ionization rate.
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In fact, using semi-analytical simulations of the absorption spectra of QSOs, Gallerani et al. (2006) claim that it is not possible to tell the difference between early and late ionization models using these observations.

The evolution of the luminosity function of Lyman $\alpha$ emitting galaxies can be used to predict the progress of reionization (Hu et al., 2005), but these studies, too, are often in conflict. Hu et al. (2002) and Malhotra & Rhoads (2004) showed that the reionization epoch was beyond $z \sim 6.6$, with Ouchi et al. (2010) predicting that reionization mainly occurred at $z > 7$. Conversely, Kashikawa et al. (2006) have suggested that reionization was not completed until much later.

Other probes of the reionization era include the study of metal absorption lines in quasars (Oh, 2002; Furlanetto & Loeb, 2003; Becker et al., 2006) and WMAP polarisation data (Page et al., 2007; Spergel et al., 2007). The latter places a weak constraint on reionization timescales with the implied optical depths to reionization. WMAP results are in conflict with a rapid reionization epoch $z \sim 6$, implying a longer, more extended process (Bennett et al., 2003; Kogut et al., 2003).

Detecting sources from even earlier periods in the universe’s history is an active and competitive area of research. Recently these investigations have included detections of the earliest galaxies (e.g. Bouwens et al., 2011; Oesch et al., 2012), and attempts to constrain the bright end of the UV luminosity function at $z > 7$ (Yan et al., 2011; Bradley et al., 2012; Schenker et al., 2012; Lorenzoni et al., 2013; McLure et al., 2013; Oesch et al., 2013). For further insight into the era of the first luminous sources, future missions such as the JWST (Gardner et al., 2006) will be required.

Mesinger (2010) warn about over interpreting observations, arguing that there is still no direct evidence that reionization completed at $z \sim 5 - 6$, although they subsequently present tentative evidence for complete reionization at $z \sim 5.5$ (McGreer et al., 2011). Fan et al. (2006) provide a comprehensive review of observational studies of reionization and their interpretative uncertainties.

1.5.1.2. The process of reionization

In the absence of substantial high-redshift observations, theoretical studies currently drive our understanding of the reionization process. These have led to a “Swiss Cheese Paradigm” of reionization (Gnedin, 2000; Barkana & Loeb, 2001; McQuinn et al., 2007), whereby fully ionized regions formed in isolated pockets that expanded in a similar way to theoretical Strömgren spheres (Strömgren, 1939). Eventually these isolated bubbles overlapped, and the universe become mostly ionized, with just a few dense neutral regions remaining at $z \sim 6$. A schematic of this process is shown in Figure 1.11.

Theoretical (Miralda-Escudé et al., 2000), semi-analytical (Gnedin & Ostriker, 1997; Cen, 2003; Benson et al., 2006) and large-scale cosmological (Iliev et al., 2006) simulations have all been used to further develop the reionization model. Various simulations
have studied the growth and overlap of bubbles of ionized hydrogen (Gnedin, 2000; Furlanetto et al., 2004; McQuinn et al., 2007; Trac & Cen, 2007; Shin et al., 2008), and many models (e.g. Gnedin, 2004; Iliev et al., 2006) claim to be able to match observations such as the electron-scattering optical depths measured by WMAP (Spergel et al., 2007).

1.5.1.3. Can massive stars reionize the universe alone?

Currently the most promising candidate for the main source of reionization is UV radiation from massive stars in the first generations of galaxies (Barkana & Loeb, 2001; Wyithe & Loeb, 2003; Wise & Abel, 2008), which, despite being short-lived (∼10 Myr), are extremely luminous sources of hydrogen-ionizing UV photons (e.g. Schaerer, 2003).

However, their predicted contribution to reionization depends strongly on the initial mass function (IMF hereafter) of early stellar populations, and the proportion of UV photons that can escape into the IGM, known as the escape fraction, $f_{\text{esc}}$. Several numerical simulations show that reionization by massive stars is only consistent with the optical depth measurements from WMAP in the case of heavy IMFs, and high escape fractions (Ciardi et al., 2003; Haiman & Holder, 2003; Sokasian et al., 2004). Furthermore, in the case of a clumpy IGM, the requirements for escape fractions may be even more restrictive (Pawlik et al., 2009). For example, Grazian et al. (2012) conclude that the number of ionizing photons emitted from galaxies at $z \sim 7$ cannot keep the Universe re-ionized in a clumpy IGM unless $f_{\text{esc}} \geq 0.3$.

Actual high-redshift escape fractions are very uncertain. Gnedin et al. (2008) find restrictively low escape fractions in simulations of gamma ray burst host galaxies and many surveys (e.g. Bunker et al., 2010; McLure et al., 2010; González et al., 2010) claim that escape fractions are too low for the observable star forming galaxies to have driven reionization. This has led to the requirement of an unseen population of dwarf galaxies for complete reionization (Ouchi et al., 2010; McLure et al., 2010; Lorenzoni et al., 2011; Bouwens et al., 2012). This prediction is supported in simulations by Razoumov & Sommer-Larsen (2010) who find suitably high escape fractions of $f_{\text{esc}} \sim 80\%$ at $z = 10$, when low luminosity dwarf galaxies are included. On the other hand, Ishida et al. (2011) estimate a significantly higher star formation rate at $z \sim 9$ than inferred from high red-
shift galaxy searches that may negate the need for low luminosity populations. Similarly both Kistler et al. (2009) and Finkelstein et al. (2012) argue that the star formation rate in observed galaxies is sufficient for reionization without the need for an unseen population.

In conclusion, observations hint that star formation is the main driver behind reionization, yet it is far from certain whether massive stars alone can keep the universe ionized, and, even if they do, it may depend on an as yet un-observed low luminosity population. Consequently, the assessment of alternative/additional sources of reionization remains highly relevant to the field.

1.5.1.4. How might X-rays affect reionization?

Cosmological reionization of neutral hydrogen may have been profoundly influenced by the presence of X-rays, leading to different growth and morphology of ionized zones compared to those created by UV sources. In the first case, the ionization cross section of neutral hydrogen decreases as approximately $E^{-3}$, so that the mean free path of a UV photon is much shorter than that of an X-ray photon. Consequently, while UV photons are effective ionizers of high density gas in galaxies, X-rays have much longer mean free paths and can escape to ionize a much larger volume (e.g. Ripamonti et al., 2008). Furthermore, recombination is likely to be more important in UV ionized, denser regions than in the low density X-ray ionized regions. This has led Dijkstra et al. (2004) to estimate that $\sim 10$ UV photons are required to ionize a single hydrogen atom compared to $\lesssim 1$ X-ray photon. Finally, photoionization by X-rays produces more energetic electrons, which can cause additional ionizations (Shull & van Steenberg, 1985).

Theoretical investigations of X-ray sources suggest that they may have been able to “pre-ionize” large volumes of the IGM, unreachable by UV sources (Oh, 2001; Venkatesan et al., 2001; Madau et al., 2004; Ricotti & Ostriker, 2004). Indeed, their penetrative power has led some authors (e.g. Glover & Brand, 2003; Madau et al., 2004; Haiman, 2011) to propose a smoother global increase in ionized fraction with a more uniform morphology than the currently favoured Swiss Cheese Paradigm (Gnedin, 2000; Barkana & Loeb, 2001; McQuinn et al., 2007). Mesinger et al. (2013) show that the presence of X-rays may result in an extended epoch during which hydrogen is $\sim 10\%$ ionized before UV photons complete the process. Finally, X-ray radiation could further influence reionization by suppressing small scale structure formation, thereby lowering recombinations (Jeon et al., 2013).

1.5.2 Heating and galaxy formation

The presence of a photoionizing background at high redshifts may have exerted significant influence on galaxy formation by, for example, suppressing radiative cooling and star formation within haloes (Benson et al., 2002) or even inhibiting the initial collapse of
baryons onto low-mass dark matter haloes (Efstathiou, 1992; Thoul & Weinberg, 1996). Stellar feedback may have also affected the extremes of the galaxy mass function, truncating and flattening its high- and low-mass ends respectively (Kauffmann et al., 1999; Cole et al., 2001; Benson et al., 2003; Tegmark et al., 2004; Bower et al., 2006; Croton et al., 2006).

Currently, most simulations of high-redshift feedback address only UV wavelengths, by depositing mechanical energy and momentum into cosmological simulations to represent radiation from massive stars, supernovae and AGN (e.g. Sales et al., 2010; Finlator et al., 2011; Stinson et al., 2013; Sobacchi & Mesinger, 2013). However, there have been several theoretical and computational studies of the heating and feedback effects of a soft X-ray background in recent years, which show that X-rays could have significantly raised the temperature of the IGM (Oh, 2001; Machacek et al., 2003). Additionally this heating may have delayed the onset of reionization by increasing the Jeans mass (Ricotti & Ostriker, 2004; Kuhlen & Madau, 2005).

A well-known effect of UV radiation in the Lyman-Werner bands ($11.2 < E_\gamma < 13.6$eV) is to photo-dissociate H$_2$ via the Solomon process. The presence of a UV background therefore inhibits the collapse of small protogalaxies ($T_{\text{vir}} < 10^4$K) for which cooling is dominated by H$_2$ (Dekel & Rees, 1987; Haiman et al., 1997; Ciardi et al., 1998; Machacek et al., 2001; Oh & Haiman, 2002). It has been suggested (Haiman et al., 2000; Oh, 2001; Machacek et al., 2003) that X-ray photons may increase the electron fraction through low level ionization, promoting the production of H$_2$ and hence offsetting photodissociation. A high-redshift X-ray background would therefore have allowed cooling to occur at higher virial temperatures, reducing $t_{\text{crit}}$ (the lowest temperature at which gas haloes can cool) and hence promoting galaxy formation. Note, however, that Glover & Brand (2003) argue that this effect is only likely to have influenced haloes with $T_{\text{vir}} > 2000$K.

1.5.3 Sources of X-rays

In the past two decades, a stream of alternative sources have been put forward as potential ionizers. In particular, many studies have discussed X-ray emission from quasars as a potential rival to the ionizing power of massive stars (Ricotti & Ostriker, 2004; Ricotti et al., 2005; Volonteri & Gnedin, 2009). However, several subsequent studies have concluded that they were unlikely to be important contributors to reionization (Dijkstra et al., 2004; Jiang et al., 2008). Shankar & Mathur (2007), for example predicted that the quasar contribution to reionization was at maximum $\sim$30%, while Srbinovskv & Wyithe (2007) estimated it to be $< 14%$. More recently, observations by Willott et al. (2010) suggest that their contribution to the intergalactic ionizing flux at $z \sim 6$ was 20-100 times lower than that required for complete ionization.

Other potential ionizing sources that have been investigated include miniquasars (Madau
et al., 2004; Zaroubi et al., 2007), high velocity structure formation shocks (Miniati et al., 2004; Dopita et al., 2011; Wyithe et al., 2011) and dark matter annihilations (Hansen & Haiman, 2004; Belikov & Hooper, 2009). In addition, massive stars themselves may have continued ionizing during their post-main sequence lives, predominantly as sources of X-rays. Oh (2001) predicted that X-ray emission from supernovae in star forming regions should be large and comparable energetically to UV emission. The role of supernovae has been explored further by Johnson & Khochfar (2011) who examined how the strong shocking of the ISM associated with supernovae leads to the production of ionising photons with harder spectra and larger escape fractions than UV photons. They estimated that such X-rays may have enhanced the fraction of hydrogen ionized by stars by \(\sim 10\%\). An additional X-ray source in the early universe may have been the accretion of the ISM onto isolated stellar mass black holes (Wheeler & Johnson, 2011).

One source of X-rays which has gained particular attention in recent years is that of stellar-mass X-ray binaries. As the focus of this thesis, these sources will be addressed in depth in Section 1.6.

### 1.5.4 Observational constraints

Several observations can be used to constrain early X-ray emission. For example, the high-redshift abundance of hard (>1keV) X-ray photons is limited by the unresolved soft X-ray background (SXB) as they were optically thin to the IGM and remain unabsorbed to the present day (Dijkstra et al., 2004; Salvaterra et al., 2007; McQuinn, 2012). These high energy X-rays would have also altered the ionization state of metal absorption lines in \(z \sim 3\) quasar spectra (McQuinn, 2012).

Upper limits on the kinetic Sunyaev-Zeldovich signal support homogenous ionization, for which X-rays are an obvious candidate (Mesinger et al., 2012; Reichardt et al., 2012; Visbal & Loeb, 2012). Furthermore, the level of X-rays can determine whether the redshifted 21cm spin-flip line appears in absorption or emission (Furlanetto et al., 2006; Mesinger et al., 2013). Feedback from X-rays must also conform to measurements of the thermal state of the IGM (Bolton et al., 2011; Raskutti et al., 2012), and timescales for \(\text{He}^+\) ionization (Hui & Haiman, 2003; Dixon & Furlanetto, 2009; McQuinn et al., 2009; Becker et al., 2011; Worseck et al., 2011).

For an in-depth study of the observational constraints placed on X-ray emission in the early universe, the reader is referred to McQuinn (2012).

### 1.6 THE ROLE OF HMXBS IN THE EARLY UNIVERSE

Accretion onto black holes in High Mass X-ray binaries is a source of ionizing energy that has been the focus of several recent studies. These sources are now believed to be
proportionally more luminous at higher redshifts than in the present day, due to an anti-
correlation of HMXB mass with metallicity (Dray, 2006; Crowther et al., 2010; Mirabel
et al., 2011). In fact, Fragos et al. (2013) have argued that these systems are likely to
dominate X-ray luminosities over AGN at $z \gtrsim 6 - 8$.

Several previous studies have highlighted the potential influence of X-ray binaries on
their surroundings. For example, their feedback effects have been addressed by Glover
& Brand (2003) and Justham & Schawinski (2012), the latter of whom predict a signi-
ficant, stochastic effect on the evolution of dwarf galaxies that could be responsible for
the diversity we see today. High Mass X-ray Binaries (HMXBs) have been evaluated as
a source of reionization in Power et al. (2009), and by Mirabel et al. (2011). Both studies
suggest that the increased mean free path of X-rays and their potential for secondary ion-
izations lead to the ionizing potential of HMXBs equalling that of their progenitors. The
X-ray energy emitted by HMXBs in the early universe, relative to stellar radiation alone,
has been quantified in both Fragos et al. (2013) and Power et al. (2013).

Qualitative reasoning and calculations of the energetics imply a significant (up to
10%) contribution from HMXBs to reionization. However, McQuinn (2012) argues,
based on assuming a hard power-law spectrum for HMXBs, that they cannot have con-
tributed to reionization without the SXB constraint being violated. Kaaret (2014), on
the other hand, predicts a much weaker constraint from the SXB, because HMXBs show
spectral curvature above 2keV. In another study, Jeon et al. (2013) computed the effects
of the very first HMXBs on their environments using a zoomed hydrodynamical simu-
lation, incorporating a fully integrated radiative transfer model. They found little or no
direct ionization of the IGM from X-rays, although they predicted a net positive effect on
reionization due to the suppression of small scale structure.

### 1.6.1 A study by Power et al (2009)

The motivation behind Chapters 2 and 3 of this thesis stems from a study carried out by
Chris Power, Celine Combet, Mark Wilkinson and Graham Wynn, my supervisor, before
my arrival at Leicester University.

Power et al. (2009) explored the possibility that HMXBs in primordial globular clusters
at high redshifts could boost the ionizing power of the cluster. Using a Monte Carlos stel-
lar population synthesis model, Power et al. (2009) created a cluster of $N = 10^6$ stars with
metallicity $Z=0$ and a Kroupa IMF, and calculated the ionizing energy produced by the
main sequence (MS) stellar and HMXB components of the model. Figure 1.12 shows
the number of ionizing photons emitted from the cluster against time, with and without
HMXBs. These were calculated by assuming that main sequence stars radiate as typical
blackbodies, and that HMXB spectra are power laws, with slope $\alpha = -1$. They showed
that after $\sim 10$ Myr, once the most massive stars have ended their main sequence lives,
the number of ionizing photons produced by HMXBs becomes dominant over the MS contribution. Overall, they found that the harder ionizing spectra, combined with the enhanced escape fraction for X-rays, implied that HMXBs could be as much as a factor of $\sim 10$ more efficient at ionizing the IGM as their MS progenitors.

The analysis in Power et al. (2009) focussed on the ionising power of HMXBs in globular clusters because (i) the inferred ages of metal poor globular clusters imply that they formed at $z \gtrsim 6$ (cf. Brodie & Strader, 2006); (ii) the relationship between the initial mass functions (IMF) of stars and the dynamical evolution of clusters is well understood (cf. Vesperini & Heggie, 1997); and (iii) the escape fraction for UV photons was likely to be large, assuming that globular clusters followed similar orbits in the past to those which they follow today (cf. Ricotti, 2002, who explored the ionizing power of globular clusters) allowing for straightforward comparison of the ionizing power of the stellar population in UV and X-rays. However, a natural extension to this project would be to broaden its...
application to general high-redshift star forming regions.

These results suggest that HMXBs may have provided an important enhancement to the ionizing power of massive stars; a finding with potentially significant consequences for hydrogen reionization. By boosting the ionizing luminosity and escape fractions of stellar clusters, HMXBs may be able to address the current inadequacies of stellar-only reionization theory, overviewed in Section 1.5.1.3. As a consequence, assessing the implications of these initial findings to the epoch of reionization has been the primary goal of my PhD.

1.6.2 Structure of the thesis

The remainder of this thesis comprises three science chapters and a final conclusions chapter, the contents of which are summarised below.

In Chapter 2, I evaluate the ionizing power of HMXBs at redshifts $z \sim 6$ by further developing the Monte Carlo stellar population model described in Power et al. (2009). Most importantly, I adopt a more realistic HMXB template spectrum, based on that of the galactic HMXB, Cygnus X-1, in its high state. I monitor the ionizing energy of the cluster as a function of time, with and without a HMXB phase, and use these results to measure the boost in ionizing power that HMXBs add to the cluster. I then derive fitting formulae for this boost that can be used as a simple implementation of HMXB feedback in numerical simulations. This is followed by a discussion of the implications of these results for the involvement of HMXBs in reionization, in which I also calculate their contribution to the present day soft-X-ray background, to confirm that they do not violate this important constraint.

In Chapter 3, I continue my investigation into role of HMXBs on their high redshift environments. Using a one-dimensional radiative transfer code, I predict the ionization and temperature profiles surrounding the coeval stellar population I modelled in Chapter 2. I consider both uniform density surroundings and a cluster embedded in a $10^8 M_\odot$ NFW halo. In the context of this starburst model, I show that HMXBs do not make a major contribution to reionization or IGM heating. Specifically, HMXBs in a constant density environment produce negligible enhanced ionization or heating, because of their high-energy SEDs and short lifetimes. For NFW profiles, penetrating HMXB photons can stall recombinations behind the initially ionized front, keeping it partially ionized for longer. I discuss the reasons behind these effects, and consider the conditions required for HMXBs to significantly influence their high-redshift environments.

In Chapter 4, I compare the orbital period distributions of black hole and neutron star LMXBs in the Ritter-Kolb catalogue (Ritter & Kolb, 2003) to show that there is statistical evidence for a dearth of black hole systems at short orbital periods ($P_{\text{orb}} < 4$ hour). I then investigate whether this could be due to black hole LMXBs being preferentially hidden.
CHAPTER 1. INTRODUCTION

from view at short orbital periods. For this purpose, I estimate the outburst properties and orbital period distribution of black hole LMXBs using two models of the transition to radiatively inefficient accretion: an instantaneous drop in accretion efficiency ($\eta$) to zero, at a fraction (f) of the Eddington luminosity ($L_{\text{Edd}}$), and a power-law efficiency decrease, $\eta \propto \dot{M}^n$, for $L < f L_{\text{Edd}}$. I show that a population of black hole LMXBs at short orbital periods can only be hidden by a sharp drop in efficiency, either instantaneous or for $n \geq 3$.

Finally, in Chapter 5, I summarise the main results of my thesis, and discuss avenues for future research.
The Influence of High Mass X-ray Binaries on the Primordial Inter-Galactic Medium: Ionizing Energy

I had a dream which was not at all a dream
The bright sun was extinguished, and the stars
Did wander darkling in the eternal space,
Rayless and pathless...

Lord Byron, Darkness
In this chapter, I begin my investigation into the influence of accretion energy from HMXBs on their high-redshift environments. For this purpose I develop a Monte Carlo stellar population simulation that includes a HMXB phase, and incorporates observationally motivated HMXB abundances and spectra. I estimate the boost that a HMXB phase can provide to the ionizing power of a massive stellar cluster and parameterise it in a form that can be inputted into cosmological simulations of the epoch of reionization.

The content of this chapter is based on the publication of Power et al. (2013) for which I was second author. My contribution to the paper was to further develop and run the stellar population synthesis model described below, to form our main results. I was responsible for writing up the results section and was also directly involved making the SXRB calculation of Section 2.4.1. The photoionization-recombination calculation of Section 2.4.2, which is not included in Power et al. (2013), is also my own work. In writing up this chapter, I rewrote all parts of the paper that were lead by other authors.

2.1 INTRODUCTION

A broad range of observational data provide compelling evidence that the Universe underwent an “Epoch of Reionization” within the first $\sim 1$ billion years after the Big Bang (e.g. Ouchi et al., 2010; Mesinger, 2010; Shull et al., 2011; McGreer et al., 2011). During this period, neutral hydrogen was “re-ionized” by a background of ultra-violet (UV) and X-ray radiation produced by the first generations of stars and galaxies (e.g. Barkana & Loeb, 2007; Robertson et al., 2010).

In Section 1.5, I summarised our current understanding of how reionization took place, and the nature of the radiating sources responsible for it. UV luminous massive stars at redshifts $z \gtrsim 6$ are widely accepted to have played a important part in the reionization of the universe (Wyithe & Loeb, 2003; Wise & Abel, 2008; Wise, 2012). However, uncertainties about the escape fraction of UV radiation in the IGM (Gnedin et al., 2008; Ouchi et al., 2010), the shape of the luminosity function (Yan et al., 2011; Bradley et al., 2012; Schenker et al., 2012; Lorenzoni et al., 2013), and the star formation rate at high-redshifts (Kistler et al., 2009; Finkelstein et al., 2012), mean that it is far from certain whether radiation from massive stars alone is responsible for complete hydrogen reionization.

However, the ionizing lives of massive stars may not be over once they leave the main sequence. Johnson & Khochfar (2011) investigated the ionizing power of supernovae and showed that they may enhance stellar ionization up to 10%. Furthermore, hydrodynamic simulations have shown that a large fraction of stars in primordial galaxies form in binaries (Krumholz et al., 2009; Stacy et al., 2010; Turk et al., 2009). This leads to a further potential source of ionizing power: a HMXB phase, during which the X-ray luminosity of the system is dominated by accretion onto the black hole or neutron star primary.

The contribution of HMXBs to reionization has received growing interest in the past
decade (Power et al., 2009; Mirabel et al., 2011; Fragos et al., 2013). Mirabel et al. (2011) found that the ionizing power of a stellar mass black hole could be greater than that of its progenitor. Haiman (2011) then used this argument to propose an entirely different process of ionization to the currently favoured “Swiss Cheese Paradigm” (Section 1.5.1.2; Gnedin, 2000; Barkana & Loeb, 2001; McQuinn et al., 2007). Instead, they predict a gradual, uniform, reionization epoch, during which a photoionizing X-ray background pervades the IGM. Other research relating to the influence of stellar mass black hole binaries on the early universe is reviewed in Section 1.6

The motivation behind this study is the publication by Power et al. (2009), which is summarised in Section 1.6.1. Power et al. (2009) modelled a globular cluster containing both main sequence stars and a coeval population of HMXBs, to compare the number of ionizing photons emitted from the cluster. Under the model constraints, they found that HMXBs could be up to 10 times more efficient at ionizing the IGM as their main sequence predecessors.

Rather than being unique to globular clusters, the formation of HMXBs should be a generic by-product of high mass star formation (cf. Helfand & Moran, 2001). This is consistent with the results of Mineo et al. (2012), who studied a statistical sample of nearby star-forming galaxies and found that luminous compact X-ray sources – with properties mirroring those of HMXBs in our Galaxy – are good indicators of recent star formation activity. It is therefore informative, and reasonable, to apply this model in a more general context, to star forming regions in the early universe.

Cosmological simulations that track the progress of reionization and galaxy evolution typically consider the feedback of massive stars, supernovae, and AGN. However, Justham & Schawinski (2012) argue that X-ray binaries should also be included in galaxy modelling, as their feedback effects are likely to be both important and qualitatively different to that of the other sources. Cosmological simulations generally model ionizing radiation by the deposition of mechanical energy and momentum into the surroundings of energetic sources. A simple way to incorporate HMXB feedback into these models would be to add additional energy into star forming regions. In this study, therefore, I aim to provide quantitative prescriptions for the boost in ionizing power that HMXBs provide to a stellar cluster, which other researchers can easily implement into their large-scale simulations.

Hard photons ($\gtrsim 1$keV) at $z \gtrsim 6$ had such long mean free paths that they would not have been absorbed locally. Instead, having redshifted to the present day, they are expected to form part of the unresolved soft X-ray background (SXB, Dijkstra et al., 2004; Salvaterra et al., 2007; McQuinn, 2012). Hence, conforming to this constraint is an essential test for any model of X-ray emission in the early universe. McQuinn (2012) used this constraint to argue that HMXBs could not have formed in sufficient quantities to contribute significantly to reionization. However, he assumed a power-law HMXB spectrum,
unlike those exhibited by sources in our own galaxy. In my models, I make no assumptions about the importance of HMXBs to reionization, yet it is still important to test them against the SXB constraint. Consequently, the choice of a plausible HMXB spectrum for this prediction is a vital part of this study, and represents a significant advancement of the model in Power et al. (2009).

In this chapter, I further develop the Monte-Carlo stellar population simulation of Power et al. (2009), to quantify the boost in ionizing power for a stellar cluster that includes a HMXB phase. I parametrise this enhancement in a form suitable for incorporating into cosmological simulations, and I discuss, qualitatively, the effect that HMXBs may have had on their high-redshift environments. In Section 2.2, I describe the stellar population synthesis, and my choices of HMXB abundance, lifetimes and spectra. Next, in Section 2.3, I run the models and provide quantitative prescriptions for their results. In Section 2.4, I then discuss the implications of these results and estimate the contribution of HMXB photons to the SXB. Furthermore, I use a static photoionization-recombination calculation to make general comments on the potential ionizing power of HMXBs. Finally, a summary of my findings can be found in Section 2.5.

### 2.2 STELLAR POPULATION SYNTHESIS

In this section, I describe the Monte Carlo stellar population code I used to estimate the ionizing energy from the cluster. This model is based on Power et al. (2009), but has been expanded to include source spectra so that the time-dependent SED of the cluster can be monitored.

#### 2.2.1 Star formation

I model an individual starburst event in which $10^6$ stars are formed instantaneously, with stellar masses in the range $0.01 \leq M_*/M_\odot \leq 100$. The population is required to fit one of three possible Initial Mass Functions (IMFs), following a Salpeter (Salpeter, 1955), Kroupa (Kroupa, 2001), or a more top-heavy Chabrier distribution (Chabrier, 2001).

I assume that all massive stars are formed in pairs, and binary parameters are then assigned following the approach of Dray (2006), with companion masses being drawn from a uniform distribution between $0.01 \leq M_*/M_\odot \leq 100$ and orbital periods distributed uniformly in logarithm between 1 and $10^4$ days.

Main sequence lifetimes are estimated using the results of Marigo et al. (2001, 2003), Schaerer et al. (1993) and Meynet & Maeder (2000) for metallicities of $Z = 0, 0.008$ and $0.02$ (i.e solar metallicity) respectively; I explore the $Z = 0$ case in the results section.
CHAPTER 2. THE IONIZING POWER OF PRIMORDIAL HMXBS

2.2.2 HMXB formation

A minimum pre-requisite for HMXB formation is that the primary must evolve to become a neutron star or black hole. From Figure 1 of (Heger et al., 2003), this requires a main sequence mass, $M_\ast \geq 8M_\odot$. Additionally, the secondary must have $M_\ast \geq 3M_\odot$ to fit the HMXB and intermediate mass binary definitions. All binary systems with these mass constraints become candidates for HMXB formation.

To become a HMXB, stars in a massive binary must not merge during the primary’s main sequence and post-main sequence evolution. This is determined by both the initial binary separation and the details of stellar evolution at low metallicities. At this point, the binary must still survive the supernova of the more massive star, which depends on its initial mass, the fraction of mass it loses, and the kick velocity during this phase. All these details are highly uncertain at high-redshifts, because the evolution of lower-metallicity stars may be very different to present day analogues. For example, they are likely to retain more of their mass because stellar winds are inefficient at low metallicities (c.f. Heger et al., 2003), meaning more primaries will become black holes (Heger & Woosley, 2010). Since less mass is lost when black holes are formed, the natal kick velocity may be lower. Consequently, more binaries may remain bound and survive to become HMXBs at high redshifts.

For my stellar population model, I estimate the remnant masses of candidate HMXB primaries using Figure 3 of Heger et al. (2003) and then calculate the revised binary parameters. If the system loses more than half of its mass in its supernova phase, it becomes unbound and does not evolve to become a HMXB; this removes $\sim 70\%$ of all systems in the case of a Kroupa IMF.

Finally, I introduce a “survival fraction” ($f_{\text{sur}}$) which indicates the likelihood of an individual candidate binary becoming a HMXB. I draw $f_{\text{sur}}$ of the remaining candidates ($\sim 30\%$ after supernova phase) at random to become HMXBs. This parameter captures the various uncertainties that further prevent massive binaries from evolving into HMXBs, and was first introduced in Power et al. (2009). With this approach, it is possible to investigate the potential effectiveness of HMXBs as sources of reionization in a robust manner, even when the factors determining the value of $f_{\text{sur}}$ are poorly understood. This parameter is further discussed in Section 2.4.3.

2.2.3 Spectral energy distribution and time dependence

The spectra and luminosity of main sequence stars and HMXBs used in the Monte Carlo stellar population synthesis are described below.
CHAPTER 2. THE IONIZING POWER OF PRIMORDIAL HMXBS

Table 2.1  Adopted mass-luminosity relation. This table displays the values of $\alpha$ and $\beta$ I adopt between the mass ranges $M_{\text{min}}$ and $M_{\text{max}}$ to determine the luminosity of main sequence stars using the mass-luminosity relation of (2.1).

<table>
<thead>
<tr>
<th>$M_{\text{min}}$</th>
<th>$M_{\text{max}}$</th>
<th>$\alpha$</th>
<th>$\beta$</th>
</tr>
</thead>
<tbody>
<tr>
<td>0.5</td>
<td>5</td>
<td>0.8</td>
<td>4.5</td>
</tr>
<tr>
<td>3</td>
<td>10</td>
<td>18</td>
<td>3.6</td>
</tr>
<tr>
<td>10</td>
<td>20</td>
<td>5.8</td>
<td>3.1</td>
</tr>
<tr>
<td>20</td>
<td>50</td>
<td>25.8</td>
<td>2.6</td>
</tr>
<tr>
<td>50</td>
<td>100</td>
<td>100.0</td>
<td>2.3</td>
</tr>
</tbody>
</table>

2.2.3.1. Main sequence stellar spectra

Stellar luminosity $L_*$ follows a mass-luminosity relation of the form,

$$\frac{L_*}{L_{\odot}} = \alpha \left( \frac{M_*}{M_{\odot}} \right)^\beta,$$  \hspace{1cm} (2.1)

where the constants $\alpha$ and $\beta$ depend on stellar mass. Their values and the stellar mass range in which they are applicable are summarised in Table 2.1. There is empirical evidence that the rate at which stellar luminosity varies with mass ($\beta = d \log L / d \log M$) decreases with increasing mass $M_*$. For example, Malkov (2007) find $\beta \sim 4.7$ at $M \sim 1 M_{\odot}$ to $\beta \sim 2.8$ at $M \sim 20 M_{\odot}$ (see their Table 6), while Vitrichenko et al. (2007) report that $\beta \sim 2.76$ over the mass range $20 \lesssim M/M_{\odot} \lesssim 50$. The values of $\alpha$ and $\beta$ in Table 2.1 reflect the decrease in $\beta$ with increasing stellar mass. The chosen values of $\alpha$ and $\beta$ are based on the functional form presented in (Malkov, 2007) for $M_* \lesssim 50 M_{\odot}$. Above $50 M_{\odot}$ the relation has been extrapolated, assuming a further reduction in $\beta$.

2.2.3.2. HMXB spectra

I assume that HMXBs radiate until the death of their secondary, with continued black-body emission from the secondary, and the emission of a characteristic spectrum from the accreting black hole/neutron star. Although these lifetimes are likely to be overestimated in terms of individual systems, whose Roche Lobe filling phases will be brief, they are reasonable, in a statistical sense, when considering the cluster as a whole; see Section 2.4.3 for a discussion.

As previously noted, the ratio of black hole to neutron star HMXBs may be higher at low metallicities (Heger & Woosley, 2010). Cantalupo (2010) find observational evidence to this effect, and Linden et al. (2010) predict that the formation rate of HMXBs, in which the primary is a black hole, could be a factor of $\sim 10$ higher than at solar metallicity. This should, in principle, shape the low-metallicity HMXB luminosity function (cf. Dray, 2006). However, a simplifying assumption is made, to draw HMXB luminosities from a Weibull distribution with a typical luminosity of $L_X \sim 10^{38}$ erg/s but capped such that
they do not accrete at super-Eddington rates; this sets an upper limit of approximately 
\[ L_X \simeq 1.26 \times 10^{38} (M_1/M_\odot) \text{ erg/s} \] on the luminosity of a HMXB with primary mass \( M_1 \).

This approach gives a distribution that is consistent with the luminosities of compact X-ray sources in nearby galaxies whose X-ray binary populations are dominated by HMXBs (cf. Figure 1 of Gilfanov et al., 2004).

In Power et al. (2009), the HMXB spectra were expected to follow a simple power-law profile with index \(-1\) between a chosen range of energies. However, for a more accurate treatment of the energy output of HMXBs, it is important to move beyond this simple approach and base model spectra on a more typical observed HMXB spectrum. My approach to modelling HMXB SEDs is based on the spectral profile of Cygnus X-1.

Cygnus X-1, a black hole with \( M = 8.7 \pm 0.8 M_\odot \) (Shaposhnikov & Titarchuk, 2007), with a supergiant companion (HDE 226868), is the brightest HMXB in our galaxy, and, as such, is one of the most extensively studied X-ray sources in the sky. The spectrum of the source fluctuates between a hard state, characterised by a luminosity of \( L(2\text{-}10\text{ keV}) \sim 3 \times 10^{36} \text{ erg s}^{-1} \) and a hard power-law spectrum \( \propto E^{-\Gamma} \), with a photon index of \( \Gamma \sim 1.7 \) (e.g. Gierliński et al., 1997), and a soft state, with an order of magnitude higher luminosity, a strong blackbody component with \( kT \sim 0.5 \text{ keV} \), and a soft power-law tail with \( \Gamma \sim 2 - 3 \) (e.g. Dolan et al., 1977; Ogawara et al., 1982). These low and hard state spectra were shown in Figure 1.6 of Chapter 1.

The higher X-ray luminosity of the soft state and its larger proportion of UV and low energy X-ray photons (which are more likely to be Compton thick to reionization, due to a higher absorption cross-section: see discussion) suggest that the soft state will contribute most significantly to the ionization of the source’s surroundings.

For this model, I adopt distinct spectral shapes for low-hard and high-soft states. However, the duration of these two states is currently uncertain; for example, Cygnus X-1 is observed to be predominantly in its low hard state, but analogous systems such as LMC X-3 (Val-Baker et al., 2007) appear to spend more time in their high soft state. Nor is it understood how the nature and duration of these states depend on factors such as metallicity. For this reason, I make the simplifying assumption that HMXB spectra do not vary in time and instead introduce a threshold in X-ray luminosity of \( 10^{37} \text{ erg s}^{-1} \), above (below) which a source is in a soft (hard) state, and model

- the low-hard state as a power-law of slope \(-0.8\) between 2-10keV, the energy range in which this spectrum profile is observed.

- the high-soft state as a composite blackbody and power-law spectrum; the power-law has a slope \(-1.1\) and is normalised such that the ratio between the blackbody and power-law components matches that of Cygnus X-1, while the black-body temperature is calculated by assuming \( (L_X/L_{\text{CygX-1}}) = (T/T_{\text{CygX-1}})^4 \).
CHAPTER 2. THE IONIZING POWER OF PRIMORDIAL HMXBS

For the cluster as a whole, this model leads to just a few (or even one) high state sources being visible at any given time, which is equivalent to the switching on and off of many HMXBs over time.

2.3 RESULTS

In the following section, I show how the ionizing power of my model stellar population evolves over the first 150 Myr after formation, and provide fits to this solution that are suitable for cosmological models. I compare and contrast results in the presence and absence of HMXBs and, where appropriate, comment on the sensitivity of these results to the assumed value of the survival fraction ($f_{\text{sur}}$) and IMF.

2.3.1 The time-dependent SED

Figure 2.1 shows the spectral energy distribution of the stellar cluster as a function of time. It is formed by a composite of stars on the main sequence, which are modelled individually as blackbodies, and HMXBs, dominated by soft-state spectra, modelled as a combination of blackbodies and power-law components as described. This plot shows a cluster with my fiducial parameter choices: a Kroupa IMF, and $f_{\text{sur}} = 1$.

The amplitude of the blackbody corresponding to stars on the main sequence decreases with time, while its peak shifts to lower energies; this reflects the evolution of the most massive stars in the cluster – which dominate the UV-luminosity – off the main sequence. Over the same period, the amplitudes of both the blackbody and power-law components of the HMXBs decrease with time.

If a Chabrier IMF is used instead of a Kroupa one, the qualitative trends remain the same. However, the top-heavy distribution results in a systematic increase in the spectrum amplitudes. The converse is true for a Salpeter IMF, in which few massive stars form. The amplitude also increases linearly with the size of the population when I vary it between $N = 10^4$ and $N = 10^6$. Finally, reducing the survival fraction, $f_{\text{sur}} < 1$ suppresses the amplitude of the HMXB blackbody and power-law contribution, while leaving the main sequence blackbody component unchanged.

2.3.2 Ionizing power over time

To assess the ionizing potential of HMXBs in this model stellar cluster, I calculate the energy available to ionize neutral atomic hydrogen in the IGM as a function of time. I evaluate this by integrating the spectra plotted in Figure 2.1 between a lower limit of 13.6 eV, the minimum photon energy required to ionize atomic hydrogen, and an upper limit of $E_{\text{lim}}$, the energy above which atomic hydrogen becomes transparent to hard X-rays. I estimate $E_{\text{lim}}$ in the same way as for Power et al. (2009), by requiring that $\sigma(E_{\text{lim}}) =$
CHAPTER 2. THE IONIZING POWER OF PRIMORDIAL HMXBS

Figure 2.1 Composite stellar cluster spectrum over time. Spectra of the total energy output of the fiducial stellar population \((N=10^6, \text{Kroupa IMF, } f_{\text{sur}} = 1)\) after 5, 50, 100, and 150 Myr (red, green, blue and purple curves respectively). The spectral shape is a combination of the stellar UV blackbody and the HMXB blackbody + power-law component based on the high-state spectrum of Cygnus X-1. Highlighted in light blue is the range of energies that are important for hydrogen ionization, from the threshold ionization energy of 13.6 eV, to 1 keV, above which, photons are optically thin to the IGM at \(z \gtrsim 6\).

\[
\frac{(H(z)/c)/n(z)}{\sigma(E)}\text{ where } \sigma(E) \text{ is the ionization cross section of neutral hydrogen, } H(z) \text{ is the Hubble parameter at redshift } z, n(z) \text{ is the mean baryon density and } c \text{ is the speed of light.}
\]

For \(z \gtrsim 6\), \(E_{\text{lim}} \gtrsim 1\text{ keV}\). This energy range is highlighted in light blue, in Figure 2.1.

Figure 2.2 shows the ionizing energy from the cluster plotted as a function of time, split into main sequence stars (solid red line) and HMXBs for five different survival fractions. At early times, there is an initial burst of ionizing energy from the most massive stars in the system, but as these sources die off, some forming HMXBs, the stellar component drops off dramatically, falling below \(10^{38} \text{ erg s}^{-1}\) after 40 Myr. Provided \(f_{\text{sur}} > 0.1\), then HMXBs become the dominant source of ionizing energy after 20-30 Myr. This boosts the ionizing power of the system, extending the time over which the cluster can radiate energetically.

The bump in ionizing power between \(20 \lesssim t \lesssim 40\) Myr is a product of the HMXB definition. For black hole or neutron star formation, a primary must exceed \(M_\ast > 8M_\odot\). Because the number densities of stars increase with decreasing mass (for all IMF choices), it is statistically more likely for more stars to be drawn at this lower mass limit. The main sequence lifetime of an \(8M_\odot\) star is \(\sim 25\) Myr (Marigo et al., 2001). An excess of HMXBs therefore form around this time, producing the observed bump. After this time, HMXBs can no longer form, so the ionizing power steadily declines over time.
Figure 2.2  Cluster ionizing energy as a function of time. The upper panel of this figure shows the instantaneous ionizing energy produced by HMXBs as a function of time for five different survival fractions, ranging from $f_{\text{sur}} = 0.01$ to 1. Also plotted (in red) is the total ionizing energy released by the main sequence stellar component of the population. In the lower panels, the cumulative ionizing power of the same cluster components are shown.
2.3.3 A parameterized ionizing energy “boost factor”

The additional ionizing energy HMXBs provide can be quantified in the form of a “boost factor” $f_b = L_{\text{HMXB}}/L_{\text{MS}}$. With this factor, one can estimate how much additional ionizing energy is deposited into the IGM by HMXBs, just by knowing how much energy is deposited by stars.

In Figure 2.3, I show how $f_b$ varies with time for a stellar population of $10^6$ stars and $f_{\text{sur}} = 1$. In the left hand panel, the ionizing energy of the cluster ($\log[\text{erg s}^{-1}]$) is plotted against time ($\log[\text{Myr}]$) as a solid red line, for a cluster with a Kroupa IMF. Salpeter and Chabrier IMFs are then assumed in equivalent plots in the middle and right-hand panels respectively. Because the ionizing power of the MS component declines sharply after $\sim 10–20$ Myr, when the most massive stars end their lives, $f_b$ grows steadily with time, peaking at $f_b \sim 1600$ after $t \sim 180$ Myr before dropping off abruptly. The dashed curves in Figure 2.3 show fits to the numerical results that agree to better than 10%. I have fitted the results with two formulae, below and above the bump at 30 Myr. These are given by

$$\log(f_b) = \begin{cases} 10^{-5} p e^{q \log t} & \text{if } t < 30 \\ a + b \log t + c (\log t)^2 + d (\log t)^3 & \text{if } t > 30 \end{cases} \quad (2.2)$$

where $t$ is population age in Myr. The fits are valid up to approximately 180 Myr, above which the HMXB contribution falls rapidly to 0 (and $f_b$ drops to 1). The parameters $p$, $q$ and $a$ to $d$ are tabulated in Table 2.2.
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Table 2.2 Parameters for the boost factor, \( f_b \).

<table>
<thead>
<tr>
<th>IMF</th>
<th>( p )</th>
<th>( q )</th>
<th>( a )</th>
<th>( b )</th>
<th>( c )</th>
<th>( d )</th>
</tr>
</thead>
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<td>Kroupa</td>
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<td>6.73</td>
<td>54.2</td>
<td>-85.4</td>
<td>45.0</td>
<td>-7.61</td>
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<tr>
<td>Salpeter</td>
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<td>66.2</td>
<td>-104</td>
<td>55.1</td>
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<tr>
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<td>46.8</td>
<td>-69.2</td>
<td>34.3</td>
<td>-5.42</td>
</tr>
</tbody>
</table>

2.4 DISCUSSION

My results show that HMXBs dominate the ionizing power of a starburst after just 20 Myr, when the most massive stars have left the main sequence. Consequently they may enhance the influence of a stellar population on its surroundings by driving further ionization and heating. In this section, I show that these results are plausible by confirming that they do not violate the soft X-ray background constraint. I then discuss the potential implications of my findings with reference to a simple static photoionization recombination equation. Finally I address the uncertainties in several model parameters, and briefly evaluate alternative approaches.

2.4.1 Observational issues: the soft X-ray background

Since the Universe is optically thin to \( \gtrsim 1 \) keV photons at \( z \gtrsim 6 \), such photons would not have been absorbed locally. Instead, they are expected to contribute to the present day soft X-ray background (SXB; the mean specific background intensity of soft X-rays, Dijkstra et al., 2004; Salvaterra et al., 2007). Hickox & Markevitch (2007) obtained a limit on the unresolved component of the SXB of \( \sim 3.4 \pm 1.4 \times 10^{-13} \text{erg s}^{-1} \text{cm}^{-1} \text{deg}^{-2} \) for the excess flux between \( 1 - 2 \) keV, which corresponds to hard X-rays at \( z \gtrsim 6 \). Determining the contribution of the \( z \gtrsim 6 \) HMXB population to this SXB component is an important method of checking that a model is compatible with observations.

Several studies have calculated the contribution of high redshift X-ray sources to the SXB (Dijkstra et al., 2004; Salvaterra et al., 2005; Zaroubi et al., 2007). I follow the approach of Dijkstra et al. (2012), who looked at the contribution of star-forming galaxies to both the soft and hard X-ray backgrounds from high redshift to the present day.

I estimate the contribution from high-redshift HMXBs to the observed SXB by evaluating

\[
\text{SXB} = \frac{\Delta \Omega}{4\pi} \frac{c}{H_0} \int_{z_{\text{min}}}^{z_{\text{max}}} \frac{\dot{\rho}_\star \mathcal{L}_X(z)}{(1+z)^2 \delta(z)} dz,
\]

(2.3)

The terms in the equation above are described in the bullet points below:

- \( \Delta \Omega \sim 3 \times 10^{-4} \text{sr} \text{deg}^{-2} \), \( c \) is the speed of light, and the Hubble parameter \( H_0 = 100h \text{km s}^{-1} \text{Mpc}^{-1} \) with \( h = 0.71 \) (Komatsu et al., 2011),
• The lower redshift bound, \(z_{\text{min}}\) is fixed at \(z = 6\).

• The upper bound, \(z_{\text{max}}\), is the earliest redshift at which HMXB formation proceeds. \(z_{\text{max}}\) is allowed to vary between \(z = 50\) and \(z = 20\), corresponding to a Universe of age between \(\sim 48\) and \(\sim 180\) Myr. This range brackets the redshifts during which the first generation of dark matter haloes formed that were sufficiently massive to support cooling and subsequent star formation (cf. Figure 1 of Glover, 2005). However, I find that my results are insensitive to \(z_{\text{max}}\).

• \(E(z) = \sqrt{\Omega_m(1+z)^3 + \Omega_\Lambda}\). I assume values of, \(\Omega_M = 0.266\) for the matter density parameter, and \(\Omega_\Lambda = 0.734\) for the dark energy density parameter (Komatsu et al., 2011).

• \(\dot{\rho}_*\) is the comoving star formation rate density in units of \(M_\odot\, \text{yr}^{-1}\, \text{Mpc}^{-3}\).

I consider four possible values of \(\dot{\rho}_*\), which account for a range of potential star formation histories:

(i) a fixed value of \(\dot{\rho}_*(z = 6) \simeq 0.002\, M_\odot\, \text{yr}^{-1}\, \text{Mpc}^{-3}\), as measured at \(z=6\) by Bunker et al. (2010),

(ii) a fixed value of \(\dot{\rho}_*(z = 6) \simeq 0.17\, M_\odot\, \text{yr}^{-1}\, \text{Mpc}^{-3}\), as measured at \(z=6\) by Tanvir et al. (2012),

(iii) a fixed value of \(\dot{\rho}_*(z = 6) \simeq 0.02\, M_\odot\, \text{yr}^{-1}\, \text{Mpc}^{-3}\), estimated at \(z=6\) using the functional form of Hopkins & Beacom (2006)

\[
\dot{\rho}_*(z) = \frac{(u + vz)h}{1 + (z/w)^x}
\]

with parameters \(u=0.017\), \(v=0.13\), \(w=3.3\) and \(x=5.3\) from Cole et al. (2001);

(iv) a redshift-dependent value of \(\dot{\rho}_*\), estimated using (2.4).

• \(\mathcal{L}_X(z)\) is the “K-corrected” X-ray luminosity per unit star formation rate in the observed energy range \(X = E_1 - E_2\). I adopt the approach of Dijkstra et al. (2012) for calculating \(\mathcal{L}_X(z)\), rewriting it as \(\mathcal{L}_X(z) = c_X K(z)\). I assume the range of values for \(c_X\) of: \(2.6 \leq c_X/10^{39}\, \text{ergs}^{-1}\, (M_\odot\, \text{yr}^{-1})^{-1} \leq 3.7\) measured by Mineo et al. (2014) for compact resolved X-ray sources in galaxies (lower limit) and unresolved galaxies in the Chandra Deep Field North and ultra-luminous infra-red galaxies (upper limit). \(K(z)\) is then calculated as

\[
K(z) = \frac{\int_{E_1(1+z)}^{E_2(1+z)} EF(E) dE}{\int_{0.5\, \text{keV}}^{8\, \text{keV}} EF(E) dE}
\]

where \(F(E)\) represents the spectrum of a single HMXB, \(E_1 = 1\, \text{keV}\) and \(E_2 = 2\, \text{keV}\).
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Evaluating (2.3), I find that the contribution of HMXBs to the soft X-ray background is of order \( \sim 5 \times 10^{-16} \text{ erg s}^{-1} \text{ cm}^{-2} \text{ deg}^{-2} \). This result is a factor of \( \sim 1000 \) smaller than the observed limit, and depends on what assumptions I make for the star formation rate. In general, the higher the star formation rate density, the larger the SXB contribution, due to the assumption that the HMXB formation rate is proportional to the star formation rate.

This result is also subject to my choice of HMXB spectrum. To explore this dependence, I also consider several alternative HMXB spectra, which are plotted in Figure 2.4. Here, the solid curve corresponds to the cumulative spectrum derived from HMXBs in the fiducial stellar population, as measured at \( t=50 \text{ Myr} \), while the red dotted, blue short-dashed and green long-dashed curves correspond to power-law spectra \( (\propto E^{-\alpha}) \) with \( \alpha = 0.5, 1.5 \) and 2.5 respectively, which have been normalised such that the total energy emitted between 1 and 2 keV is the same for both composite and power-law spectra. The normalisation is not vital for the SXB calculation however, rather it is the shape that is important.

If HMXBs emit power-law spectra, I find that their SXB contribution \( \sim 2.15 \times 10^{-16} \text{ erg s}^{-1} \text{ cm}^{-2} \text{ deg}^{-2} \) for \( \alpha=0 \) compared to \( \sim 11.4 \times 10^{-16} \text{ erg s}^{-1} \text{ cm}^{-2} \text{ deg}^{-2} \) for \( \alpha=2 \), i.e. softer spectra make larger contributions to the SXB. Nevertheless, this contribution is still a factor of \( \sim 100 \) smaller than the observed limit.

2.4.2 A static photoionization-recombination equilibrium calculation

The importance of HMXBs to reionization cannot be argued on energetics calculations alone, as X-rays are likely to have a different effect on ionization compared to the UV photons from massive stars.

Since the absorption cross-section of neutral hydrogen decreases sharply with photon energy \( (\sigma \propto E^{-3}) \), X-rays can travel further though their surroundings than UV photons before they are absorbed. UV photons may therefore ionize denser galactic regions, while X-rays may penetrate into the less dense IGM. This means that recombination rates are likely to be much higher in UV ionized regions. Indeed, Dijkstra et al. (2004) estimate that \( \sim 10 \) UV photons are required to ionize a single hydrogen atom, compared to \( \ll 1 \) X-ray photons. X-ray ionization is also expected to prompt further ionization by producing energetic electrons. (Shull & van Steenberg, 1985). As a simple visual way to qualify these alternate effects, I solve a static version of the photoionization-recombination equilibrium equation to demonstrate the equilibrium solution that can be reached when timescales are not accounted for. For this purpose, I have used the calculation described in Zaroubi & Silk (2005); Zaroubi et al. (2007), which I write in the form:

\[
\alpha_{\text{HI}}^{(2)} n_{\text{H}}^2 (1 - x_{\text{HI}})^2 = \Gamma(r) n_{\text{H}} x_{\text{HI}}. \tag{2.6}
\]

\( \Gamma(r) \) 1I treat the unnormalised cumulative spectrum as indicative of the shape of a typical HMXB.
The total energy emitted between 1 and 2 keV is the same for both the composite and power-law spectra, respectively. The composite HMXB spectrum is as measured at $t = 50$ Myr in Figure 2.1, while the power-law spectra are normalised such that the total energy emitted between 1 and 2 keV is the same for both the composite and power-law spectra.

Here, $\alpha_{HI}^{(2)} = 2.6 \times 10^{13} (T / 10^4 K)^{-0.85} \text{cm}^3 \text{s}^{-1}$ is the recombination cross-section for a gas temperature $T$ (c.f. Kaplan & Pikelner, 1970), where I use $T = 10^4 K$. I evaluate the mean number density of hydrogen, $n_H \sim 1.9 \times 10^{-7} (1 + z)^3 \text{cm}^{-3}$ (Spergel et al., 2007) at $z = 10$. The ionization rate per hydrogen atom at distance $r$ from the source is given by

$$\Gamma(r) = \int_{E_0}^{\infty} \sigma(E) \mathcal{N}(E;r) \left[ 1 + \frac{E}{E_0} \phi(E,x_e) \right] \frac{dE}{E}, \quad (2.7)$$

where $\sigma_H(E) = \sigma_0 (E/E_0)^3$ is the bound-free absorption cross-section for hydrogen, with $\sigma_0 = 6 \times 10^{-18} \text{cm}^2$ and $E_0 = 13.6 \text{eV}$. $\mathcal{N}(E;r)$ is the number of photons per unit time per unit area at a distance $r$ from the source.

$$\mathcal{N}(E,r) = e^{-\tau(E,r)} \frac{F(E)}{4\pi r^2} \text{cm}^{-2} \text{s}^{-1}, \quad (2.8)$$

where the optical depth is evaluated as


Figure 2.5 Ionization profiles determined using a static photoionization-recombination equilibrium calculation. In the left-hand panel, the hydrogen neutral fraction is plotted against comoving distance from two model stellar clusters, without (solid red line) and with (solid blue line) a HMXB phase. These profiles were calculated by solving a static photoionization-recombination equilibrium calculation after inputting the radiation emanating from the clusters at 5 Myr. The right panel shows solutions to the same calculations at 100 Myr.

\[
\tau(E; r) = \int_0^r n_H x_H \sigma(E) dr,
\]

(2.9)

and the spectrum, \( F(E) \), is inputted from our Monte Carlo model of the stellar population. The term \([1 + (E/E_0) \phi(E, x_e)]\) accounts for secondary ionizations, where

\[
\phi(E, x_e) = 0.3908 \left[ 1 - (1 - x_e^{0.4092})^{1.7592} \right]
\]

(2.10)

and \( x_e = 1 - x_H \) (Shull & van Steenberg, 1985).

The equations are solved at 5 and 100 Myr in the left and right-hand panels of Figure 2.5, which show the hydrogen neutral fraction as a function of distance. The stellar-only model is plotted in red, and a model with a HMXB phase is shown in blue.

At 5 Myr, the stellar-only model fully ionizes a region of size \( \sim 2.4 \) Mpc, beyond which the medium sharply returns to neutral. When the equations are solved, instead, at 100 Myr, the ionized front is much smaller, extending only to \( \sim 50 \) kpc, as the most massive stars have left their main sequence and the ionizing energy is much lower.

The effect of a HMXB phase can be seen in both figures. The more penetrating X-rays do not increase the size of the fully ionized zones, but travel further into the IGM, leading to low levels of partial ionization. At 5 Myr, HMXBs do not significantly change the ionizing profiles because their contribution to the total ionizing energy emanating from the cluster is much lower than the main sequence one. However, after 100 Myr, HMXBs
change the shape of the ionization front, and create a smooth, extended, transition from the ionized to neutral IGM. They cause low levels of ionization beyond the main ionization front, which is in line with arguments from Haiman (2011), who suggest that HMXBs partially ionize the IGM and make the process of reionization smoother.

2.4.2.1. Important note: limitations of the static photoionization-recombination equation

I include the above calculation for illustrative purposes to show the different effects of UV and X-ray ionizing photons on the surrounding IGM. However, these static calculations are extremely limited for the following reasons;

(i) firstly, the fully ionized zones are over estimated because the spectrum at \( r \) is calculated as \( F(E) \exp^{-\tau} / 4\pi d^2 \) meaning the photons used to ionize the material within \( r \) have not been removed,

(ii) furthermore, ionization history has not been accounted for in the equilibrium calculation, which is why the ionized volume appears to recede over time, rather than recombing as a whole. The recombination cross-section, \( \alpha_{HI}^{(2)} \), is constant at all distances for a constant temperature, so recombination is likely to be more uniform.

(iii) Finally, note that this is an equilibrium solution that does not account for the timescales involved. It may be that this equilibrium cannot be reached during the cluster’s lifetime (this is discussed further in Chapter 3).

These issues mean that a static calculation is not sufficient for making absolute conclusions about the influence of X-rays on reionization, despite the use of similar arguments in recent literature (e.g. Zaroubi & Silk, 2005; Zaroubi et al., 2007). In Chapter 3, I will use a time-dependent radiative transfer code to run a more concrete version of the same calculation. I will compare the results to this calculation in Section 3.4.1.3.

2.4.3 Model limitations

My Monte Carlo stellar population synthesis relies on several assumptions that are observationally uncertain. Four of the most important uncertainties are addressed below.

**HMXB Luminosities** – In this study, HMXBs are given a typical X-ray luminosity of \( L_X \sim 10^{38} \text{erg s}^{-1} \). This is considerably higher than that of local HMXBs in our galaxy and the neighbouring Large and Small Magellanic Clouds. For example, there is just one source in the Small Magellanic Cloud with this luminosity (Shtykovskiy & Gilfanov, 2005).
Chapter 2. The Ionizing Power of Primordial HMXBs

The primary mass at formation ($M_{\text{prim}}$) versus the lifetime of the secondary once the primary has gone supernova for all model binaries in the cluster (red crosses). Also plotted is the population that survive the supernova to become HMXBs (green crosses). The blue line shows the median lifetime for HMXBs in ten logarithmic mass bins between 8 and 100 $M_\odot$.

However, this choice can be justified by reasoning that typical high-redshift environments are more likely to mirror populations in galaxies with high star formation rates (see e.g. Figure 1 of Gilfanov et al., 2004) and bursts of star formation from recent merger events in, e.g. the Antennae (Fabbiano et al., 2001) and the “Cartwheel” (Wolter & Trinchieri, 2004), both of which exhibit X-ray luminosities above $L_X \sim 10^{38}$ erg s$^{-1}$. If these are more analogous to HMXB forming regions in the early universe, then the model luminosity choice is reasonable.

The Survival Fraction – The survival fraction, $f_{\text{sur}}$, is a convenient way of capturing the uncertainties that surround the evolution of massive binaries into HMXBs. However, its value remains highly uncertain, as it is likely to depend on parameters that were different in the early universe. For example, $f_{\text{sur}}$ is likely to increase for stars with lower metallicities, since models of single stars with low metal content show that they collapse directly with no energetic natal kicks, and end as black holes (Heger et al., 2003; Meynet & Maeder, 2005; Gregory et al., 2009).

Estimation of $f_{\text{sur}}$ by comparison to more local observations is difficult. There are around 50 HMXB candidates in “The Antenna” star forming region, which has a mass of between $10^4$ and $10^6$ $M_\odot$ (Fall et al., 2005). In my fiducial stellar population model, with
a Kroupa IMF, I predict the formation of approximately 300 HMXBs per $10^5 \, M_\odot$, with the number of HMXBs scaling linearly with mass. If 10% of these systems have luminosities in excess of $L_X \sim 10^{38}$ erg s$^{-1}$, as expected for my chosen distribution, there would be around 30 bright HMXBs in a $10^5 M_\odot$ cluster, or 300 in $10^6 M_\odot$ cluster. Comparing the Antennae to the former case gives $f_{\text{sur}} \sim 1$ or for the latter, $f_{\text{sur}} \sim 0.1$. The mass and the spatial distribution of the Antennae is highly uncertain, meaning this comparison does not constrain $f_{\text{sur}}$. However, it shows that the $0.1 < f_{\text{sur}} < 1$ range broadly fits observations.

More detailed modelling in the future could potentially improve these estimates for $f_{\text{sur}}$. For example, one approach might be to populate a model Antennae galaxy with HMXBs based on its cluster distribution and to make mock observations for comparison (as suggested in Power et al., 2009).

### HMXB Lifetimes

In this stellar population model, HMXBs are expected to radiate until their companion leaves the main sequence. This leads to a range of lifetimes that are shown in Figure 2.6. Here, the lifetime and primary mass of all the binary systems in a $10^6$ star model, with a Kroupa IMF, have been plotted as red crosses. Also plotted, in green, are the primary masses and lifetimes of all the HMXBs that survive their supernova phase (with $f_{\text{sur}} = 1$). In blue, I also plot the median lifetime of HMXBs in ten logarithmically spaced mass bins between 8 and $100 \, M_\odot$. For systems that survive to become HMXBs, the lifetime is typically between 1-10 Myr, decreasing for more massive primaries.

Although these values are not unreasonable, it is likely that they are overestimates. HMXBs are probably not luminous for the entire lifetime of the secondary, as their Roche Lobe filling phases may be brief. However, rather than accurately modelling the lifetime of every individual binary, the code aims to monitor the HMXB population in a statistical sense. My model produces mostly low luminosity X-ray binaries, with one, or a few, brighter sources present at any given time, broadly matching observations of HMXBs in local stellar clusters. However, it is still important to recognise that this is an optimistic scenario.

### No Continuous Star Formation

This model treats star formation as a single burst over a short timescale. While this may be relevant for, e.g. globular clusters, star formation is likely to be a continuous process within young galaxies. Ongoing star formation may produce a different SED to that predicted in my models.

Fragos et al. (2013) have developed stellar population synthesis models for a similar purpose to this study. They also track the spectral energy distribution of a population, which includes a HMXB phase, but they assume ongoing star formation. The results of their models are summarised as a set of parameterizations for the X-ray luminosity as a function of star formation rate. These are complementary to my own findings, and either parameterization can be used in cosmological simulations depending on the set-up of the
model.

2.5 SUMMARY

In this chapter, I have developed a Monte Carlo stellar population model comprising main sequence stars formed in a single burst followed by a later HMXB phase. With observationally motivated choices for HMXB abundances, spectra and lifetimes, I argue that my models provide a good estimate of HMXB emission in a plausible context.

The code monitors the time dependent SED of the cluster, and its total ionizing power. I find that, while massive main sequence stars dominate the cluster’s ionizing output for the first 20 Myr, HMXBs take over after this time, provided their survival fraction is greater than 0.1. These results suggest that HMXBs may be able extend the energetic lifetime of a cluster, enhancing its ionizing power. I have parameterised this enhancement as a boost factor, which depends on the mass of the cluster. This can be incorporated into cosmological simulations which already model stellar feedback, so that they can include the effect of HMXBs.

I have made the important verification that my models do not violate the SXB constraint, finding that they fall 100-1000 times below the observed value, depending on my choice of HMXB spectrum. Solving a static photoionization recombination calculation, I have described general arguments supporting the additional ionizing effect that HMXBs are likely to have, as a precursor to the more rigorous radiative transfer calculation that will be the focus of Chapter 3. Finally, I have assessed the validity of various parameters in my model, and noted some of its limitations.
The Influence of High Mass X-ray Binaries on the Primordial Inter-Galactic Medium: Radiative Transfer

A light is gone from yonder sky
A star has left its sphere;
The beautiful, and do they die
In yon bright world as here?

[...]

Didst thou fade gradual from the time
The first great curse was hurled,
Till lost in sorrow and in crime
Star of our early world?

Letitia Elizabeth Landen, The Lost Star
CHAPTER 3. IONIZATION AND HEATING BY PRIMORDIAL HMXBS

In this chapter, I move beyond the energetics arguments of the previous chapter, to predict the actual ionizing and heating effects of HMXBs in the early universe. To assess the ionization state and temperature of the local IGM surrounding a stellar cluster with a HMXB phase, I apply a one-dimensional radiative transfer code to the time-dependent cluster SED modelled in Chapter 2. Although my previous chapter showed that HMXBs enhance the ionizing radiation emanating from a stellar cluster, here, I show that this leads to insignificant effects on the surrounding IGM. My results counter the general arguments in favour of high-redshift feedback from X-ray binaries that have been made in recent literature.

3.1 INTRODUCTION

X-ray radiation may have profoundly influenced its surroundings in the distant past, shaping the universe we see today. As described in Section 1.5.2, the presence of an X-ray background at high redshifts could have increased the fractional ionization of protogalactic gas, offsetting the effects of photodissociation from UV feedback, and allowing gas haloes to cool at higher virial temperatures (Haiman et al., 2000; Oh, 2001; Machacek et al., 2003). The timescales and progress of the cosmological reionization of hydrogen may have also been significantly affected by the presence of a photoionizing X-ray background. For example, X-rays may be able to “pre-ionize” large volumes of the IGM that less penetrating UV radiation is unable to reach (Oh, 2001; Venkatesan et al., 2001; Madau et al., 2004; Ricotti & Ostriker, 2004). Mesinger et al. (2013) found that this process will leave the IGM \( \sim 10\% \) ionized for an extended epoch before UV photons complete the process. This argument counters the currently favoured “Swiss Cheese Paradigm” for hydrogen reionization (Gnedin, 2000; Barkana & Loeb, 2001; McQuinn et al., 2007), which claims that UV radiation forms isolated ionized pockets around sources that slowly spread and combine. Instead, an X-ray background may drive a uniform, global increase in ionized fraction (Glover & Brand, 2003; Madau et al., 2004; Haiman, 2011). X-rays could further promote reionization by suppressing small scale structure and hence reducing recombinations (Jeon et al., 2013). Furthermore, X-ray heating may have delayed the onset of reionization by increasing the Jeans Mass (Ricotti & Ostriker, 2004; Kuhlen & Madau, 2005).

HMXB mass anti-correlates with metallicity (Dray, 2006; Crowther et al., 2010; Mirabel et al., 2011), so these systems were likely to have been more luminous at high redshifts. In fact, Fragos et al. (2013) recently used stellar population models to show that they dominated the X-ray luminosity of early galaxies over AGN at \( z \gtrsim 6 - 8 \). In Chapter 2, I estimated the contribution that X-rays from HMXBs could make to the ionizing power of a primordial stellar cluster. I found that X-rays dominate the ionizing energy of a starburst after just 20 Myr, so could have provided an important boost to reionization. The
X-ray energy emitted by HMXBs in the early universe, relative to stellar radiation alone, has been quantified in both Fragos et al. (2013) and my most recent study, described in Chapter 2 and published in Power et al. (2013).

Although qualitative energetics arguments suggest that HMXBs may have significantly boosted the ionizing power of their stellar progenitors, studies that account for HMXB spectral energy distributions are less conclusive. McQuinn (2012), for example, use a simple power-law HMXB spectrum to argue that their contribution to ionization must be negligible unless the SXB constraint is violated. Additionally, Jeon et al. (2013) studied the feedback effects of primordial HMXBs using zoomed hydrodynamical simulations, incorporating a fully integrated radiative-transfer algorithm. They found a net positive effect on reionization due to the suppression of small-scale structure, but no direct increase in ionization by HMXB X-rays. The discrepancy between these studies, and among general HMXB ionization predictions, highlight the need to move beyond energetics arguments when assessing the feedback potential of HMXBs; full radiative transfer calculations, with accurate HMXB spectra, are required to predict the proportion of energy that directly contributes to heating and reionization, and the timescales over which it acts.

The Monte Carlo stellar population model I described in Chapter 2 – with observationally motivated HMXB abundances, lifetimes, and, importantly, spectra – provide the ideal basis for an improved study of the influence of HMXBs within a population of stars. I have already shown that my models are not in conflict with the SXB constraint. I also made a simple model using a static calculation to estimate the effect of HMXBs on ionizing their surroundings, showing that X-rays may partially ionize additional IGM material. However, as I explained then, these static calculations are severely limited, and still constitute a qualitative argument.

I now use one dimensional radiative transfer calculations to estimate the influence of HMXBs in a realistic stellar cluster on their environment, with no presuppositions about their significance to reionization. In Section 3.2, I briefly describe the stellar population model and the radiative transfer calculation used in this chapter, each of which have been discussed previously in this thesis. I then use these models in Section 3.3 to predict the ionization and temperature profiles surrounding a high-redshift stellar cluster, with and without a HMXB component, in both constant and varying density environments. In Section 3.4, I discuss the factors that influence my results, and the conditions required for HMXBs to have a strong influence on their surroundings, before drawing my final conclusions in Section 3.5.
3.2 METHODS

In this study, I use my stellar population synthesis model described in Chapter 2 to simulate a stellar cluster forming in an instantaneous burst, containing both main sequence stars and HMXBs. I monitor the total energy released from this cluster, and enter it into a one-dimensional time-dependent radiative transfer code, first described in Bolton & Haehnelt (2007), to estimate the ionization state and temperature of the surroundings. These models are described below.

3.2.1 Stellar population synthesis

A full description of my model can be found in Section 2.2. I model the time dependent SED of a $10^6$ star cluster forming in a single starburst event. Stars are found in binary systems, following a Kroupa IMF and I assume a metallicity of $Z = 0$ for estimating their lifetimes. HMXBs form when the primary leaves its main sequence, provided it loses less than 50% of its mass during supernova phase, with a further survival fraction $f_{\text{sur}} = 1$. The choices above are described and justified in the previous chapter.

In Chapter 2, I modelled HMXB spectra as a blackbodies with a high energy power-law tail. I continue to use this as my fiducial spectrum model, but additionally consider three alternative power-law HMXB spectra. I use these to check the sensitivity of my results to spectrum choice, and also because HMXB spectra have been simplified as power-laws in other studies, so these models provide the opportunity for comparisons to their results. The four spectrum choices are defined below:

1. For my fiducial model, I base HMXB spectra on the galactic HMXB, Cygnus X-1, which alternates between a high luminosity soft state and and low luminosity hard state. In my model, if $L_X > 10^{37}$ erg s$^{-1}$, I use a soft state, composite black body and power-law spectrum (c.f. Dolan et al., 1977; Ogawara et al., 1982). The blackbody temperature is calculated by assuming $(L_X/L_{\text{CygX-1}}) = (T/T_{\text{CygX-1}})^4$, and the power-law has a slope of $\alpha = -1.1$, normalised such that the ratio between the blackbody and power law components matches that of Cygnus X-1. For $L_X < 10^{37}$ sources, I model a hard state power-law spectrum with $\alpha = -0.8$ between 2 and 10keV (c.f. Gierliński et al., 1997, 1999).

The total SED of the HMXBs in the cluster at 20 Myr is plotted as a green dashed line in Figure 3.1. The contribution from hard state spectra is negligible, so this represents a sum of the most luminous soft state HMXB composite black body and power law spectra. This model will subsequently be referred to as a BB+PL spectrum.

2. A power law spectrum, $F(E) = E^{-\alpha}$, and $\alpha = 0.5$ in the observed X-ray range of
Figure 3.1 
Choice of HMXB spectrum. Example spectrum snapshots from my Monte Carlo stellar population models, taken at 50 Myr, assuming a Kroupa IMF and $f_{\text{sur}} = 1$. The combined spectrum from all the main sequence stars is plotted in red. The other lines show the net spectral output from the HMXBs for four different choices of HMXB spectrum. My fiducial black body and hard power law spectrum, based on that of Cygnus X-1, is plotted as a light green dashed line. Power law HMXB spectra with $\alpha = 0.5, 1.5,$ and 2.5 are plotted as dark blue dotted, purple dotted and light blue dot-dashed lines respectively.

3. A power law spectrum with $\alpha = 1.5$, plotted in Figure 3.1 as a fine purple dotted line.

4. A power law spectrum with $\alpha = 2.5$, plotted in Figure 3.1 as a light blue dot-dashed line.

My model output is a global time-dependant SED which is then used as an input for my radiative transfer calculations.

3.2.2 1D radiative transfer code

The line-of-sight radiative transfer implementation I use in this chapter was written by Dr Jamie Bolton, University of Nottingham, and is described fully in Section 1.4.2, as well as in Appendix B of Bolton & Haehnelt (2007). It is based on a photon conservation algorithm developed by Abel et al. (1999), which was extended by Bolton et al. (2004) to incorporate multiple frequencies. The code accepts a time-varying SED ranging from $13.6 \text{eV} < E < 4 \text{keV}$ – the energy range important for hydrogen and helium ionization.
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The environment, which comprises these two elements, follows a one-dimensional density profile of the user’s choice.

3.3 RESULTS

The ionizing energy output of a single stellar population is initially dominated by its most massive stars, until they end their short main sequence lives. However, this may not be their only contribution to the ionizing power of the cluster; if they evolve to become accreting black hole systems, they can then provide further ionizing energy to their surroundings. Figure 2.2 showed the stellar cluster’s total ionizing energy (13.6eV < E < 1keV) against time, separated into a stellar component (red line) and a HMXB component (all other lines), where the energy produced by the HMXB component depends on the fraction of main sequence binaries that evolve to become HMXBs (30% × \( f_{\text{sur}} \); cf. Section 2.2). In this Chapter, I showed that the ionizing energy from HMXBs becomes dominant after 20-30 Myr, provided \( f_{\text{sur}} \gtrsim 0.1 \).

In the following section, I use the time-dependent SED of the same model cluster to study the ionization and heating of its surroundings. I compare my results for models with no post-MS systems to those with HMXBs (\( f_{\text{sur}} = 1 \)), and comment on the additional effect a HMXB phase may have on the cluster’s environment.

3.3.1 Heating and ionization of a constant density medium

In the first case, I consider the propagation of ionizing energy from my model cluster directly into the IGM, using the one dimensional radiative transfer code described in Section 1.4.2. I assume a uniform IGM density equal to the average matter density at a given redshift. I then follow the cluster’s emission from its birth, at \( z = 14.5 \), for 200 Myr, until \( z = 10 \). The IGM density decreases over this time due to cosmological expansion.

3.3.1.1. Ionization profiles

Snapshots from my radiative transfer calculation, at 5, 50 and 100 Myr (top, middle and bottom rows) after the birth of the cluster, are plotted in Figure 3.2. In the left-hand and middle columns, abundances are plotted against distance for a cluster with no HMXBs (left) and one with a HMXB phase (middle). The neutral hydrogen fraction is plotted as a function of distance from the cluster, as a solid red line. Similarly the neutral helium fraction and \( \text{He}^+ \) abundance are plotted as blue dot-dashed and green dashed lines respectively.

In the stellar-only models (left column), the hydrogen and helium ionization fronts begin by expanding rapidly into the gas. As the number of ionizing photons close to the
Figure 3.2 Heating and ionization from a model stellar cluster surrounded by a uniform density medium: including a HMXB phase. The profiles assume a uniform density medium equal to the redshift dependent average background density, for a 200 Myr simulation ending at $z = 10$. In the left-hand (no HMXB phase) and middle (HMXB phase) columns, the hydrogen neutral fraction (solid red line), helium neutral fraction (light blue dot-dashed line) and helium$^+$ abundance (green dashed line) are plotted against distance. These are shown at 5, 50 and 100 Myr after the birth of the stellar cluster in the upper, middle and lower panels respectively. The right-hand panels plot temperature against distance for both models: the stellar-only simulation (dark blue dotted line) and the model containing a further HMXB phase (fine purple dotted line).
outer edges drops, recombinations begin to stall the ionization fronts, so that their maximum extent, at 5 Myr, is $\sim 0.7$ Mpc. Since the recombination rate is fairly uniform within the fully ionized zone, the neutral fraction then gradually increases to approximately 25% at 50 Myr. By 100 Myr, the neutral fraction of the gas behind the ionization fronts has increased to $\sim 40\%$. An inner highly ionized region close to the source of radiation exists for hydrogen at both 50 and 100 Myr, though at later times it has significantly reduced in size and depth. A similar feature can be seen in the helium profile, which disappears by 100 Myr. Except in very local regions, there is very little He$^{2+}$. This gradually increases within the ionization front up to 100 Myr, where around 40% of the He$^+$ is ionized.

The inclusion of HMXBs (middle column) has very little effect on any of the ionization profiles. Close to the source of radiation ($< 100$ kpc), there is some additional ionization of hydrogen and helium at late times, but the size and depth of the main ionization fronts are identical to the purely stellar case. At early times, the harder photons from X-rays cause a very local ionized He$^{2+}$ region close the cluster, which recombines at later times.

3.3.1.2. Temperature profiles

Also plotted in Figure 3.2, in the right-hand panels, are the temperature profiles surrounding the two model clusters at three different times; 5, 50 and 100 Myr (top to bottom panels) after the birth of the cluster. The profile for a main sequence stellar-only model is shown as a dark blue dashed line, and the model with a HMXB phase is plotted with a fine dotted purple line. Within 5 Myr, the region behind the ionization front is heated to $\sim 10^4$K, and cools slowly, reaching $\sim 2500$K after 100 Myr. The temperature profiles with and without HMXBs are similar in the main ionized zone, but after 50 Myr, they diverge in both the innermost region and just outside the ionization front. HMXBs are responsible for increased heating local to the source of radiation due to hydrogen and helium photo-heating close to the source. At large distances, beyond the ionization front, a slightly extended heated region is visible, associated with the low level of He$^+$ photo-heating by hard X-ray photons at these distances.

3.3.1.3. Different spectra

The ionizing and heating potential of HMXBs is likely to depend on their spectra. HMXBs with my fiducial model spectrum, a composite blackbody and high energy power law (BB+PL), had negligible influence on their surroundings when added to the radiation already present from main-sequence stars. I now test the alternative, power-law, spectra shown in Figure 3.1, with $\alpha = 2.5, 1.5$ and 0.5. Since the power law spectra extend from 200eV to 30keV, different slopes lead to different proportions of HMXB energy falling within the useful ionizing range, from $E_{\text{min}} = 13.6$eV, to $E_{\text{max}} \approx$ a few kpc, above which
Figure 3.3  Heating and ionization from a model stellar cluster surrounded by a uniform density medium: choice of HMXB spectrum. Plots comparing the effect of my choice of HMXB spectrum on the ionization and temperature profiles surrounding a model stellar cluster embedded in a uniform density medium. Calculations were made using a one dimensional radiative transfer code, inputting the radiation from the cluster over 200 Myr, ending at z = 10. The left and right-hand panels show results at 5 and 100 Myr respectively after the birth of the cluster. Plotted as a solid red line is the output for a stellar-only simulation. Models including HMXBs are then plotted for a black body + power law spectrum (dashed green line) and three power law spectra; $\alpha = 0.5$ (dark blue dotted line), $\alpha = 1.5$ (fine purple dotted line) and $\alpha = 2.5$ (light blue dot-dashed line). From top to bottom, the panels show the hydrogen neutral fraction, helium neutral fraction, helium$^+$ abundance and gas temperature as a function of distance from the cluster.
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Photon mean free paths are too long for significant interaction with the local IGM. In fact, while 98% of the energy for $\alpha = 2.5$ falls within the 13.6-4keV energy range used in my radiative transfer code, just 84.5% and 30.9% of the HMXB energy from spectra with $\alpha = 1.5$ and 0.5 respectively fall within these bounds.

Figure 3.3 shows the resulting ionization and temperature profiles generated by re-running my one dimensional radiative transfer calculation with these different HMXB spectra. Results are shown at 5 Myr and 100 Myr on the left and right panels respectively. On the upper panels, the hydrogen neutral fraction is plotted as a function of distance for five different models; my two previous models, without (solid red line) and with (green dashed line) HMXBs, assuming a BB+PL spectrum, and my three new power-law models with spectral slopes, $\alpha$, of 0.5 (dark blue dotted line), 1.5 (fine purple dotted line) and 2.5 (light blue dot-dashed line). At 5 Myr, there is no discernible difference in the ionization state of hydrogen between the models, since stellar radiation still dominates the ionizing power. At 100 Myr there are minor differences depending on HMXB spectra within 100 kpc of the radiation source. HMXBs with softer spectra are able to keep these nearby regions marginally more ionized than those with harder spectra. On the next row, the same results are shown for helium ionization, where HMXBs have the same effect; deeper local ionized fractions for higher values of $\alpha$. He$^+$ abundance is plotted against distance on the third row. At 5 Myr, the local He$^+$ ionization is higher for softer spectra. At both 5 and 100 Myr, some singly ionized helium is present outside the main ionization front, with low levels of secondary ionization. Here, harder spectra, such the BB+PL model, and a power law with $\alpha = 0.5$, produce higher He$^{2+}$ fractions.

The temperature profiles for the different HMXB spectra are shown in the lower panels of Figure 3.3. Softer power-law spectra, with larger values of $\alpha$ increase the temperature of the inner 0.1 Mpc to a greater extent than the BB+PL model. They also have a visible effect on the temperature outside the ionization front, raising it to 150K and 400K for $\alpha = 1.5$ and $\alpha = 2.5$ respectively at 1 Mpc from the cluster after 100 Myr. Softer power-law HMXB spectra lead to higher temperatures because a larger proportion of their energy falls below $E_{\text{max}}$. HMXBs with particularly soft spectra may therefore be able to extend the volume heated by a typical starburst.

3.3.2 Heating and ionization from HMXBs within a galactic halo

In a constant density medium, a HMXB phase does not provide any significant additional heating and ionization above that released during the main sequence phase of a single starburst, unless the HMXB spectrum is particularly soft. This result predicts the effect of releasing all of the radiation from a stellar cluster directly into the IGM. Since starbursts form within galaxies, it is unlikely that all of their radiation escapes into the IGM; in a high column density galactic environment, X-rays may be able to penetrate regions which
Figure 3.4  Heating and ionization from a model stellar cluster embedded in a $10^8$ $M_\odot$ NFW halo: including a HMXB phase. The stellar cluster simulation ran for 200 Myr, ending at $z = 10$. On the left-hand (no HMXB phase) and middle (HMXB phase) panels, hydrogen neutral fraction (solid red line), helium neutral fraction (light blue dot-dashed line) and helium$^+$ abundance (green dashed line) are plotted against distance from the cluster. Panels on the upper, middle and lower rows show these profiles at 5, 50 and 100 Myr respectively. In the right hand panels, the gas temperature has been plotted against distance for the stellar-only simulation (dark blue dotted line) and the model containing HMXBs (fine purple dotted line).
UV radiation cannot. To study the effect of such a density profile, I now use a simple model with a stellar cluster placed at the centre of a galaxy whose gas follows a radially dependent NFW profile out to a fixed radius, beyond which the IGM is at mean density:

$$\rho(r) = \rho_{\text{crit}} \delta_{\text{char}} \left( \frac{r}{R_s} \right)^{\alpha} \left( 1 + \frac{r}{R_s} \right)^{-(3-\alpha)}.$$  (3.1)

Here, $R_s$ is the scale radius; $\rho_{\text{crit}}$ is the mean density of the Universe, which I take to be $2.7755 \times 10^{11} \text{M}_\odot \text{Mpc}^{-3}$; $\delta_{\text{char}}$ is the characteristic overdensity, and $\alpha$ is the inner asymptotic slope of the profile.

For X-rays to propagate further than UV radiation, the central regions need to be dense enough to block the UV, but not so dense that the X-rays are also prevented from escaping. I force this scenario using an example halo with a chosen virial mass of $10^8 \text{M}_\odot$ and a concentration of 10. These parameters define both the virial radius and the scale radius $R_s$. I model a “cored” galaxy, with $\alpha = 0$, and assume that the baryon density follows the dark matter density, with a baryon fraction of 16%. The halo extends to $\sim 6$ kpc; at greater radii, the medium returns to the constant IGM density used in the last model.

### 3.3.2.1. Ionization profiles

In Figure 3.4, I plot the hydrogen, helium and He$^+$ ionization profiles produced by a cluster without and with HMXBs (left and middle panels respectively), assuming a BB+PL HMXB spectrum. These are plotted at 5, 50 and 100 Myr after the birth of the cluster, from top to bottom rows.

In the first 5 Myr, the initial burst of UV radiation from the most massive stars in the cluster dominates its energy output. Since it is able to escape the inner regions of the galaxy, the ionization profiles at this time are the same, with and without HMXBs. In the stellar-only model, the ionized region produced by the most massive main sequence stars has fully recombined by 50 Myr, and the volume remains neutral from then onwards.

In contrast, the presence of HMXBs stalls the recombinations, and preserves the ionized region throughout the cluster’s lifetime, so that it is still 30% ionized at 100 Myr. This is because the longer mean free path of X-rays from HMXBs allows them to propagate further through their surroundings than UV energy from stars. This is shown in Figure 3.5, where the optical depth of UV and X-ray radiation to hydrogen, neutral helium and He$^+$, are plotted from top to bottom rows as a function of distance. These optical depths are calculated for radiation propagating through the IGM at 5 and 50 Myr (left and right-hand panels) after the birth of a cluster containing a HMXB phase, with an ionization state corresponding to rows 1 and 2 of the middle column of Figure 3.4. X-ray energies (blue dotted lines) are represented at 200eV, and UV energies (solid red lines) correspond...
Figure 3.5 Heating and ionization from a model stellar cluster embedded in a $10^8 M_\odot$ NFW halo: optical depths. The optical depths of UV (solid red line) and X-ray (dotted blue line) photons are plotted as a function of distance for a stellar cluster containing a HMXB phase embedded in an NFW density profile with $M_{\text{halo}} = 10^8 M_\odot$. The top panels show the optical depth of photons with the UV threshold energy (13.6eV) and X-rays (200eV) to neutral hydrogen at 5 Myr and 50 Myr in the left and right panels respectively. For reference $\tau = 1$ is shown as a green dashed line. Similarly optical depths are shown for neutral helium and He$^+$ in the middle and bottom panels, assuming UV ionization threshold energies of 24.5eV, and 54.4eV respectively, alongside X-ray energies of 200eV.
to the threshold ionization energy for the relevant elements: 13.6eV for hydrogen, 24.5eV for neutral helium, and 54.4eV for He$^+$. Additionally, $\tau = 1$ is plotted as a green dashed line for comparison. Please note that the initial rise of $\tau$ from zero is not visible in these plots because of size of the galaxy ($\sim 6\text{kpc}$) is too small to feature.

At 5 Myr, UV radiation is optically thin to hydrogen in the local IGM, but $\tau_{H_0}$ steadily increases with distance, surpassing 1 at $\sim 0.45 \text{ Mpc}$, and sharply rising to $\sim 10^3$ outside the initial ionization front. X-ray radiation, however, is optically thin to hydrogen throughout the originally ionized volume, and $\tau_{H_0}$ is only greater than 1 beyond 0.8 Mpc from the cluster. Even at early times, therefore, X-rays are able to propagate more freely through the ionized region, countering recombinations. After 50 Myr, both UV and X-ray radiation is optically thick to hydrogen. However, a small proportion of X-ray radiation ($\exp(-\tau) \sim 10^{-7}$) is able to propagate through the IGM. The optical depth of UV radiation is a factor of $10^3$ higher, so it is almost completely confined within the host galaxy. Although low, this fractional propagation of X-rays helps slow the recombination timescale for the ionized region.

Similar results can be seen for helium, although $\tau_{He^0} \gg \tau_{H_0}$. Both UV and X-ray radiation are optically thick to He$^+$ at 5 and 50 Myr. However the X-ray optical depth is a factor of $\sim 50$ lower than the UV value, so a small fraction of He$^+$ can be ionized when HMXBs are included in the starburst model.

Although X-rays enhance the lifetime of the ionized regions created by an initial burst of UV radiation, they do not extend the ionization front itself.

### 3.3.2.2. Temperature profiles

The temperature plots in the right hand panel of Figure 3.4 show identical profiles at 5 Myr, exceeding $10^4$K in the innermost regions, and dropping off sharply at the edge of the ionization front. In the stellar-only model, as the ionized fraction rapidly falls, recombinations are the dominant cooling mechanism. However the timescales for cooling are too long for thermal equilibrium to be reached. Therefore, although fully neutral within 50 Myr, the gas temperature remains at $\sim 10^4$K, for the simulation lifetime. The partially ionized IGM that remains when HMXBs are included has a considerably larger electron fraction. Consequently, Inverse Compton cooling becomes the dominant cooling mechanism, followed closely by Bremsstrahlung cooling. Therefore, although the heating rates are increased when HMXBs are included, the gas temperature decreases at a greater rate than in the stellar-only case, dropping to $5 \times 10^3$ K after 100Myr $^1$.

$^1$Note that the computed temperature is actually the electron temperature rather than the gas temperature. The temperature of neutral hydrogen, for example, could be much lower if there has been insufficient time for it to reach equilibrium via scattering
Figure 3.6  Heating and ionization from a model stellar cluster embedded in a $10^8 \, M_\odot$ NFW halo: choice of HMXB spectrum. Ionization state and temperature profiles of the gas surrounding five model stellar clusters containing (i) just stars, (solid red lines), (ii) both stars and HMXBs, assuming my fiducial blackbody and hard power law (BB+PL) HMXB spectrum (green dashed lines), and (iii-v) both stars and HMXBs, assuming a power law HMXB spectrum, with $\alpha = 0.5, 1.5$ and $2.5$ (dark blue dotted, fine purple dotted, and light blue dot-dashed lines respectively). Profiles were calculated using a one dimensional radiative transfer code whereby radiation was propagated into a medium following an NFW density profile with $M_{\text{halo}} = 10^8 \, M_\odot$. From top to bottom panels, hydrogen neutral fraction, helium neutral fraction, helium$^+$ abundance and gas temperature are plotted as a function of distance from the cluster. Panels in the left-hand column contain outputs from 5 Myr after the birth of the cluster, while those in the right-hand column plot the same for 100 Myr.
3.3.2.3. Comparison between spectra

In rows 1-3 of Figure 3.6, the ionization profiles of hydrogen, helium and He$^+$ are compared for four different HMXB spectra, at 5 Myr (left) and 100 Myr (right) after the birth of the stellar population, where the same NFW profile has been assumed. The stellar-only model is plotted as a solid red line and the BB+PL spectrum HMXB model is plotted as a green dashed line. In addition, clusters containing HMXBs with power law spectra, where $\alpha = 2.5, 1.5$ and 0.5, are plotted as dark blue dashed, purple dotted and light blue dot-dashed lines respectively. At 5 Myr there is no difference in the ionization profiles of hydrogen or helium, since the ionizing output is dominated by main sequence stellar sources. After 100 Myr, however, some differences can be seen between HMXB spectra within 0.2 Mpc from the source. Models with soft spectra marginally increase the ionized fraction of material close to the source of radiation. The different shapes of the He$^+$ profiles are comparable to those in the constant density case (c.f. Figure 3.3), with harder spectra ionizing a larger proportion of He$^{2+}$ ions.

A comparison between the temperature profiles surrounding the $10^8$ M$_\odot$ halo with different HMXB spectrum choices is shown in the bottom panels of Figure 3.6. Altering the HMXB spectra leads to broadly similar effects to those seen in the constant density case (Figure 3.3), with softer spectra producing higher local temperature peaks and increased heating beyond the main ionization front at late times.

3.4 DISCUSSION

Although I have shown that the inclusion of a HMXB phase in the evolution of a stellar cluster can extend the lifetime of its surrounding ionized volume, this effect is only visible when column densities are high enough to prevent the escape of UV radiation from the cluster. When all of the cluster’s radiation is released directly into the IGM, X-rays from HMXBs have little effect on the ionization and heating of their surroundings. In both density profile choices, HMXBs were unable to extend the size of the ionization front.

These results do not support the recent arguments in favour of HMXBs having a significant effect on their surroundings at high redshifts. X-rays from HMXBs do not “pre-ionize” distant regions or markedly raise temperatures, except in the case of particularly soft power-law HMXB spectra. In this section, I aim to explain these findings by answering the following two questions: firstly, why, in these scenarios, is there such a negligible effect from HMXBs? Secondly, under what conditions could HMXBs have an important influence on their surroundings and are these constraints physically reasonable?
3.4.1 Why do HMXBs in my model have insignificant effects on their surroundings?

I have modelled the HMXBs in my stellar population synthesis with careful comparisons to observed systems and population statistics. These include making reasoned choices for HMXB spectra, lifetimes and abundances. These parameters all influence the predicted contribution of HMXBs to the radiation emitted by the cluster. The lack of additional ionization and heating of the IGM by the inclusion of HMXBs can therefore be discussed in the context of these selections. The main reasons for the insignificant influence of HMXBs on their surroundings can be explained as follows.

3.4.1.1. HMXB radiation is dwarfed by that of its progenitors

The red lines in the upper panels of Figure 3.7 show the effect of HMXB radiation on its surroundings when all main sequence stellar radiation is removed from the simulation. Here I use a power law HMXB spectrum, normalised between 0.2-30keV, with $\alpha = 1.5$, as plotted in Figure 3.1. From left to right, the plots show the hydrogen neutral fraction, helium neutral fraction, and gas temperature as a function of comoving distance. A uniform density, with the same conditions as those in Section 3.3.1 has been assumed. The red solid lines show results at 5 Myr, and the red dashed lines show the profiles after 100 Myr. These plots show that HMXBs can locally ionize hydrogen up to $\sim 60\%$ ($\sim 90\%$), to distances of 100(400) kpc at 5 (100) Myr after the birth the cluster. They heat their local IGM to temperatures greater than $10^4$ K. The gas then cools slowly in the inner regions over time, but the temperature front moves outwards, raising the IGM temperature to 100K at $> 1$ Mpc after 100 Myr.

These results, if shown on their own, would predict a high influence of HMXBs on their surroundings. However, all HMXBs have stellar progenitors, and form with other stars. Therefore, in the context of a stellar cluster, rather than as an isolated source, HMXBs have relatively low influence. This can be seen in the red lines on the bottom panel of Figure 3.7, where the radiation of both HMXBs and main sequence stars is now included. In this case, the influence of stellar radiation, in particular the short-lived massive stellar progenitors of HMXB systems, is so high that the HMXB effect looks minimal and unimportant in comparison. This comparison warns against considering HMXBs out of context, as they would have been born into a region already fully ionized by their stellar progenitors.

3.4.1.2. HMXBs do not radiate in useful frequency ranges

Also plotted in blue in Figure 3.7 are the neutral hydrogen, neutral helium and temperature profiles of a constant density IGM, when the HMXB energy is emitted over an alternat-
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Figure 3.7 The effect of HMXB spectral energy range on their heating and ionizing potential. Upper panels: Hydrogen neutral fraction, helium neutral fraction and temperature are plotted as functions of distance from the cluster for two model stellar populations emitting only HMXB radiation (the stellar component has been removed at 5 Myr, solid lines) and 100 Myr (dashed lines). Both models assume a power law HMXB spectrum with $\alpha = 1.5$, normalised between 200 eV and 30 keV (red lines), and between 13.6 eV and 4 keV (blue lines). Lower panels: Ionization and temperature profiles from the same simulations, but with stellar radiation included.
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ive frequency range. In the upper panels, the main sequence stellar radiation has been removed. Here the HMXB energy is normalised between 13.6eV - 4keV, the range used by my radiative transfer model. In this case, a much larger proportion of HMXB energy is in the useful ionizing range. As shown in the top left panel, HMXBs with these spectra could fully ionize hydrogen to a distance of 100 kpc after 5 Myr (solid blue line). The ionized zone then expands over time, with material out to 400 kpc remaining significantly ionized at 100 Myr (blue dashed line). Similarly, the helium ionization front extends to 100 kpc at 5 Myr, and propagates further as material behind the front slowly recombines. Local temperatures are also higher with this lower frequency spectrum.

In the blue lines on the lower figures, HMXB radiation is also moved into the 13.6eV - 4keV range, but the main sequence stellar component is included. Lowering the frequency of HMXB radiation results in higher hydrogen and helium ionized fractions, particularly after 100 Myr. The ionization fronts also extend marginally further. After 100 Myr, the temperatures are also higher, within \( \sim 1 \) Mpc for this altered HMXB spectrum.

These plots highlight the importance of the spectrum choice. The fact that HMXBs emit energy in X-rays (\( > 200 \)eV) is an important factor in determining their ionizing and heating potential. For this reason they do not add much to the ionizing and heating influence of a stellar cluster, despite their high luminosities.

3.4.1.3. Ionization timescales are too long

The reason that X-rays from HMXBs have little effect on their surroundings can be seen in Figure 3.8. Here, the ionization and recombination timescales for the models in Figure 3.2 on page 76 have been calculated as:

\[
\text{Photoionization Timescale : } t_{\text{ion}} = \Gamma^{-1} \quad (3.2)
\]
\[
\text{Recombination Timescale : } t_{\text{rec}} = (\alpha n_e)^{-1} \quad (3.3)
\]

Where \( \Gamma \) is the photo-ionization rate (s\(^{-1}\)), \( \alpha \) is the recombination rate (cm\(^3\)s\(^{-1}\)) and \( n_e \) is the electron density (cm\(^{-3}\)) as defined in Appendix B4 of Bolton & Haehnelt (2007).

The left column shows the stellar-only case, while a model containing a HMXB phase with a BB+PL spectrum is plotted in the right-hand panels. These columns correspond to the left and middle panels of Figure 3.2 respectively. Ionization timescales are plotted as solid red lines and recombination timescales are plotted as blue dashed lines. These are shown at 5, 50 and 100 Myr after the birth of the cluster in the top, middle and bottom panels respectively. Also plotted, for context, as a horizontal green dotted line, is the age of the universe at \( z=10 \) (\( \approx 500 \) Myr), which was calculated for \( H_0 = 0.7, \Omega_M = 0.27 \) and \( \Omega_{\text{vac}} = 0.73 \) using the web-based cosmology calculator described in Wright (2006).

At 5 Myr, for a main-sequence stellar radiation only model (left-hand panels), the ionization timescale is extremely short (\(< 1 \) Myr) out to \( \sim 0.7 \) Mpc – the edge of the
Figure 3.8  **Ionization and recombination timescales.** Plotted in the upper, middle and lower panels are the hydrogen ionization (solid red lines) and recombination (green dashed lines) timescales at 5, 50 and 100 Myr respectively after the birth of a model stellar cluster. Results are shown for a simulation containing stellar radiation only (left panels) and a model including both stars and HMXBs (right panels) assuming a BB+PL HMXB spectrum (see Figure 3.1). These profiles were calculated using a one-dimensional time-dependant radiative transfer code over 200 Myr of the cluster’s lifetime, assuming a uniform, redshift-dependent, density profile, with a final redshift of $z = 10$. Also plotted for comparison, as a dotted green line, is the age of the universe at $z = 10$. 
fully ionized volume – before rising sharply to $> 10$ Gyr. The recombination rate, which depends on the number of free electrons, is $\sim 60$ Myr behind the ionization front and also increases rapidly at 0.7 Mpc. The timescale profiles for recombination do not change significantly in the next 100 Myr. However, 50 Myr into the cluster’s lifetime, only the inner 100kpc has an ionization timescale of $< 500$ Myr, and after 100 Myr, the front has receded to just 50 kpc. Ionization is therefore slower than recombination at almost all distances within the initially ionized region so it slowly recombines.

The right-hand profiles show the equivalent results for a cluster containing both stars and HMXBs. Within 0.7 Mpc, the ionization and recombination timescales at 5, 50 and 100 Mpc are very similar to the stellar-only case. However, at 5 Myr, during the era of the most luminous HMXBs, both the ionization and recombination timescales are reduced to $\sim 1$ Gyr at the edge of the ionization front, gradually rising at greater distances. The ionization timescale is approximately two times longer than the recombination timescale. This implies that, after $> 10^3$ Myr, an equilibrium between ionization and recombination could result in low level partial ionization of the distant IGM. However, this equilibrium cannot be reached in the lifetime of the cluster, so this solution could never form.

In Chapter 2, I solved a static version of the photoionization-recombination equations, following the qualitative reasoning used by several authors to assess the influence of X-rays on their primordial surroundings. It showed a much more notable effect of HMXBs on their surroundings than those found using a full radiative-transfer model. This is because it did not take account of ionization timescales.

The upper panels of Figure 3.9 show the hydrogen neutral fraction against time obtained from this static calculation, while the lower panels show the ionization and recombination timescales. In the upper panels, the stellar-only model is plotted in red, and a model with a HMXB phase, using a BB+PL HMXB spectrum, is shown in blue. At both 5 and 100 Myr, the addition of HMXB radiation leads to low level ionization of the distant IGM, beyond the fully ionized volume produced by the main sequence stellar radiation.

The lower panels show the ionization (solid lines) and recombination (dashed lines) timescales for the two different models. The red lines correspond to the stellar-only model, and the blue to the model comprising both stars and HMXBs. In the stellar-only case, both the photoionization and recombination timescales rise rapidly beyond the ionization front, and negligible partial ionization is visible in the upper panels. However, at late times, when HMXBs are included, both the recombination and photoionization timescales are reduced beyond the main ionized volume. Here, at 100 Myr and a distance of 0.1Mpc, the photoionization rate is only a factor of 10 higher than the recombination rate. This accounts for the partially ionized equilibrium visible in the upper panels.

Of course, all the timescales outside the ionization fronts are longer than $\sim 500$ Myr, the age of the universe at $z = 10$, which implies that the partially ionized equilibrium at these distances is unreachable. I include this discussion to highlight the limitations...
of simple calculations and general arguments used to argue in favour of alternative or additional ionizing sources. A boost in ionizing energy does not necessarily translate to increased ionization, as the spectrum of the new source needs to be considered; in the context of my isolated cluster model, X-rays may require too long a timescale to have a significant influence because the sources are short lived.

### 3.4.2 Can HMXBs have an important feedback effect?

I have shown that, in the context of my starburst model, radiation from HMXBs does not produce significant additional ionization and heating of the IGM, either in a constant or varying density medium. Indeed, in the NFW density case, including HMXBs can only prolong the lifetime of previously ionized structures, which actually results in lower IGM temperatures. I now discuss the fate of the HMXB photons in my models, and the conditions required for their feedback effects to be non-negligible.
CHAPTER 3. IONIZATION AND HEATING BY PRIMORDIAL HMXBS

3.4.2.1. Escape fractions

One of the arguments in favour of an X-ray contribution to reionization cites their long mean free paths, compared to UV photons (e.g. Mirabel et al., 2011). This means photons are able to escape further into the IGM rather than being absorbed in their immediate surroundings.

In the upper panels of Figure 3.10, I plot the escape fractions as a function of distance for my constant density IGM model, where the escape fraction is defined as

\[ f_{\text{esc}}(r) = \frac{\int_{13.6\text{eV}}^{1\text{keV}} F(E; r)dE}{\int_{13.6\text{eV}}^{1\text{keV}} F(E; 0)dE}, \]  

i.e. the proportion of ionizing energy emitted from the cluster that reaches a distance \( r \).

Five models are plotted; a stellar-only simulation (solid red line) and four models including a HMXB phase with differing HMXB spectra; a BB+PL spectrum (dashed green line) and power law spectra with \( \alpha = 2.5 \) (dark blue dashed line), 1.5 (purple dotted line) and 0.5 (light blue dot-dashed line). The escape fractions are plotted at 5 Myr and 100 Myr after the birth of a stellar cluster, in the left and right-hand panels. On the second and third rows, this escape fraction is split into a UV (13.6eV-200eV) and X-ray (200eV-1keV) component respectively.

At 5 Myr, the total escape fractions drop to below 20% beyond the ionization front for all models, as radiation from main sequence stars dominates the cluster SED. In the middle panel, it is clear that all of the UV energy is absorbed within this region. However, the bottom panel shows that 90% of the X-ray energy escapes for all models including HMXBs, with more energy escaping when the HMXB spectrum is harder. Note that the X-ray radiation from main sequence stars is negligible, so its divergence from the other lines in the bottom panel is unimportant.

After 100 Myr, radiation from HMXBs dominates the ionizing emission from the cluster, so more pronounced differences are seen in the total escape fractions between models with and without a HMXB phase. In the stellar-only model, the escape fraction decreases sharply with distance from the source, with all energy absorbed within 0.2 Mpc. In the models involving a HMXB phase, however, most of the energy escapes, because the SED is much harder. The UV escape fraction is greater for harder spectra, and the X-ray escape fraction remains similar to the 5 Myr snapshot, with > 80% of the X-rays escaping beyond 1 Mpc in all models including a HMXB phase.

These escape fraction plots show that, as theory suggests, a significant proportion of HMXB energy penetrates deep into the IGM. It does not make an impression on the ionization state of the material at these distances because of long ionization timescales. However, should the X-ray background build up to greater levels, these timescales could be reduced and an effect may be visible.
Figure 3.10 Heating and ionization from a model stellar cluster surrounded by a uniform density medium: escape fractions. The escape fraction – the fraction of energy in a given frequency range that remains unabsorbed – is plotted as a function of distance from the cluster for different stellar population models. A constant redshift-dependent density profile has been assumed, and the calculations have been made using a one dimensional radiative transfer code, over a cluster lifetime of 200 Myr, ending at \( z = 10 \). The upper panels show the total 13.6eV-4000keV escape fraction for five different models at 5 Myr (left panel) and 100 Myr (right panel) after the birth of the cluster. The models plotted are (i) stellar radiation only (solid red line), (ii) stars and HMXBs, with a blackbody + hard power law HMXB spectrum (our fiducial model; green dashed line) and (iii-iv) power law models with \( \alpha = 0.5, 1.5 \) and 2.5 (dark blue dotted, fine purple dotted and light blue dot-dashed lines respectively). The panels on the second and third rows split this escape fraction into UV (13.6-200eV) and X-rays (200-4000eV) respectively.
CHAPTER 3. IONIZATION AND HEATING BY PRIMORDIAL HMXBS

3.4.2.2. Greater HMXB number densities

I have established that the ionization timescales for the X-rays from HMXBs are too long for them to exert a major effect on the ionization state of their distant surroundings, despite their high escape fractions. A higher number density of HMXBs might lower these timescales. This could be achieved in three ways;

(i) increasing the HMXB to main sequence star ratio,

(ii) modelling a larger starburst, so that more HMXBs can form,

(iii) continuous star formation, allowing for more uniform X-ray production.

My estimates of the proportion of stellar binaries that survive the primary’s supernova phase to become HMXBs (30%) are observationally motivated. Additionally, I have assumed $f_{\text{sur}} = 1$ throughout this paper, thus neglecting any other evolutionary reasons for HMXBs failing to form. Therefore, increasing the ratio of HMXBs to main sequence stars, as in option (i), would be implausible within the context of my simulations.

In Figure 3.11, I investigate (ii) by increasing the size of my stellar clusters from $10^6$ to $10^{10}$ stars. Here, the total energy from the cluster with and without a HMXB phase is propagated through a constant density medium, as described in Section 3.3.1. From top to bottom, hydrogen neutral fraction, helium neutral fraction, $\text{He}^+$ abundance, temperature and ionization timescale are plotted against distance for a stellar-only model (red lines) and a stellar and HMXB model, with a BB+PL HMXB spectrum and $f_{\text{sur}} = 1$ (blue lines). These are plotted at 5 Myr after the birth of the cluster (solid lines), and 100 Myr (dashed lines). Four different cluster sizes are shown from left to right, containing $10^7$, $10^8$, $10^9$ and $10^{10}$ stars respectively. Note that the horizontal scale is increased for more massive clusters, to incorporate the larger ionized volumes.

Since a greater X-ray background density results in shorter ionization timescales, the relative influence of stars and HMXBs on their surroundings is non-linear. Therefore, although stars dominate the heating and ionizing influence of the cluster on the IGM, HMXBs show more significant effects as the cluster size increases. In a $10^7$ star cluster, the hydrogen ionization profiles (top left plot) show little difference at 5 Myr. However, although the ionization front is not extended at this time, after 100 Myr, a partially ionized zone remains that is marginally larger than the originally ionized region. There is no matching increase in the volume of ionized helium, but recombinations are also quenched by HMXB radiation. As the number of stars in the cluster increases, towards the right-hand panel of Figure 10, the effect of HMXBs on their environment increases relative to that of the main sequence stellar component. In particular, both the depth and extent of the partially ionized regions after 100 Myr increase. However, at no point is the outer edge of the ionization front extended because the ionization timescales (bottom panels), while much reduced by the inclusion of HMXBs, remain prohibitively long.
Figure 3.11 Heating and ionization from larger stellar clusters in a uniform density medium: boosting HMXB ionization and heating power. From left to right, clusters contain $10^7$, $10^8$, $10^9$ and $10^{10}$ stars. Their combined emission is propagated through a constant density medium, for 200 Myr, ending at $z = 10$. From top to bottom, the hydrogen neutral fraction, helium neutral fraction, helium$^+$ abundance, gas temperatures and ionization timescales have been plotted as a function of distance from the clusters for two models; without (red lines) and with (blue lines) HMXBs. Solid lines show the profiles at 5 Myr after the birth of the cluster, and dashed lines show the profiles after 100 Myr.
CHAPTER 3. IONIZATION AND HEATING BY PRIMORDIAL HMXBS

The temperature at 5 Myr is dominated by stellar radiation for all cluster sizes. However, the effects of HMXBs are increasingly significant as cluster size increases. Close to the cluster, and beyond the ionization front, a HMXB phase results in higher temperatures. Yet, within the main ionized region, the higher cooling rate of the partially ionized regions means that including HMXBs results in lower temperatures, as seen in the NFW density results of Figure 3.4.

In my model, I treat star formation as a single burst over a short timescale. Although this may be relevant for e.g. globular clusters, star formation is likely to be a continuous process within young galaxies. HMXBs are unable to sustain a sufficient ionizing background for significant ionization and heating under these conditions, due to their brief lifetimes. However, ongoing HMXB formation may be able to maintain a low level X-ray ionizing background, pre-ionizing the distant IGM before UV radiation from main sequence stellar sources penetrates these distances. Quantification of these effects is beyond the scope of my starburst model.

3.4.3 Mechanical feedback and local effects

I note that HMXBs may be able to exert influence on the IGM via alternative mechanisms that have not been considered in this study. For example, several HMXBs are known to drive large scale outflows (e.g. SS 433; Blundell et al., 2001). These alternative feedback processes are addressed in Justham & Schawinski (2012).

Furthermore, while this study may rule out any strong feedback effect of HMXBs on the IGM, I have not addressed the radiative and mechanical effects that HMXBs may exert on their host galaxies. Significant feedback within these dense environments may be plausible.

3.5 CONCLUSIONS

I have investigated the influence of HMXBs on the high redshift IGM, using my model of a coeval stellar population, comprising main sequence stars formed in a single burst followed by a later HMXB phase. I argue that my model provides a plausible representation of HMXBs in the early universe, due to observationally motivated choices for HMXB spectra, abundance and lifetimes. Using a one dimensional radiative transfer code, I have predicted the ionization and temperature profiles surrounding the population. I have considered the interaction of the stellar radiation from my model cluster through a uniform density IGM from $z = 14.5$ to $z = 10$, and through a $10^8 M_\odot$ one-dimensional NFW profile, with and without a HMXB phase.

For a constant density IGM, HMXBs produce negligible enhanced ionization, except in the case of extreme HMXB spectra. This is because their high energy SEDs and brief
lifetimes leave insufficient time for an X-ray ionized background to build up. Radiation from massive main sequence stars raises the temperature to $10^4$ K, but including a HMXB phase only increases the heating rate in the inner $\sim 0.1$ Mpc. In the NFW profile case, UV radiation from the most massive stars initially ionizes a large IGM volume, which later recombines as the reduced stellar radiation is unable to penetrate the high density galactic core. HMXB photons stall recombinations behind the front, keeping it partially ionized for longer, but do not extend the outer edge of the ionized zone. The partially ionized regions have an increased number of free electrons, so enhance both Inverse Compton and Bremsstrahlung cooling. Consequently, HMXBs actually increase the cooling rate in these regions, halving the gas temperature after 100 Myr. In neither case is thermal equilibrium reached during the simulation’s lifetime, so cooling is ongoing.

Although I have shown that HMXBs do not have a strong influence on the IGM in the context of my starburst model, they may nevertheless have a significant local influence within their host galaxies. Additionally, X-rays from HMXBs had a high escape fraction in both density cases. Consequently, ongoing HMXB formation may be able to maintain a low level X-ray ionizing background, which could partially ionize the distant IGM.
Short Period X-ray Binaries: the Transition to Radiatively Inefficient Accretion

Madman, thou errest. I say, there is no darkness but ignorance, in which thou art more puzzled than the Egyptians in their fog.

William Shakespeare, Twelfth Night
In this chapter, the focus of my thesis shifts from a hypothetical population of X-ray binaries in the early universe to the black hole population in our own Galaxy. I investigate the nature of the transition between radiatively efficient and inefficient accretion flows, and the effects this may have on the detectability of the lowest luminosity black hole accretors in our Galaxy.

4.1 INTRODUCTION

Low Mass X-ray Binaries (LMXBs) produce X-rays through the accretion of matter onto their primaries (see reviews by e.g. Tanaka & Lewin, 1995; van Paradijs & McClintock, 1995; McClintock & Remillard, 2006; Done et al., 2007). All black hole and many neutron star LMXBs are transient, experiencing occasional outbursts during which their mass accretion rate and X-ray luminosity increase by several orders of magnitude. These outbursts have typical durations spanning days to months.

In Section 1.3.2, I showed that the lightcurves of LMXB outbursts decay exponentially (e.g. Tanaka & Shibazaki, 1996; Chen et al., 1997), unless the disc is entirely irradiated (see Section 4.3.2). Their mass accretion rate was characterised by King & Ritter (1998) as:

$$
\dot{M} \approx R_D \nu \rho \exp \left( -\frac{3\nu t}{R_D^2} \right) .
$$

Here $\rho$ is the disc density and $\nu = \alpha_h c_s H$ is the viscosity, where $\alpha_h \sim 0.1$ is the hot state viscosity parameter (Shakura & Sunyaev, 1973), $c_s$ is the sound speed, and $H$ is the vertical scale height of the disc. The maximum extent of the disc, $R_D$, is limited by the size of the Roche Lobe of the primary and can be related to the orbital period ($P_{\text{orb}}$). For typical mass ratios ($q = M_2/M_1 \leq 0.1$), $R_D \approx 1.77 \times 10^{10} m_1^{1/3} P_{\text{orb}}^{2/3} (h) \text{ cm}$, where $P_{\text{orb}}(h)$ is the orbital period in hours and $m_1$ is the primary mass in solar masses. Both the peak accretion rate ($\dot{M}_{\text{max}} \sim R_D \nu \rho$) and outburst duration ($t_o \sim M_D(R_D) / \dot{M}_{\text{max}}$, where $M_D$ is the disc mass; c.f. Section 4.3) depend inherently on the orbital period.

The radiative efficiency of accretion in an LMXB is given by $L = \eta \dot{M} c^2$, where $\eta \sim 0.1$ for a radiatively efficient flow through a thin accretion disc (see e.g., Frank et al., 2002). At low accretion rates, cooling becomes inefficient, and a radiatively inefficient advection dominated accretion flow (ADAF) can occur (Ichimaru, 1977; Rees et al., 1982; Narayan & Yi, 1995b), which is described in Section 1.2.4. In this model, the accreting gas has a very low density, leading to an optically thin flow that cannot cool efficiently within an accretion time. The viscous energy is stored in the gas as thermal energy rather than being radiated away, and is advected onto the central compact object. The transition from radiatively efficient to inefficient flow is expected to take place once the accretion luminosity reaches a few percent of the Eddington luminosity ($L_{\text{Edd}}$). For black hole primaries, which lack a hard surface, $\eta \rightarrow 0$ and the accretion energy is carried with
the mass flow into the hole or transferred elsewhere, e.g. as radio jets or mechanical outflows. Neutron star primaries do not experience a drop in accretion efficiency because the advected energy must always be radiated from the stellar surface. Hence, at sufficiently short periods (and thus small disc radii and peak accretion rates), black hole LMXBs will undergo fainter, shorter outbursts than comparable neutron star systems, making them more difficult to detect.

The relationship between orbital period and peak outburst luminosity for LMXBs is well established (e.g. Shahbaz et al., 1998; Portegies Zwart et al., 2004). Wu et al. (2010) studied the outburst luminosities of a sample of transient LMXBs observed by the Rossi X-ray Timing Explorer (RXTE). They showed that there was no discernible difference in the orbital period-peak outburst luminosity relation of black holes compared to neutron star primaries, when luminosities were measured in Eddington units. However, they suggest that the two populations may diverge at short periods. Along similar lines, Meyer-Hofmeister (2004) note that the low peak outburst luminosities in short period black hole LMXBs can cause them to remain in a low luminosity – hard spectral state, rather than entering the high luminosity – soft state expected for radiatively efficient accretion. This breed of outburst, where sources remain in the low-hard state throughout, has been observed in several sources (Brocksopp et al., 2001, 2004).

In addition to establishing a relation between orbital period and peak outburst luminosity, Wu et al. (2010) note that the absence, in their dataset, of black holes with orbital periods of \( \leq 4 \) hr, may be caused by radiatively inefficient accretion lowering the peak outburst luminosities of these systems. In this chapter, I investigate this hypothesis, using the Ritter-Kolb catalogue (Ritter & Kolb, 2003) to show that there is statistical evidence for a dearth of black holes in LMXBs at short orbital periods. I suggest that this is caused by the increasing importance of radiatively inefficient accretion in black hole systems lowering not only peak outburst luminosities but also outburst durations and X-ray duty cycles. Additionally, I investigate the nature of the transition to radiatively inefficient accretion, modelling it as an instantaneous change to \( \eta = 0 \) at a fraction \( f \) of the Eddington luminosity \( L_{\text{Edd}} \) and as a power law decrease, \( \eta \propto \dot{M}^n \), below \( fL_{\text{Edd}} \). In Section 4.2, I present the sample of systems and the statistical evidence for the lack of black hole LMXBs at short orbital periods. In Section 4.3, I study the effect of a radiative efficiency switch on the peak luminosities and outburst timescales of black hole LMXBs. I determine the conditions under which these effects can hide a short orbital period black hole population in Section 4.4. This is followed by a discussion of my findings and my conclusions.

### 4.2 ORBITAL PERIOD DISTRIBUTION ANALYSIS

If a switch to inefficient accretion in black hole LMXBs at low luminosities reduces their outburst luminosities and timescales, it follows that black holes will be more difficult to
Figure 4.1  Orbital period distributions of neutron star and black hole LMXBs. Samples 1, 2 and 3 as described in text. The vertical dashed line separates sources with periods above and below 80 minutes.

Table 4.1  Comparison between black hole and neutron star LMXB orbital periods using A-D statistics: In the left hand column, orbital periods of the entire black hole sample (18 sources) is compared to those of all NS LMXBs (51 sources) in row 1, NS LMXBs with orbital periods greater than 80 mins (40 systems) in row 2, all transient NS LMXBs (25 systems) in row 3, and all transient NS LMXBs with orbital periods greater than 80 mins (21 systems) in row 4. The other columns contain the same comparisons but all NS and BH with orbital periods above and below 0.5 days are removed from the middle and right column samples respectively. Results with a $p$-value below 0.01 (high confidence) are highlighted in bold.

<table>
<thead>
<tr>
<th>Black Holes compared to:</th>
<th>A-D Probability (all)</th>
<th>$P_{\text{orb}} &lt; 0.5 \text{ d}$</th>
<th>$P_{\text{orb}} &gt; 0.5 \text{ d}$</th>
</tr>
</thead>
<tbody>
<tr>
<td>all NS-LMXBs</td>
<td>$1.2 \times 10^{-3}$</td>
<td>$8.7 \times 10^{-4}$</td>
<td>0.17</td>
</tr>
<tr>
<td>all NS-LMXBs: $P_{\text{orb}} &gt; 80 \text{ min}$</td>
<td>0.016</td>
<td>$7.2 \times 10^{-3}$</td>
<td>0.17</td>
</tr>
<tr>
<td>Transient NS-LMXBs</td>
<td>$1.8 \times 10^{-3}$</td>
<td>$5.5 \times 10^{-4}$</td>
<td>0.29</td>
</tr>
<tr>
<td>Transient NS-LMXBs: $P_{\text{orb}} &gt; 80 \text{ min}$</td>
<td>$9.6 \times 10^{-3}$</td>
<td>0.020</td>
<td>0.29</td>
</tr>
</tbody>
</table>
detect than neutron stars with comparable accretion rates. Such an effect may be observed as a divergence in the populations of black hole systems compared to those of neutron stars at short orbital periods. I searched for this within the catalogue of cataclysmic variables, LMXBs and related objects compiled by Ritter and Kolb (Ritter & Kolb, 2003), by selecting 3 LMXB samples, defined as:

- Sample 1: all black hole LMXBs, excluding any unconfirmed candidate systems; (17 systems)
- Sample 2: all neutron star LMXBs, both transient and persistent, but excluding any globular cluster systems (which are likely to have very different evolutionary histories to Galactic binaries); (51 systems)
- Sample 3: only the transient neutron star LMXBs from Sample 2. (25 systems)

In all samples, extragalactic sources were removed. In addition, I removed the following systems which are mislabelled in the catalogue:

(a) Globular cluster sources (not labeled GC in catalogue): J1748-2446 (Papitto et al., 2011), J1748-2021 #1 (Altamirano et al., 2008), J1748-2021 #2 (Altamirano et al., 2010), NGC 104-X-5 (Edmonds et al., 2002), J1623-2631 (Kaluzny et al., 2012), NGC 104-X7 (Edmonds et al., 2002), J1910-5959 #1 (Kaluzny & Thompson, 2009) and NGC 104-W37 (Heinke et al., 2005)

(b) Black hole candidates (not labeled BH? in catalogue): V1408 (Russell et al., 2010), V4134 Sgr (Kaaret et al., 2006) J1242+3232 (Carpano et al., 2007), J1752-0127 (Durant et al., 2009)

A full list of the sources used in these samples can be found at the end of this chapter in Table 4.2.

Before comparing the black hole and neutron star LMXB populations, I checked whether there was a difference in the orbital period distribution of transient and persistent neutron star systems. For an irradiated disk, a source is persistent if \( \dot{M}_{\text{ext}} > \dot{M}_{\text{crit}} \), where \( \dot{M}_{\text{ext}} = -\dot{M}_2 \) is the mass accretion rate at the edge of the disc and

\[
\dot{M}_{\text{crit}} = 9.5 \times 10^{14} \xi^{-0.36} \alpha_{0.1}^{0.04+0.01 \log \xi^{-3}} \int_0^{2.39-0.1 \log \xi^{-3}} R_{10}^{-0.64+0.08 \log \xi^{-3}} g \, \text{s}^{-1}
\]

from Coriat et al. (2012) with an irradiation constant \( \xi = 10^{-3} \xi^{-3} \). For \( \xi = 10^{-3} \), \( \dot{M}_{\text{crit}} \propto R_{10}^{2.38} \propto P_{\text{orb}}^{1.6} \). This proportionality implies a difference in the orbital periods of transient and persistent LMXBs, with more persistent systems at shorter orbital periods. However, an Anderson-Darling test of Samples 2 and 3 gives a 0.60 probability that they are the same, indicating that both are drawn from the same distribution.
CHAPTER 4. SHORT PERIOD X-RAY BINARIES

The presence of both persistent and transient systems at higher orbital periods can be reasoned by the following two arguments. Firstly, for $P_{\text{orb}} > 2$ hours, $M_{\text{ext}} \propto P_{\text{orb}}^{5/3}$, from (1.7), meaning $M_{\text{ext}}$ and $M_{\text{crit}}$ are increasing at very similar rates with orbital period. Consequently, the considerable scatter in $M_{\text{ext}}$ for observed systems is likely to lead to the presence of both persistent and transient systems at all orbital periods. Secondly, at large orbital periods, persistent sources with $M_{\text{ext}} > M_{\text{crit}}$ will be among the brightest LMXBs and so may be more likely to be observed, despite the increase in proportion of transient systems.

Although there is no significant difference in the period distribution of transient and persistent LMXBs, I compare both Samples 2 and 3 with the black hole population as a control to test for sampling biases caused by the lack of persistent black hole sources.

In the upper panel of Figure 4.1, I plot the normalised histogram of the log$P_{\text{orb}}$ distribution of samples 1, 2 and 3. There is no discernible difference between the two neutron star samples, but the black hole sample shows a distinct lack of sources with orbital periods below $\sim 0.1$ days. There are in fact 25 (12) neutron star LMXBs in sample 2 (3) with orbital periods below the minimum black hole period of 0.17 days (4.1 hr). The cumulative histogram in the bottom panel of Figure 4.1 further highlights this difference.

The vertical dashed line in Figure 4.1 separates sources with orbital periods above and below 80 minutes. Below this (conservative) limit, the mass-period relation requires that these systems have either sub-stellar or compact secondaries (King, 1988). I compare the black hole sample to the neutron star sample including and excluding sources below this limit, as a check for differences caused by system evolution.

I use the Anderson-Darling test to determine whether the populations of the black hole and neutron star LMXBs are different. The A-D test is similar in purpose to the more familiar Kolmogorov-Smirnov test, and can be used to test whether two samples are consistent with being drawn from the same parent distribution. However, the A-D has better sensitivity to differences between distributions, especially near their tails, compared with the K-S test (Feigelson & Babu, 2012). The A-D test results are shown in Table 4.1. The 4 rows represent comparisons of the black hole sample with (a) sample 2, (b) sample 2 for $P_{\text{orb}} > 80$ min, (c) sample 3 and (d) sample 3 for $P_{\text{orb}} > 80$ min. The A-D $p$-values are shown in the left hand columns; these correspond to the probability of observing a difference as (or more) extreme than actually observed, on the assumption the two parent populations are the same. In addition to the $p$-values for the full samples, I show those for sources with orbital periods above and below $P_{\text{orb}} = 0.5$ days; this breaks the black hole population roughly in half (8 and 9 sources respectively). Results with $p < 0.01$ are highlighted in bold.

These results suggest that all four neutron star samples have different orbital period distributions to the black hole sample across the full period range. This holds when I compare systems with $P_{\text{orb}} < 0.5$ days. However, if I remove all sources with periods <
0.5 days, I find no significant difference in the period distributions. From this it can be concluded that the black hole and neutron star samples differ at the short end of the orbital period distribution.

I tested the robustness of this conclusion by adding/removing a few short period systems from each sample. In order to reduce the difference between the samples to the point where the A-D test reports no statistically significant result (i.e. \( p > 0.05 \)), one would have to remove the eight shortest period NS systems (from the NS transient sample) or include four additional BH systems with periods < 0.1d. Consequently, the make up of the samples would have to change quite considerably in order to influence this result.

### 4.3 Predicted Effects of a Transition to Radiatively Inefficient Accretion

At high luminosities, a radiative efficiency of \( \eta = 0.1 \) is a reasonable approximation to black hole and neutron star LMXB outbursts. However, below a fraction \( f \) of \( L_{\text{Edd}} \) (typically a few percent; Abramowicz et al., 1995; Maccarone, 2003), accretion onto black holes becomes radiatively inefficient and \( \eta \to 0 \). The efficacy of this switch to radiatively inefficient accretion, in hiding a population of short period black hole LMXBs, depends on the nature of the transition. Here, I investigate two possible forms:

**A** a sharp, instantaneous switch to \( \eta = 0 \) at \( L \leq f L_{\text{Edd}} \),

**B** a gradual reduction in \( \eta \), where below \( L \leq f L_{\text{Edd}} \):

\[
\eta = 0.1 \left( \frac{\dot{M}}{f M_{\text{Edd}}} \right)^n
\]

(4.3)

based on theory by Narayan & Yi (1995b). I use \( n = 1 \) in subsequent sections, as used in Coriat et al. (2012), but discuss larger values of \( n \) in Section 4.5.2.

A reduction in \( \eta \) at low luminosities will reduce the observed duration of an outburst and, in the case of short orbital period systems, may render the entire outburst undetectable.

#### 4.3.1 Peak Luminosity

In the early stages of LMXB outbursts, the accretion rate is known to decay exponentially, following (4.1). I approximate the disc radius as \( R_D \approx 0.7 R_{L1} \), where \( R_{L1} \) is the Roche radius of the primary:

\[
\frac{R_{L1}}{a} = \frac{0.46q^{-2/3}}{0.6q^{-2/3} + \ln(1 + q^{-1/3})}
\]

(4.4)
(Eggleton, 1983), and \( a \) is the binary separation, given by:

\[
a = 3.53 \times 10^{10} \frac{m_1^{1/3}}{(1 + q)^{1/3}} P_{\text{orb}}^{2/3} \text{ cm.} \tag{4.5}
\]

The secondary mass is fixed at 0.4\( M_\odot \), based on typical values in the Ritter-Kolb catalogue (see Section 4.4.3.2 for further discussion of secondary masses). Following King & Ritter (1998), I take \( \rho \) to be \( \sim 10^{-8} \text{ g cm}^{-3} \). They show that \( \rho \) is independent of radius, meaning that this value is suitable for short period systems. For neutron stars, I use a value of \( 2 \times 10^{-8} \) as \( \rho \) depends on the surface density (see equation 4.10) which is proportional to \( M^{-0.35} \).

The viscosity is taken as \( \nu = \alpha_k c_s H = \alpha_k c_s (H/R) R_D \) where \( \alpha_k \sim 0.1, c_s \approx 10 \sqrt{(T/10^4 \text{ K})} \) km s\(^{-1} \) with \( T \), the hydrogen ionisation temperature, \( \sim 6500 \text{K} \). The alpha prescription above suggests that \( \nu \) is linearly proportional to \( R_D \). However, observationally, there is significant uncertainty in its value and its dependence on \( R_D \). King & Ritter (1998) choose a value of \( \nu = 10^{15} \text{ cm}^2 \text{s}^{-1} \) for a disc radius of \( 1 \times 10^{11} \text{ cm} \). At orbital periods of \( \sim 0.1 \) days, \( R_D \sim 5 \times 10^{10} \text{ cm} \). Adopting this as a typical radius, I use \( \nu = 5 \times 10^{14} \text{ cm}^2 \text{s}^{-1} \). This choice, while consistent with earlier work, also yields peak luminosities in reasonable agreement with the observational results in Wu et al. (2010): \( L_{\text{peak}} \sim (\text{a few} \times 10^{36}) \) erg s\(^{-1} \) for periods of a few hours.

At \( t = 0 \), the mass accretion rate, \( \dot{M}_{\text{peak}} = \rho \nu R_D \), so the peak luminosity is \( L_{\text{bol,peak}} = \eta c^2 \rho \nu R_D \). Since I wish to compare our estimates to X-ray observations, I convert this bolometric luminosity into the 2-10 keV value by introducing a correction factor, \( f_{\text{corr}} \), such that \( L_{\text{peak}} = L_{\text{bol,peak}} / f_{\text{corr}} \). For both neutron star and black hole systems, I calculated a simple spectrum from an accretion disc of specified inner/peak temperature absorbed by a column density of \( 5 \times 10^{21} \text{ cm}^{-2} \); a typical level of Galactic absorption to an X-ray binary. I used this to compute the ratio of bolometric flux (0.002-30 keV), after removing the effect of the absorption, to the flux in the 2-10 keV band including absorption. For this calculation I used XSPEC v12.8 (Arnaud, 1996). In the neutron star case, I used a TBabs absorption model (Wilms et al., 2000), whereas I modelled the black hole spectrum with a discBB model, since, in outburst, its spectrum is dominated by the disc and needs to be treated as a multi-temperature blackbody. A black hole LMXB in outburst has an inner disc temperature of 0.5-1.0 keV, giving \( f_{\text{corr}} \sim 8.5 - 2.5 \). For neutron star LMXBs, with higher disc temperatures and thermal emission from their surface and boundary layer, \( f_{\text{corr}} \sim 1.3 - 1.2 \). Based on these estimates, I use a correction factor of 4 for black holes (corresponding to a temperature of \( \sim 0.7 \text{ keV} \)), and 1.3 for neutron stars.

The uppermost panels of Figure 4.2 show the peak luminosities \( (L_{\text{peak}}) \) as a function of orbital period for black holes of mass \( M_1 = 8M_\odot \) using three different values of \( f \): 0.01, 0.03 and 0.05, plotted as dark blue dotted, purple dotted and light blue dot-dashed lines respectively. Both a neutron star (solid red line) and a black hole system (green dashed
Figure 4.2 Outburst properties of LMXBs under Case A and Case B transitions to inefficient accretion. Peak 2-10 keV outburst luminosity (upper panels), observable outburst timescale (middle panels) and X-ray duty cycles (lower panels) vs orbital period. The dark blue (thick dotted), purple (fine-dotted) and light blue (dot-dashed) lines show black hole LMXBs with a primary mass of $8 M_\odot$ and radiative efficiency switch fractions $f = 0.01$, 0.03 and 0.05 respectively. Also plotted (solid red line) is a neutron star system with $M_1 = 1.4 M_\odot$ and (dashed green line) a $8 M_\odot$ black hole system with no switch to radiative inefficiency for comparison. The case A transition (a sharp switch to inefficiency, described in the text) is shown in the left hand panels while the case B (smooth) switch is shown on the right. The observable timescales are defined using a luminosity of $L \sim 10^{36}$ erg/s to signify the "end" of the outburst. This is calculated from the limiting flux, 10 mCrab, with a source distance of 8 kpc. The timescales in the middle panels are replotted for three different distances in Figure 4.3.

Case A, where $\eta$ drops sharply to 0 at $\dot{M} \leq f \dot{M}_{\text{Edd}}$, is plotted in the left panel. Here, $f \sim 0.03$ produces a cut-off in the peak outburst luminosity at orbital periods below $\sim 0.1$ days, matching the observed cut-off of low-orbital period black hole systems in the observational data. The gradient before the cut-off is $2/3$, since $L_{\text{peak}} \propto M_{\text{max}} \propto R_D \propto P_{\text{orb}}^{-2/3}$. The gradients are consistent with the observed correlation in Wu et al. (2010). The different normalisations of the neutron star and black hole $L_{\text{peak}} - P_{\text{orb}}$ relations are due to alternate choices of $f_{\text{corr}}$ for the two types of system.

Case B, with a power-law decrease in $\eta$ below $fL_{\text{Edd}}$, with $n = 1$, is plotted in the

---

Footnote: Farr et al. (2011) found that wider distributions of black hole mass were also consistent with data; this would need to be taken into account if a range of masses were used.
right-hand panel. In this case, the peak luminosity-orbital period relation steepens at short orbital periods when the transition to radiatively inefficient accretion sets in. Here, the factor of $\dot{M}$ in $\eta$ means that $L_{\text{peak}} \propto \dot{M}^2 \propto R_D^2 \propto P_{\text{orb}}^{4/3}$. While a gradual switch to radiatively inefficient accretion causes the peak luminosities of black hole systems to fall more rapidly at short periods, it does not produce large differences between the two types of system at periods $\sim 0.1$ days, unless $f > 0.05$. In this case, the black hole and neutron star gradients diverge below 0.3 days. Observations show a similar gradient and normalisation of peak luminosities down to $\sim 0.1$ days, although the predicted effect could be hidden.

Since these simple analytic models only track emission from the disc, note that the sharp change in luminosity at $fL_{\text{Edd}}$ does not necessarily produce an equally strong observable signature. Sources with peak luminosities below the efficiency threshold are likely to exhibit low-hard outbursts, with power-law spectra that are not dominated by the disc blackbody. In this case, my choice of spectrum, and hence, $f_{\text{corr}}$ is no longer valid and it is likely to be larger by factor of $\sim 5$ (see Maccarone, 2003). This change in $f_{\text{corr}}$ may go some way towards shielding the drop in disc luminosity from the overall bolometric luminosity, which does not change dramatically as the disc moves into a hard state (Zhang et al., 1997; Belloni et al., 2006).

### 4.3.2 Outburst timescales

From a time $T$ onwards, after the initial exponential fall off, King & Ritter (1998) show that the mass infall rate follows a linear decay, obeying:

$$\dot{M} = \left(\frac{3v}{B_m}\right)^{1/2} \left[M_{h}(T) - \left(\frac{3v}{B_n}\right)^{1/2} (t - T)\right]$$

(4.6)

where $T$ is the time taken for the irradiated radius to drop below the disc radius, and is given by:

$$T = \frac{R_D^2}{3v} \log \left(\frac{B_n\nu \rho}{R_D}\right)$$

(4.7)

and $M_h(T)$ is the irradiated mass at time $T$:

$$M_h(t) = \rhoR_d^3 \exp\left(-\frac{3vt}{R_D^2}\right).$$

(4.8)

$B_m$ is defined by $R_h^2 = B_m\dot{M}$ where $R_h$ is the irradiation radius of the disc. The value of $B_m$ depends on $m$: 1 for neutron stars with a hard surface and 2 for black holes without. In these simple calculations, I use a value of $B_m = 10^5$ for both systems. I note that the uncertainties in my chosen parameters affect the exact results of these calculations, but not the main conclusions of my study; these are addressed in Section 4.5.1.
Figure 4.3  Observable timescales at three distances from the source. Observable timescales ($t_{\text{det}}$) are plotted for a black hole LMXB (green dashed line) and a neutron star LMXB (solid red line) as a function of orbital period at three different distances from the source. In the upper panel, the distance is 5 kpc, meaning the outbursts become undetectable at a limiting flux of $7 \times 10^{35} \text{erg s}^{-1}$. For the middle panel, at a distance of 10 kpc, the limiting flux is $3 \times 10^{36} \text{erg s}^{-1}$, and in the bottom panel the plots are made for a distance of 15 kpc which corresponds to a limiting flux of $6 \times 10^{36} \text{erg s}^{-1}$.
The outburst ends when $\dot{M} = 0$; with $t_o = T + t_{\text{visc}}/3$. However, observations are flux limited, and therefore the outburst appears to end earlier, once its signal drops below a limiting flux, $F_{\text{lim}}$. I define the observable outburst duration, $t_{\text{det}}$, to be the time at which the luminosity drops below $L_{\text{lim}} = 4\pi d^2 F_{\text{lim}}$. Since the stellar density is highest at the Galactic centre, it is reasonable to assume that this is the location of the highest concentration of LMXBs, so that $d \simeq 8$ kpc. Using the daily exposure sensitivity of the RXTE, I define $F_{\text{lim}}$ as 10 mCrab, and calculate $L_{\text{lim}}$ to be $\sim 10^{36}$ ergs$^{-1}$. My use of the RXTE survey to define $F_{\text{lim}}$ is explained in Section 4.4.1.

In the middle panels of Figure 4.2, I plot the predicted observable outburst durations. For Case A, outbursts are terminated abruptly when the luminosity falls below either $L_{\text{lim}}$ or $f L_{\text{Edd}}$. As well as the sharp cut-offs corresponding to the peak luminosity graphs in the top panels, note that timescales diverge from the expected duration by an increasing factor as the orbital period drops.

For Case B transitions, a decrease in outburst duration is only evident at orbital periods $\lesssim 0.5$ days. Here, a more gradual divergence from the expected outburst timescales occurs. At the minimum observed black hole orbital period of $\sim 0.1$ days, the expected duration of a black hole outburst is approximately halved for $f = 0.03$. However, it is still longer than the duration of a neutron star outburst, and should not significantly affect detectability.

The relative observable outburst duration of black hole and neutron stars is affected by their distance from us, which determines the minimum luminosity, $L_{\text{lim}}$ at which they can be detected. This is illustrated in Figure 4.3, where the observable timescales ($t_{\text{det}}$) are plotted for a black hole LMXB (green dashed line) and a neutron star LMXB (solid red line) as a function of orbital period, at distances of 5, 10 and 15 kpc in the upper, middle and lower panels respectively. These correspond to limiting luminosities of $7 \times 10^{35}$, $3 \times 10^{36}$ and $6 \times 10^{36}$ erg s$^{-1}$ respectively. At orbital periods greater than $\sim 0.2$ days, the outburst duration for black hole systems is longer than that for neutron star systems in all three cases. However, at low orbital periods, neutron star LMXBs may be visible for longer. This is because neutron stars have larger observable peak luminosities (smaller $f_{\text{corr}}$). Over time, the neutron star luminosity falls more rapidly than the black hole case (neutron stars have a shorter e-folding times as they have smaller disc radii). If the outburst luminosity of a neutron star LMXB at a given orbital period drops below $L_{\text{lim}}$ before it drops below the outburst luminosity of an equivalent black hole system, then the neutron star system will be observed for longer.

4.3.3 Duty cycles

In the lower panels of Figure 4.2 I estimate the outburst duty cycle, which I define as $t_{\text{det}}/(t_q + t_o)$, where $t_o$ and $t_q$ are the outburst and quiescent timescales respectively. The
Figure 4.4 Outburst probabilities and recurrence times. Upper panel: Outburst probabilities as a function of orbital period for neutron star and black hole \((M_1 = 8M_\odot)\) LMXBs, plotted as solid red, and green dashed lines respectively. Lower Panel: Outburst recurrence times plotted against orbital period for neutron star and black hole LMXBs. Also shown, as a blue dotted line, is the mission length of the RXTE ASM survey; 15 years.

The quiescent timescale is estimated as \(t_q = M_D/|\dot{M}_2|\), where \(\dot{M}_2\) is the mass transfer rate from the secondary. I take \(M_D\), the disc mass, as:

\[
M_D = \int_0^{R_D} 2\pi R_D \Sigma_{\text{max}}(R) dR. \tag{4.9}
\]

Cannizzo et al. (1988) give \(\Sigma_{\text{max}}\) as:

\[
\Sigma_{\text{max}} = 11.4 \left(\frac{R}{10^{10}\text{cm}}\right)^{1.05} \left(\frac{M_1}{M_\odot}\right)^{-0.35} \alpha_c^{-0.85} \text{g cm}^{-2} \tag{4.10}
\]

where \(\alpha_c \sim 0.01\) is the cold state viscosity parameter. Using (4.10) and integrating I obtain:

\[
M_D = 2.4 \times 10^{21} \alpha_c^{-0.85} \left(\frac{M_1}{M_\odot}\right)^{-0.35} \left(\frac{R_D}{10^{10}\text{cm}}\right)^{3.05} \text{g}. \tag{4.11}
\]

Using the entire disc mass in the calculation of the quiescent timescale is a good approximation at short orbital periods, where X-ray irradiation is expected to result in the
accretion of all or most of the disc. The rate at which the disc is replenished with mass after an outburst is given by \(-\dot{M}_2\), which is calculated using the formulae of King (1988):

\[
-\dot{M}_2 = \begin{cases} 
10^{-10} \left( \frac{P(h)}{2} \right)^{-2/3} M_\odot \text{yr}^{-1} & \text{if } P(h) < 2 \\
6 \times 10^{-10} \left( \frac{P(h)}{3} \right)^{5/3} M_\odot \text{yr}^{-1} & \text{otherwise.}
\end{cases}
\] (4.12)

The equations above represent mass transfer driven by gravitational radiation \((P(h) < 2)\) and magnetic breaking (otherwise).

It is clear from the lower panels of Figure 4.2 that in Case A, the X-ray duty cycle of black hole LMXBs \(\rightarrow 0\) at periods below 0.1 day for all \(f > 0.01\). In Case B, the duty cycle is also noticeably reduced at short orbital periods. In Section 4.4.3.3., I compare this model to one with a constant \(\dot{M}_2\) and find no differences in my conclusions.

### 4.4 REPRODUCING THE OBSERVED DEARTH OF BH-LMXBS AT SHORT ORBITAL PERIODS.

Here, I test the effect of the switch to ADAF in hiding a population of short orbital period black hole systems by predicting the relative detection probabilities of neutron star to black hole LMXBs as a function of orbital period.

#### 4.4.1 Observational requirements

The All Sky Monitor (ASM) on board RXTE constantly surveyed the sky from March 1996 until January 2012 and discovered many of the known LMXBs in our Galaxy (Wen et al., 2006). The instrument cycled through a series of \(\sim 90\) s dwells on different areas of the sky in a stochastic pattern, so that a randomly chosen source was scanned typically 5-10 times per day (Levine et al., 1996).

The limiting flux, below which a source is undetectable, depends on the instrument used and the exposure time. An all-sky survey is more likely to see a transient object than an instrument with a narrow field of view. For this reason, I use the RXTE for all the observational comparisons in this study. The approximate daily sensitivity of the ASM is 10 mCrab (where 1 mCrab \(\sim 2.4 \times 10^{-11}\) ergs\(^{-1}\) cm\(^{-2}\)). I therefore assume an outburst is detectable above \(F_{\text{lim}} = 10\) mCrab, and that the observed outburst duration is \(t_{\text{det}} = t(F > F_{\text{lim}})\).

For a transient source to have been detected by the RXTE, it must have had at least one outburst during the lifetime of the survey. This depends on the recurrence time, \(t_{\text{rec}} = t_q\), and the length of the survey \((t_{\text{survey}} = 15\) years\), and can be estimated as:
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\[
P(\geq 1 \text{ outburst}) = \begin{cases} 
\frac{t_{\text{survey}}}{t_{\text{rec}}} & \text{if } t_{\text{survey}} < t_{\text{rec}}, \\
1 & \text{otherwise.}
\end{cases} \quad (4.13)
\]

Provided an outburst occurred during the survey time, it must have been bright enough, and have lasted long enough, for the survey to have detected it. For an outburst to have been unambiguously detected it must have been visible for at least 1 day (private communication, A. Levine & R. Remillard). If \( t_{\text{det}} < 1 \) day, the likelihood of it being observed \( (P_{\text{obs}}) \) decreases with the decay timescale. I therefore parameterise \( P_{\text{obs}} \) as:

\[
P_{\text{obs}} = \begin{cases} 
t_{\text{det}}/1 \text{ day} & \text{if } t_{\text{det}} < 1 \text{ day,} \\
1 & \text{otherwise.}
\end{cases} \quad (4.14)
\]

The total probability of having an outburst observed by the ASM is therefore:

\[
P_{\text{det}} = P_{\text{obs}} \times P(\geq 1 \text{ outburst}) \quad (4.15)
\]

and depends on both the mass of the system (i.e. whether it is a black hole or neutron star) and on its orbital period.

A further constraint on objects forming part of this analysis is whether they have been followed up optically to measure orbital parameters and identify the nature of the primary. I make the assumption here that the fraction of candidate objects that are currently unidentified or poorly studied is independent of the type of primary and its orbital period. In this case, this additional factor will not affect my results.

### 4.4.2 Relative detection probabilities of a population of LMXBs at 8 kpc

Initially, I make detection probability calculations for LMXBs located at the Galactic centre, approximately 8 kpc away, as assumed in Section 4.3. In this case, the outburst timescales \( t_{\text{det}} \) correspond to those in the middle panels of Figure 4.2. For orbital periods where \( t_{\text{det}} > 1 \) day, \( P_{\text{obs}} = 1 \) and \( P_{\text{det}} = P(\geq 1 \text{ outburst}) \). In the upper panel of Figure 4.4, outburst probabilities for black hole and neutron star LMXBs are plotted as a function of orbital period, using solid red and green dashed lines respectively. \( P(\geq 1 \text{ outburst}) \) is a function of \( t_{\text{rec}} \), which is plotted in the panel below. \( P(\geq 1 \text{ outburst}) \) drops below 1 only when \( t_{\text{rec}} \) increases above \( t_{\text{survey}} \) (blue dashed line). The discontinuity in \( t_{\text{rec}} \) at 2 hours marks the switch in mass transfer process from magnetic breaking to gravitational radiation in (4.12). Above this value, \( t_{\text{rec}} = M_d/|M_2| \propto R_d^3/P_{\text{orb}}^{5/3} \propto P_{\text{orb}}^{1/3} \), and below, \( t_{\text{rec}} \propto R_d^3/P_{\text{orb}}^{2/3} \propto P_{\text{orb}}^{8/3} \). Although the discontinuity in \( t_{\text{rec}} \) produces a negative spike in the black hole outburst probability curve, this feature does not affect my results (see Section 4.3.3 for a discussion).
Figure 4.5  Black hole LMXB detection probabilities and predicted orbital period distributions, at a fixed distance of 8kpc. Upper panels: the ratio of detection probabilities as a function of orbital period, of black hole ($M_1 = 8M_\odot$) to neutron star LMXBs, at 8kpc from the Earth. The solid red line shows the ratios with no inefficiency switch during outburst. The green (dashed) and dark blue (thin dotted) lines include a switch to inefficiency in the black hole systems following case A (sharp switch) and case B (smooth switch) respectively. Results are shown for a switch condition of $f = 0.03$ in the left hand panels, and $f = 0.05$ on the right. Lower panels: the predicted orbital period distributions, plotted as normalised cumulative histograms, using the same line styles as the panels above. Also plotted are the observed black hole (light blue dashed line) and neutron star (purple dotted line) distributions.
In Figure 4.5 (upper panels), I plot the relative detection probabilities with \( f = 0.03 \) (left panel) and \( f = 0.05 \) (right panel). The solid red line shows the ratio of black hole to neutron star detection probabilities without a switch to ADAF, and the green dashed and dark blue dotted lines show the results for sharp and smooth switches (Case A and Case B) respectively. At large orbital periods, when \( P_{\text{obs}} = 1 \), the ratios plotted are simply the ratios of \( P(\geq 1 \text{ outburst}) \) shown in the upper panels of Figure 4.4. When there is no efficiency switch, the black hole detection probability drops to zero at \( \sim 0.01 \) days, as \( t_{\text{det}} < 1 \) day. When an efficiency switch is included, this drop in detection probability ratios occurs at higher orbital periods, of \( \sim 0.05 \) \((0.07)\) days with a smooth switch, and \( \sim 0.2 \) \((0.3)\) days with a sharp switch when \( f = 0.03 \) \((0.05)\).

It is now possible to use this detection probability approximation to estimate how the switch to radiatively inefficient accretion affects the orbital period distributions of black hole LMXBs. By convolving the relative detection probability with the observed period distribution of neutron star LMXBs, one can predict an expected distribution for black hole systems. I use Sample 2, weighting each neutron star LMXB by the relative detection probability at that orbital period, assuming the population distribution of periods is the same as that observed for neutron stars. These predicted distributions are shown in the bottom panels of Figure 4.5 alongside the observed distributions for black hole and neutron star LMXBs from Figure 4.1. It is clear that the case A efficiency switch is able to reproduce the absence of black hole LMXBs below 0.1 days when \( f = 0.03 \), as well as to some extent, predicting the shape of the black hole orbital period distribution. As a simple confirmation, I use these results to estimate the number of black holes expected to have orbital periods below 0.17 days, in a sample the same size as Sample 1 (17 sources). I find it to be zero for a sharp switch, provided \( f \geq 0.03 \). It is particularly interesting that this transition lies in the expected range of a few percent of \( L_{\text{Edd}} \), as any other reason for the relative population difference would give arbitrary values of \( f \). Switch B does not produce a significant difference between the black hole and neutron star orbital period distributions, predicting 7 \((6)\) black holes below 0.17 days, for \( f = 0.03 \) \((0.05)\).

### 4.4.3 Relative detection probabilities of the Galactic population of LMXBs

If all LMXBs are 8 kpc from the Earth, \( P_{\text{obs}} \) is either 0 or 1, depending on their orbital period. In our Galaxy however, the probability of a source being detected while in outburst depends on its distance, as closer sources will be visible for longer. Therefore, when observing the LMXBs in our Galaxy, \( P_{\text{obs}} \) may range between 0 and 1, depending on the fraction of sources which are detectable. To address this issue, I now repeat my calculation using a Galactic distribution of LMXBs.

The Milky Way stellar population comprises a disc, bulge, and spheroidal component. Grimm et al. (2002) studied the population of LMXBs in our Galaxy and modelled them.
using these three components, finding a ratio of 2:1:0.8 of the mass of LMXBs in the disc : bulge : spheroid respectively. The spheroidal component accounts for globular cluster LMXBs, which are not part of this study, so I do not include this in my calculations. I distribute LMXBs in a disc:

\[ N_{\text{disc}}(r,z) \propto \exp\left(-a_{\text{disc}}\right) \]  \hspace{1cm} (4.16)

where

\[ a_{\text{disc}} = \frac{r_m}{r} + \frac{r}{r_d} + \frac{|z|}{r_z}, \]  \hspace{1cm} (4.17)

and a bulge, following the parameterisation of Binney et al. (1997)

\[ N_{\text{bulge}}(x,y,z) \propto \left(1 + \frac{a_{\text{bulge}}}{a_0}\right)^{-1.8} \exp\left(\frac{a_{\text{bulge}}^2}{a_m^2}\right), \]  \hspace{1cm} (4.18)

where

\[ a_{\text{bulge}} = x^2 + \frac{y^2}{\beta^2} + \frac{z^2}{\zeta^2}. \]  \hspace{1cm} (4.19)

I use parameters of \( r_m = 6.5 \text{kpc}, r_d = 3.5 \text{kpc}, r_z = 0.41 \text{kpc}, a_m = 1.9 \text{kpc}, a_0 = 0.1 \text{kpc}, \beta = 0.5 \text{kpc} \) and \( \zeta = 0.6 \text{kpc} \). I normalise both distributions by containing them within 16 kpc, the outer edge of our Galaxy, and by requiring a ratio of 2:1 of binaries in the Galactic disc, and bulge, assuming direct proportionally between number ratios and the mass ratios calculated in Grimm et al. (2002).

On the left panel of Figure 4.6, I plot this distribution, with 500 binaries, where \( x, \) and \( y \) are the cartesian coordinates in the plane of the galaxy, and on the right panel,
Figure 4.7 Black hole LMXB detection probabilities and predicted orbital period distributions for sources following a galactic spatial distribution. Upper panels: the ratio of detection probabilities as a function of orbital period, of black hole ($M_1 = 8M_\odot$) to neutron star LMXBs, distributed so as to follow the Galactic disk and bulge population. The solid red line shows the ratios with no inefficiency switch during outburst. The faint red line shows the same, but for a galactic distribution containing only a disc, and no bulge. The green (dashed) and dark blue (thin dotted) lines include a switch to inefficiency in the black hole systems following case A (sharp switch) and case B (smooth switch) respectively. Results are shown for a switch condition of $f = 0.03$ on the left hand panels, and $f = 0.05$ on the right. Lower panels: the predicted orbital period distributions, plotted as normalised cumulative histograms, using the same line styles as the panels above. Plotted as light blue dashed and purple dotted lines are the observed black hole and neutron star distributions respectively.

this is converted into a normalised histogram of their distance from the Earth (located at $x = 8 \text{kpc}$, $y = 0 \text{kpc}$). The vertical lines on this plot represent the distances out to which black hole LMXBs with orbital periods of 0.01 (green dashed line) and 0.1 (purple dotted line) days are guaranteed to be observed by RXTE when in outburst, i.e. $P_{\text{obs}} = 1$. These two lines bracket a significant proportion of the distribution, implying that short orbital period systems in this distance range are unlikely to be observable. In particular it is clear that an efficiency switch, shortening the outburst time, may significantly affect the likelihood of observing short orbital period systems.

4.4.3.1. Relative detection probabilities

In order to recreate the detection probabilities and predictions of Figure 4.5 using a distribution in LMXB distances, I run a Monte Carlo simulation of LMXBs in the Galaxy.
I create 200 000 LMXBs (100 000 each of black hole and neutron stars) with the same orbital period. These systems are distributed in 3 dimensions over the Galactic disc and bulge by randomly selecting $r$, $\theta$ and $z$ so that they follow the distributions in (4.16) and (4.18) respectively. The detection probability is then evaluated for each system and averaged to give an overall detection probability of LMXBs at that orbital period. I run models for orbital periods ranging from 0.01 to 100 days, until a convergence of 0.1% has been reached.

The results are plotted in Figure 4.7, in the same style as Figure 4.5. Additionally, in the upper panels, I plot the detection probability ratios with no efficiency switch for disc LMXBs only as a faint red line. This relation differs from the equivalent ratios plotted in Figure 4.5 (solid red line), because $P_{\text{obs}}$ now takes fractional values. Below $\sim 0.2$ days, the most distant black hole LMXBs in the galaxy become undetectable, because they have lower peak luminosities than neutron stars (cf. Figure 4.2). This causes the black hole LMXB detection probabilities, and hence the relative detection probabilities, to decrease. However, below $\sim 0.04$ days, the ratios begin to increase again. This is because the neutron star detection probability ratio drops more steeply than the black hole one. This effect is evident in Figure 4.3. From 15 kpc (bottom panel), to 5 kpc (top panel), the observable outburst timescale for short orbital period systems increases for both neutron star and black hole LMXBs, as the limiting luminosity becomes lower. However, the rate of increase for black holes is more significant. Therefore, as systems at greater distances, like 15 kpc and 10 kpc become undetectable, the black hole LMXBs become more observable than neutron star systems. The same shape can be seen in the disc + bulge result (solid red line), but here, another feature is present: a sharp decrease in the relative detection probabilities at an orbital period of $\sim 0.02$ days. This is due to the black hole bulge population becoming undetectable.

While the results for the sharp switch case are the same as in Figure 4.5, the smooth efficiency switch now alters the predicted black hole orbital period distribution more significantly. However, these results predict 6 (5) black hole systems below 0.17 days when $f = 0.03$ (0.05), so a smooth switch cannot explain the absence of observed short orbital period black hole LMXBs. While the sharp switch is still more successful at reproducing the observed distribution, it appears that, at higher values of $f$ (> 0.05), the smooth switch may have a similar effect.

4.4.3.2. Varying $M_2$

In the previous plots, $M_2$ has been fixed at 0.4 $M_\odot$. However, the actual dependence of $M_2$ on the systems’ other properties is uncertain, and it is likely that $M_2$ increases with orbital period. For this reason, I repeat my models with secondary masses that adhere to the main-sequence, Roche Lobe filling relation, $M_2 = 0.1P_{\text{orb}}$ (King & Ritter, 1998), with the results plotted in Figure 4.8. Here, the relative detection probability curves have a
4.4.3.3. The discontinuity in $\dot{M}_2$

The discontinuity (at $P_{\text{orb}} \sim 0.1$ days) in the duty cycles plotted in the bottom panels of Figure 4.2, is due to the switch in accretion mode from magnetic braking to gravitational radiation as described in (4.12), and also produces a feature in the relative detection probability plots. Since this is in the relevant orbital period range for this study, I also checked that my results were not sensitive to this switch by redoing my models with (a) a constant $\dot{M}_2$ and (b) no gravitational radiation mode at short orbital periods. These changes alter $P(\geq 1 \text{ outburst})$, which describes the shape of the relative detection probability plots. However, they have no effect on $P_{\text{obs}}$, which depends on the accretion efficiency switch. Therefore my results, and the effect of different types of accretion switch on the predicted
black hole orbital period distributions, are unaltered.

4.4.3.4. A non-constant recurrence time

As observations do not show consistent recurrence times for single sources, I have tested
the resilience of my results to allowing a range of recurrence times at each orbital period.
In the extreme case, the time spent in quiescence after each outburst is independent of any
previous event. In this case I can assume a Poisson distribution of outbursts so that:

\[ P(\geq 1 \text{ outburst}) = 1 - e^{-1/\tau} \]  

(4.20)

where \(1/\tau = t_{\text{survey}}/t_{\text{rec}}\).

This affects the shape of the black hole to neutron star detection ratio plots (upper
panels of Figure 4.5), but not enough to change the results in the lower panels, so it does
not affect my conclusions.

4.5 DISCUSSION

Within the framework of the simple outburst model above, it is clear that the lack of
known black hole LMXBs at short orbital periods can be explained by a transition to
radiatively inefficient accretion, provided that the switch to inefficient accretion is sharp.
I assume a fixed primary mass as well as constant disc densities and viscosities. All of
these are likely to vary between observed systems and would produce differences in the
details of the orbital period distributions. A more robust treatment aimed at reproducing
the observed period distributions would require detailed population synthesis calculations
incorporating these factors. Some of the major causes of uncertainty are discussed below.

4.5.1 Choice of parameters

The results above are sensitive to my choices of \(\rho\), \(\nu\), and the black hole \(f_{\text{corr}}\), which I
investigate in Figures 4.9, 4.10 and 4.11. In Figure 4.9 I plot the same six panels as in
Figure 4.2. The solid red line represents an 8\(M_\odot\) black hole system assuming no switch to
inefficiency at low luminosities. The other three lines show the effect of varying \(\rho\). Each
line assumes \(f = 0.03\), and the same \(\nu\) as used previously \((5 \times 10^{14} \text{ cm}^2 \text{ s}^{-1})\). The values
of \(\rho\) shown are, in order of decreasing density, \(5 \times 10^{-8} \text{ g cm}^{-3}\) (dark blue dotted line),
\(1 \times 10^{-8} \text{ g cm}^{-3}\) (purple, fine dotted line line) and \(5 \times 10^{-9} \text{ g cm}^{-3}\) (light blue dot-dashed
line). Both the peak luminosity relation and the orbital period at which the sharp switch
(switch A) produces a cut-off in the distribution, are affected by changes in density. The
smooth switch (Case B) also varies. By requiring that the gradient and peak luminos-
ity normalisation match observational expectations, it is possible to rule out the higher
density cases.
Similarly, in Figure 4.10, I plot the same systems, with $f = 0.03$, but this time I keep the density at its original value $(1 \times 10^{-8} \text{ g cm}^{-3})$ and allow $\nu$ to vary. I choose values of $1 \times 10^{14} \text{ cm}^2 \text{ s}^{-1}$ (dark blue dotted line), $5 \times 10^{14} \text{ cm}^2 \text{ s}^{-1}$ (purple fine dotted line) and $1 \times 10^{15} \text{ cm}^2 \text{ s}^{-1}$ (light blue dot-dashed line). Varying $\nu$ has a similar effect to varying $\rho$ for the peak luminosities. However, in addition, the timescales, which are inversely proportional to $\nu$, are extended by decreasing $\nu$.

In Figure 4.11, $\nu$ and $\rho$ are fixed at their original values, and I vary $f_{\text{corr}}$, the correction factor applied to the bolometric luminosity of the black hole systems, to estimate the X-ray luminosity. This depends on the inner disc temperature, which is likely to range from 0.5 to 1 keV. I plot values of 2.5 (dark blue dotted line), 4 (purple fine dotted line) and 8.5 (light blue dot-dashed line), which correspond to disc temperatures of 1, 0.7 and 0.5 keV respectively. Varying $f_{\text{corr}}$ alters the luminosities of the sharp switch case, and all three panels corresponding to the smooth switch case, in a similar way to varying $\nu$ and $\rho$. Since the inefficiency switch luminosity ($\sim 2 \times 10^{37} \text{ ergs s}^{-1}$) is higher than the limiting luminosity at all values of $f_{\text{corr}}$ shown, varying this parameter does not affect the timescale or duty cycle plots when a sharp efficiency switch is implemented.

It is clear that my choices of $\rho$ and $\nu$, which are inherently uncertain, significantly affect the feasibility of the sharp switch to inefficiency explaining a short orbital period black hole deficit. In addition, $f_{\text{corr}}$, which depends on the black hole temperature, is likely to affect the normalisation of my results. However, in all cases the smooth switch scenario only produces a "turn-off" the known peak luminosity-orbital period relation at higher orbital periods than observed.

### 4.5.2 Varying the switch steepness

I now investigate how sharp the switch needs to be to produce the observed cut-off. In Figure 4.12, I vary the dependence of $\eta$ on $\dot{M}$ such that:

$$\eta = 0.1 \left( \frac{\dot{M}}{f\dot{M}_{\text{Edd}}} \right)^n. \quad (4.21)$$

A larger value of $n$ more closely approximates the sharp switch scenario. Here I show the relative detection probabilities on the top panel for $n = 1$ (green dashed), $n = 3$ (dark blue dotted) and $n = 5$ (purple dotted) for $f = 0.03$ on the left panel and $f = 0.05$ on the right-hand panel. Indices of $n \gtrsim 3$ appear to fit the black hole distribution as well as the sharp switch itself, provided $f = 0.05$, with a prediction of 1(0) observable black hole systems below 0.17 days for $n = 3(5)$. 
Figure 4.9  The effect of disc density on outburst properties. Panels follow those in Figure 4.2. A black hole of $8 \, M_\odot$ with no efficiency switch is plotted in solid red. The other 3 lines represent black hole systems ($M_1 = 8M_\odot$) with an efficiency switch at $f = 0.03$. $\rho$ takes values $5 \times 10^{-8} \, g \, cm^{-3}$ (dark blue dotted line), $1 \times 10^{-8} \, g \, cm^{-3}$ (purple dotted line, as in Figure 2) and $5 \times 10^{-9} \, g \, cm^{-3}$ (light blue dot-dashed line).

Figure 4.10  The effect of disc viscosity on outburst properties. Panels follow those in Figures 4.2 and 4.9. A black hole of $8 \, M_\odot$ with no efficiency switch is plotted in solid red. The other 3 lines represent black hole systems ($M_1 = 8M_\odot$) with an efficiency switch at $f = 0.03$. Here $\nu$ is varied between $1 \times 10^{14} \, cm^2 \, s^{-1}$ (dark blue dotted line), $5 \times 10^{14} \, cm^2 \, s^{-1}$ (purple dotted line), and $1 \times 10^{15} \, cm^2 \, s^{-1}$ (light blue dot-dashed line).
Figure 4.11  The effect of \( f_{\text{corr}} \) on outburst properties. Panels follow those in Figures 4.2 and 4.9. A black hole of 8 \( M_\odot \) with no efficiency switch is plotted in solid red. The other 3 lines represent black hole systems (\( M_1 = 8M_\odot \)) with an efficiency switch at \( f = 0.03 \). In this case, \( f_{\text{corr}} \) takes values of 2.5 (dark blue dotted line), 4 (purple dotted line), and 8.5 (light blue dot-dashed line), which correspond to disc inner temperatures of 1, 0.7 and 0.5 keV respectively.

4.5.3 Observational issues

The above discussion is based on detections possible with the RXTE ASM, which discovered several new black hole transients during its 15 year mission (e.g. XTE J1118+480, XTE J1650-500, XTE J1550−56). However, there have been, and still are, many other surveys capable of discovering BHB transients, each with different sensitivity levels, energy coverage, time resolution and sampling. Indeed, RXTE also routinely performed scans of the Galactic bulge region with the PCA detectors, capable of detecting fainter sources but reduced angular and temporal coverage compared to the ASM. These scans discovered the BH candidate XTE J1752−223, among others.

Presently, the Burst Alert Telescope (BAT) on the Swift mission detects known and new transients in the 15-150 keV range, making it sensitive to the bright low/hard states of black hole X-ray binaries (BHB) during outburst in addition to its higher sensitivity in Crab units. This has been operating since 2004 and has discovered several new systems including Swift J1753.5−0127, a short period (3.24 hr) black hole candidate (Zurita et al., 2008). Since 2009, the MAXI monitor has been scanning almost the entire sky in the 2-20 keV X-ray band every 96 minutes and has detected several new transients. Both Swift
Figure 4.12  The effect of varying the switch steepness, $n$. Upper panels: the ratio of detection probabilities as a function of orbital period, of black hole ($M_1 = 8M_\odot$) to neutron star LMXBs. The solid red line shows the ratios with no inefficiency switch during outburst. The other 3 lines (green dashed, dark blue dotted, purple dotted) show the same result with a smooth case B switch and steepness $n = 1, 3,$ and $5$ respectively, as described in the text. Results are shown for a switch condition of $f = 0.03$ on the left hand panels, and $f = 0.05$ on the right. Lower panels: the predicted orbital period distributions, plotted as normalised cumulative histograms, using the same line styles as the panels above. Also plotted are the observed black hole (purple dotted line) and neutron star (light blue dot-dashed line) distributions.
and MAXI detected a new transient MAXI J1659+152, later identified as another short period BHB (2.4 hr; Kuulkers et al., 2013).

These, together with other past and current missions, increase the sensitivity to BHBs compared to using the RXTE ASM alone, increasing the chance of detecting outbursts in Galactic BHBs except at very low peak luminosity and/or very long recurrence times (>10 years).

4.5.3.1. Comparison to the Ritter-Kolb catalogue

I also note that in calculating the relative detection probabilities, I have neglected some of the bias that is likely to exist in the Ritter-Kolb catalogue. Systems with confirmed orbital periods must have had an optical follow-up after their initial X-ray detection. Such a follow-up may be unsuccessful if the X-ray flux is very uncertain, or if the optical/IR counterpart is too faint. Both of these factors are biased against short period systems. I justify neglecting these issues by arguing that they are likely to affect black hole and neutron star systems equally, so would not enter into the ratio calculations.

4.5.4 A hidden population of short orbital period black holes?

The fact that the values of \( f \) fall in the expected theoretical range provides good support for the premise that the transition to radiatively inefficient accretion underpins the lack of black hole LMXBs at short periods. Alternative arguments for the lack of black hole systems (based on system evolution, selection effects, etc) would produce arbitrary values of \( f \), and the fact that the cut-off period for black hole LXMBs coincides with theoretical expectations for radiatively inefficient accretion would be purely coincidental.

Despite this support for a sharp change in \( \eta \), observations do not show a sudden change in bolometric luminosity at the efficiency switch boundary. In addition, observational evidence (see e.g. Zhang et al., 1997; Homan et al., 2005) suggests a smooth switch between efficient and inefficient flows. These factors suggest that, rather than being lost to the black hole, accretion power may be shifted out of the X-ray band, through a sharp change in \( f_{\text{corr}} \). This may mean that systems become detectable in other wavelengths, making it harder for short orbital period black hole LMXBs to “hide” through inefficient accretion. Consequently, there is scope for a future study in which the detection of hard X-ray dominated LMXB spectra could be assessed.

It is worth noting that black hole LMXBs with \( P_{\text{orb}} \sim 0.1 \) day would be expected to produce low luminosity (\(~ \text{few} \% L_{\text{Edd}}\)) and short (\(~ 1 \) day) outbursts. These systems would have low detection probabilities and would be rare, but may be observable in locations where formation conditions were favourable. It may be possible that short period black hole LXMBs form some fraction of the low-luminosity, short-duration, very fast X-ray transients observed in the Galactic Centre (e.g. Wijnands & van der Klis, 2000;
Maccarone & Patruno, 2013).

### 4.6 CONCLUSIONS

In summary, I have shown that there is some evidence for a statistically significant absence of black hole LMXBs with $P_{\text{orb}} \leq 0.1 \text{d}$ in the Ritter-Kolb catalogue (Ritter & Kolb, 2003), and that the neutron star and black hole distributions appear to be different at short orbital periods.

I have investigated whether this could be due to black hole LMXBs being preferentially hidden from view at short orbital periods. Using a simple model of LMXB outbursts, I have demonstrated that at short orbital periods, where the peak luminosity drops to the threshold for radiatively inefficient accretion, black hole LMXBs have lower peak luminosities, outburst durations and X-ray duty cycles than comparable neutron star systems. These factors can combine to severely reduce the detection probability of short period black hole LMXBs relative to neutron star systems. The transition to radiatively inefficient accretion is able to explain the lack of black hole LMXBs at $P_{\text{orb}} \leq 0.1 \text{d}$ with transition accretion rates in the accepted range of a few percent of $L_{\text{Edd}}$, providing the switch to inefficiency is sharp.

However, if the switch to inefficiency is gradual, it is unlikely that radiatively inefficient accretion could explain the period distribution. This would imply that some other observational bias is preventing the detection of short orbital period black holes compared to neutron stars, or that no short orbital period black hole systems exist.
### Table 4.2  List of LMXBs with known orbital periods

<table>
<thead>
<tr>
<th>Object Name</th>
<th>RA</th>
<th>Dec</th>
<th>Type</th>
<th>Transient</th>
<th>$P_{\text{orb}}$</th>
<th>Source</th>
<th>Sample</th>
</tr>
</thead>
<tbody>
<tr>
<td>V1487 Aql</td>
<td>19 15 11.5</td>
<td>+10 56 45 l</td>
<td>BH</td>
<td>Yes</td>
<td>30.8</td>
<td>Neil et al. (2007)</td>
<td>(i)</td>
</tr>
<tr>
<td>V404 Cyg</td>
<td>20 24 03.8</td>
<td>+33 52 03 l</td>
<td>BH</td>
<td>Yes</td>
<td>6.4714</td>
<td>Casares &amp; Charles (1994)</td>
<td>(i)</td>
</tr>
<tr>
<td>V4641 Sgr</td>
<td>18 19 21.6</td>
<td>−25 24 25 l</td>
<td>BH</td>
<td>Yes</td>
<td>2.8173</td>
<td>Orosz et al. (2001)</td>
<td>(i)</td>
</tr>
<tr>
<td>V1033 Sco</td>
<td>16 54 00.1</td>
<td>−39 50 45 l</td>
<td>BH</td>
<td>Yes</td>
<td>2.62120</td>
<td>Greene et al. (2001)</td>
<td>(i)</td>
</tr>
<tr>
<td>BW Cir</td>
<td>13 58 09.9</td>
<td>−64 44 05 l</td>
<td>BH</td>
<td>Yes</td>
<td>2.54451</td>
<td>Casares et al. (2009)</td>
<td>(i)</td>
</tr>
<tr>
<td>V821 Ara</td>
<td>17 02 49.4</td>
<td>−48 47 23 l</td>
<td>BH</td>
<td>Yes</td>
<td>1.7557</td>
<td>Hynes et al. (2003)</td>
<td>(i)</td>
</tr>
<tr>
<td>V381 Nor</td>
<td>15 50 58.8</td>
<td>−56 28 35 l</td>
<td>BH</td>
<td>Yes</td>
<td>1.542033</td>
<td>Orosz et al. (2011)</td>
<td>(i)</td>
</tr>
<tr>
<td>IL Lup</td>
<td>15 47 08.3</td>
<td>−47 40 11 l</td>
<td>BH</td>
<td>Yes</td>
<td>1.116407</td>
<td>Orosz (2003)</td>
<td>(i)</td>
</tr>
<tr>
<td>V2107 Oph</td>
<td>17 08 14.1</td>
<td>−25 05 32 l</td>
<td>BH</td>
<td>Yes</td>
<td>0.521</td>
<td>Remillard et al. (1996)</td>
<td>(i)</td>
</tr>
<tr>
<td>GU Mus</td>
<td>11 26 26.6</td>
<td>−68 40 32 l</td>
<td>BH</td>
<td>Yes</td>
<td>0.432602</td>
<td>Orosz et al. (1996)</td>
<td>(i)</td>
</tr>
<tr>
<td>QZ Vul</td>
<td>20 02 49.4</td>
<td>+24 14 11 l</td>
<td>BH</td>
<td>Yes</td>
<td>0.344087</td>
<td>Ioannou et al. (2004)</td>
<td>(i)</td>
</tr>
<tr>
<td>V616 Mon</td>
<td>06 22 44.4</td>
<td>−00 20 45 l</td>
<td>BH</td>
<td>Yes</td>
<td>0.323014</td>
<td>González Hernández &amp; Casares (2010)</td>
<td>(i)</td>
</tr>
<tr>
<td>J1650-4957</td>
<td>16 50 00.8</td>
<td>−49 57 45 l</td>
<td>BH</td>
<td>Yes</td>
<td>0.3205</td>
<td>Orosz et al. (2004)</td>
<td>(i)</td>
</tr>
<tr>
<td>MM Vel</td>
<td>10 13 36.3</td>
<td>−45 04 32 l</td>
<td>BH</td>
<td>Yes</td>
<td>0.285206</td>
<td>Filipenko et al. (1999)</td>
<td>(i)</td>
</tr>
<tr>
<td>V406 Vul</td>
<td>18 58 41.7</td>
<td>+22 39 30 l</td>
<td>BH</td>
<td>Yes</td>
<td>0.274</td>
<td>Corral-Santana et al. (2011)</td>
<td>(i)</td>
</tr>
<tr>
<td>V518 Per</td>
<td>04 21 42.8</td>
<td>+32 54 27 l</td>
<td>BH</td>
<td>Yes</td>
<td>0.212160</td>
<td>Webb et al. (2000)</td>
<td>(i)</td>
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<tr>
<td>KV UMa</td>
<td>11 18 10.9</td>
<td>+48 02 13 l</td>
<td>BH</td>
<td>Yes</td>
<td>0.16995</td>
<td>González Hernández et al. (2008)</td>
<td>(i)</td>
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<tr>
<td>J1744-2844</td>
<td>17 44 33.1</td>
<td>−28 44 27 l</td>
<td>NS</td>
<td>Yes</td>
<td>11.8367</td>
<td>Finger et al. (1997)</td>
<td>(ii),(iii)</td>
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<tr>
<td>0042+3244</td>
<td>00 44 48.8</td>
<td>+33 00 33.2</td>
<td>NS</td>
<td>Yes</td>
<td>11.6</td>
<td>Watson &amp; Ricketts (1978)</td>
<td>(ii),(iii)</td>
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<tr>
<td>V1333 Aql</td>
<td>19 11 15.9</td>
<td>+00 35 06 l</td>
<td>NS</td>
<td>Yes</td>
<td>0.78950</td>
<td>Chevalier &amp; Ilovaisky (1998)</td>
<td>(ii),(iii)</td>
</tr>
<tr>
<td>J0556-3310</td>
<td>05 56 46.3</td>
<td>−33 10 26 l</td>
<td>NS</td>
<td>Yes</td>
<td>0.684</td>
<td>Cornelisse et al. (2012)</td>
<td>(ii),(iii)</td>
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<td>Dec</td>
<td>Source</td>
<td>Comment</td>
<td>Ref.</td>
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<tr>
<td>V822 Cen</td>
<td>14 58 21.9</td>
<td>-31 40 07.1</td>
<td>NS</td>
<td>Yes</td>
<td>0.629052</td>
<td>Casares et al. (2007) (ii),(iii)</td>
<td></td>
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<tr>
<td>QX Nor</td>
<td>16 12 42.9</td>
<td>-52 25 23.1</td>
<td>NS</td>
<td>Yes</td>
<td>0.5370</td>
<td>Wachter et al. (2002) (ii),(iii)</td>
<td></td>
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<tr>
<td>J1749-2808</td>
<td>17 49 31.8</td>
<td>-28 08 05.2</td>
<td>NS</td>
<td>Yes</td>
<td>0.367369</td>
<td>Markwardt &amp; Strohmayer (2010) (ii),(iii)</td>
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<tr>
<td>J1745-2901</td>
<td>17 45 35.7</td>
<td>-29 01 34.11</td>
<td>NS</td>
<td>Yes</td>
<td>0.347960</td>
<td>Hyodo et al. (2009) (ii),(iii)</td>
<td></td>
</tr>
<tr>
<td>V2134 Oph</td>
<td>17 02 06.4</td>
<td>-29 56 44.1</td>
<td>NS</td>
<td>Yes</td>
<td>0.296505</td>
<td>Oosterbroek et al. (2001) (ii),(iii)</td>
<td></td>
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<tr>
<td>LZ Aqr</td>
<td>21 23 14.5</td>
<td>-05 47 53.1</td>
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We are all in the gutter, but some of us are looking at the stars.

Oscar Wilde, Lady Windermere’s Fan
CHAPTER 5. CONCLUSIONS

In this thesis, I set out to assess the effects of accretion onto stellar mass black holes in binary systems in two contexts. Firstly, I sought to evaluate their influence on their high-redshift environments during the epoch of cosmological reionization. Secondly, after finding a dearth of short-period black hole binaries in our Galaxy, I aimed to investigate whether they were truly absent, or just hidden from our view due to inefficient accretion.

In this chapter, I will review my findings, conclusions, and the wider ramifications of both of these studies. I will also evaluate future research directions for each project, before making my final summary.

5.1 THE INFLUENCE OF HMXBS ON THE PRIMORDIAL INTER-GALACTIC MEDIUM

Accretion onto black holes in High Mass X-ray Binaries may have been an important source of energy in the early universe. However, the extent to which they influenced their surroundings has been highly debated. While some authors have suggested that they could contribute to cosmological reionization by \( \gtrsim 10\% \), others have claimed that their influence must be negligible, to prevent the violation of current observational constraints.

In my thesis, I have developed an accurate model of the time-varying spectral energy distribution arising from an evolving stellar cluster with a HMXB phase, which I argue is ideal suited for resolving this debate. Because my model makes no assumptions about the global influence of HMXBs, it has been possible to make an unbiased assessment of the energetic boost that a HMXB phase could add to a stellar ionizing population. Additionally, by modelling plausible HMXB spectra, I have been able to move beyond the qualitative energetics arguments used by some authors when tackling this issue.

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In Chapter 2, I introduced my Monte Carlo stellar population synthesis, which is a further development of the globular cluster model presented by my collaborators in Power et al. (2009). My most major development to this code was to model the cluster’s time-dependent SED, using observationally motivated source spectra. The HMXB spectrum was based on that of Cygnus X-1 in its high state, and is more characteristic of black hole binary spectra than the power-law models that have been used in other studies. With this code I monitored the total energy output of the cluster with, and without, a HMXB phase. In the hydrogen-ionizing energy range of \( 13.6 < E < 1 \text{keV} \), I found that HMXBs dominated the energetic output of a cluster after just 20 Myr, once the most massive stars had left the main sequence. Over the entire lifetime of the cluster, HMXBs contributed up to 10% of the population’s ionizing power. I parameterised my results as simple fitting formulae that predicted the enhancement in ionizing energy provided by HMXBs as a function of stellar mass. These formulae can now be used in cosmological simulations, as a method for implementing HMXB feedback. Finally, I showed that these models do not
CHAPTER 5. CONCLUSIONS

violate the unresolved soft X-ray background constraint, to which all $\gtrsim 1$ keV photons at $z \gtrsim 6$ are expected to contribute.

In Chapter 3, I used the model results from the previous chapter to make my own investigation into the influence of HMXBs on the primordial IGM. For this work, I used a one-dimensional radiative-transfer model, to predict the ionization and temperature profiles of the IGM surrounding a stellar cluster. To begin with, I modelled a constant density environment between $z = 14.5$ and $z = 10$. This effectively assumed that all of the energy from the cluster was able to escape from its galactic environment, into the IGM. In this case I found that including a HMXB phase resulted in negligible enhanced ionization and heating. This is because their short lifetimes, and high-energy SEDs, were insufficient for the build up of an X-ray ionizing background. I then considered a cluster embedded in a $10^8 M_\odot$ NFW halo. Here, the most massive stars initially ionized a large IGM volume. However, in the stellar-only case, the region rapidly recombined as later UV radiation was unable to escape the high-density galactic core. HMXBs, emitting penetrating X-rays, were able to keep these regions ionized for longer, although they did not increase the ionized volume. The increased number of free-electrons in the resulting partially ionized regions promoted further cooling via Bremsstrahlung and Inverse Compton processes. This meant that a HMXB phase actually led to lower IGM temperatures.

To conclude, in the context of this model, I have found that HMXBs do not make a major contribution to the reionization or heating of the IGM. This is in contrast to recent studies, which predict a much stronger influence. Although HMXBs boost the ionizing energy by $\lesssim 10\%$, their short lifetimes, and hard spectra, lead to a low X-ray background that is insufficient to ionize hydrogen or helium over short timescales.

5.1.1 Future research prospects

Despite this somewhat negative result, there are still several avenues through which HMXBs may have influenced their primordial galactic environments. HMXBs may have had a significant influence within their host galaxies, particularly the high-density regions that UV radiation is unable to penetrate. I have also found that the X-ray escape fractions are very high, meaning that sufficient build up of an ionizing X-ray background in the IGM may have been possible in the case of ongoing star formation. Both of these options warrant further study. In this section, I highlight the future research avenues that could be taken, and discuss some important improvements that could be made to my stellar cluster models.

5.1.1.1. Model improvements and extensions

There are several uncertainties related to the binary microphysics within my stellar population models, which it would be feasible to improve. These include the HMXB spectrum,
which is currently non-time varying. In reality, HMXBs are known to evolve through different spectral states associated with a range of luminosities, so it would be useful to model these in more detail. Additionally, in Section 2.4.3, I discussed some of the most uncertain simulation parameters, which included my assumptions of HMXB luminosities, lifetimes, and “survival fractions” – the proportion of candidate systems which survive beyond their supernova phase to become HMXBs. I have checked that my results are reasonable by some general comparisons between these models and observed HMXB populations, but there are further observational comparisons I could make to improve their accuracy. For example, some star formation regions, such as “the Antennae” (Fall et al., 2005; Fabbiano et al., 2001) have been well studied to the extent that the X-ray luminosity, SED, and number of bright HMXBs within the region are well-known. Developing a stellar population simulation that matched these details by seeding HMXBs into a model environment, and making mock observations, may potentially improve our understanding of the most uncertain model parameters.

In moving towards a more observationally-motivated stellar population model, it may be informative to base HMXB properties on the observed redshift-evolving SFR-X-ray luminosity relationship (e.g. Mineo et al., 2014; Tremmel et al., 2013; Basu-Zych et al., 2013). The HMXB luminosity function has been characterised by Mineo et al. (2012) and has since been used by Fragos et al. (2013) to characterise the energetic influence of HMXBs at high redshifts. In Section 3.4.2, I suggested that X-rays from ongoing HMXB formation may be able to maintain a low-level ionizing background capable of “pre-ionizing” the distant IGM. Using this SFR-$L_X$ relation, it would be possible to make my stellar population model time-dependent, tracing the universal star formation rate, to investigate this scenario.

Another possible development to my model might be to account for other sources of ionizing radiation within the stellar population. For example, Johnson & Khochfar (2011) predict a 10% contribution to reionization by supernovae. As part of the HMXB formation process, a natural expansion of the model may be to include a supernova phase into the binary evolution. There have also been recent investigations of the accretion of the ISM onto isolated stellar mass black holes (Wheeler & Johnson, 2011). This may be a further source of energy within the cluster, worthy of consideration.

5.1.1.2. The effects of HMXBs on their local environments

While this study has found that HMXBs exerted little effect on the high-redshift IGM, they may yet have profoundly influenced the local environments within their host galaxies. Hence, this may be an important future research focus. While one can justify treating a stellar population as a single source when calculating its effects on the distant IGM, a spatial distribution of sources would be required for studying their local effects. Also, within these environments, it may be necessary to account for alternative HMXB feed-
back mechanisms, such as their ability to drive large scale kinetic outflows (Blundell et al., 2001; Gallo et al., 2005). Finally, a three-dimensional, density-varying galactic environment would also need to be modelled. Consequently, studying the effects of HMXBs in a galactic environment is likely to be beyond the means of a one-dimensional radiative-transfer code. Instead, it would require three-dimensional hydrodynamic simulations that incorporate radiative transfer, such as those used in Jeon et al. (2013).

### 5.2 DETECTING SHORT ORBITAL PERIOD BLACK HOLE LMXBS IN OUR GALAXY

In Chapter 4, I compared the orbital period distributions of the confirmed black hole and neutron star LMXBs in our Galaxy, using data from the publicly available Ritter-Kolb catalogue (Ritter & Kolb, 2003). I found that there was a statistically significant dearth of short-orbital period black hole LMXBs ($P_{\text{orb}} < 4$ hour) if the two populations were assumed to follow the same distribution. This could be due either to a true absence of such systems in the Galaxy, or to black hole LMXBs being preferentially hidden from view at short orbital periods. In this study, I investigated the latter possibility.

Using a simple model of LMXB outbursts, I showed that, at short orbital periods, where the peak luminosity drops below the threshold for radiatively inefficient accretion, black hole LMXBs have lower peak luminosities, outburst durations and X-ray duty cycles than comparable neutron star systems. I studied two types of efficiency switch: (i) a sharp switch where $\eta = 0$ below a fraction $f$ of the Eddington Luminosity, $L_{\text{Edd}}$, and (ii) a power-law efficiency decrease of $\eta \propto \dot{M}^n$, for $L < fL_{\text{Edd}}$. By monitoring the outburst properties in a model population of black hole LMXBs, I found that those with short orbital periods became undetectable to the RXTE All Sky Survey, only in the case of an instantaneous switch or for $n \geq 3$.

My results suggest that there may be an unobserved short orbital period black hole population, depending on the properties of radiatively inefficient flows. Otherwise it implies that there is either some other observational bias or that there is a genuine absence of such sources that must be explained via evolutionary arguments.

#### 5.2.1 Future research prospects

Rather than using a simple model of outbursts following analytical prescriptions, it may be possible to further develop the model by monitoring luminosities and timescales using a one-dimensional disk code, solving the disk equations numerically. Predictions such as time spent in quiescence are particularly uncertain and could benefit from a more detailed analysis.
As observations do not show a sharp change in bolometric luminosity at the efficiency switch boundary, it is possible that, rather than being advected onto the black hole, accretion power may instead be shifted away from the soft X-ray band at the onset of inefficiency. In this situation, they may become detectable in other wavelengths, such as hard X-rays. Hence, observations by telescopes such as INTEGRAL could be used to find them. A future research direction might be an assessment of the detectability of radiatively-inefficient accretion systems using this instrument. This may narrow the scope of radiatively inefficient accretion to hide short-orbital period LMXBs, and place tighter constraints on the Galactic population.

It is also possible that, despite low detection probabilities, black hole systems may still be detectable in environments where formation conditions are favourable. In recent years, a new transient population known as Very Faint X-ray Transients, which have low peak luminosities and short outburst durations, have been discovered (Sakano et al., 2005; Wijnands et al., 2006; Maccarone & Patruno, 2013). Some fraction of this population may be black hole LMXBs with peak outburst luminosities close to, or below, the threshold for inefficient accretion. Analysing this population would be an important future research trajectory.

5.3 FINAL COMMENTS

My thesis work can be summarised by two main conclusions. Firstly, in our own galaxy, I have discovered statistical evidence for a lack of short period black hole LMXBs, provided they are assumed to follow the neutron star LMXB population. I propose that this may be due to inefficient accretion flows preferentially hiding the black hole population. However, this is only possible if the switch occurs sharply at a threshold mass accretion rate. Otherwise, alternative observational or evolutionary explanations are required. Secondly, I predict that HMXBs had little ionizing or heating effects on the IGM at $z \gtrsim 6$, countering recent favourable arguments. However, they may exert significant local influence on their host galaxies.

Although my conclusions can be considered to be negative results, it is nevertheless important to rule out possible theories about the progress of reionization, and the sources responsible for it. Arguably this leads to a more solid understanding of a crucial period in the universe’s history.


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