Black-Hole Binaries As Relics Of Gamma-Ray Burst/Hypernova Explosions

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Abstract of the Dissertation

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The Collapsar model, in which a fast-spinning massive star collapses into a Kerr black hole, has become the standard model to explain long-soft gamma-ray bursts and hypernova explosions (GRB/HN). However, stars massive enough (those with ZAMS mass $\gtrsim (18-20)M_\odot$) to produce these events evolve through a path that loses too much angular momentum to produce a central engine capable of delivering the necessary energy. In this work I suggest that the soft X-ray transient sources are the remnants of GRBs/HNe. Binaries in which the massive primary star evolves a carbon-oxygen burning core, then start to transfer material to the secondary star (Case C mass transfer), causing the orbit to decay until a common-envelope phase sets in. The secondary spirals in, further narrowing the orbit of the binary and removing the hydrogen envelope of the primary star. Eventually the primary star becomes tidally locked and spins up, acquiring enough rotational

energy to power up a GRB/HN explosion. The central engine producing the GRB/HN event is the Kerr black hole acting through the Blandford-Znajek mechanism.

This model can explain not only the long-soft GRBs, but also the subluminous bursts (which comprise $\sim 97\%$ of the total), the long-soft bursts and the short-hard bursts (in a neutron star, black hole merger). Because of our binary evolution through Case C mass transfer, it turns out that for the subluminous and cosmological bursts, the angular momentum Ω is proportional to $m_D^{3/2}$, where m_D is the mass of the donor (secondary star). This binary evolution model has a great advantage over the Woosley Collapsar model; one can "dial" the donor mass in order to obtain whatever angular momentum is needed to drive the explosion.

Population syntheses show that there are enough binaries to account for the progenitors of all known classes of GRBs.

To my family.

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Chapter 1

Introduction.

1.1 Gamma-Ray Bursts and Hypernovae

Gamma-ray bursts (GRBs) are sudden flashes of γ rays with energies in the range of approximately 1 keV (for X-Ray Flashes, XRFs) to a few MeV. They were first discovered at the end of the 1960's. Military satellites, intended to look for nuclear-weapon-test gamma-ray signatures (Vela for the USA and Konus for the USSR [138]), made the shocking observation that most of the gamma-ray events they were detecting were not coming from the Earth beneath them, but rather from the skies above them, and, furthermore, they did not match the signature of a nuclear weapon explosion.

In the following years, many theories were developed in order to explain the origin of the GRBs, at some point the number of theories outnumbered the number of observed events. Originally those theories advocating a Galactic origin gained strength, given that the large amount of energy they emit seemed too large for any cosmological source. However, a Cosmological origin was suggested by the following properties [111]: a spatially isotropic (but not homogeneous) distribution, a number density that increased slowly when their intensity decreased (but not as a $-3/2$ power law [96]) and some gravitationally lensed events (similar to quasars).

Further studies and observations have verified that they have, indeed, an isotropic distribution, as well as shown that the observed event rate is about a GRB per day, and they have a bimodal distribution on their duration (T_{90}), with a short-duration maximum at ~ 0.3 seconds, a minimum at ~ 2 seconds and a long-duration maximum at ~ 30 seconds (see, e.g., plot 11.3 in [138] for BATSE observations).

It later became clear that the explosions were not spherically symmetric, like those of supernovae. Instead, most of the energy is focused in highly relativistic jets, and we can only observe the GRBs which are pointed in our direction. This realization solved the problem of the excessive amount of energy required in a spherical explosion.

Today, there are two main models explaining their high-energy radiation, duration and cosmological origin. The first one, aiming to explain the Short-Hard GRBs (or SHBs), involves the merger of two compact objects; that is, they are produced either by the coalescence of a Double Neutron Star (DNS) binary, or a Black Hole, Neutron Star (BH-NS) binary. In chapter 7 we will make the case for a BH-NS merger (as opposed to a DNS merger) by adapting our model for Long-Soft GRBs (LS-GRBs) to reproduce the energies associated with the SHBs using the same central engine.

The second model explains the origin of the LS-GRBs as failed Supernovae (SNe), that is, even though a SN might be launched, the shockwave will fail to leave the core of the star, thus the core will collapse into a BH. When such a BH is formed, given some initial conditions, this may still produce a very energetic explosion which initially launches a LS-GRB and later produces a hypernova (HN) explosion.

Now, hypernovae (HNe) are stellar explosions about an order of magnitude more energetic than SNe. While a normal SN may have a luminous energy of about 1 bethe[1] (10^{51} ergs), a typical HN explosion has a luminous energy of several bethes. They are usually Type I_b or I_c, that is they show no hydrogen lines and they may or may not show helium lines. They were predicted to be an observational counterpart to LS-GRBs by [145], and later modeled by [88]. They will be the subject of the next section (1.2).

1.2 Collapsar Model

As recently as the 1990's the origin of the sources of the GRBs was not confirmed to be cosmological, extra-galactic, not even outside the solar system.

[1]also found in the literature as foe, an acronym for "ten to the Fifty One Ergs"

However the isotropy in their spatial distribution was a large factor favoring the cosmological-origin models. Nevertheless [145] discussed the idea of merging compact objects (either BH-NS or DNS) and proposed the idea of failed SNe (of Type $I_{b/c}$) as two mechanisms capable of producing a non-spherically symmetric but extremely energetic source of neutrino beams which eventually would produce the observed GRBs. These two ideas competed against many other ones as well as each other in order to explain the source of the GRBs, and it was not until the late 1990's, early 2000's that some observational evidence supported the idea the LS-GRBs came from "failed SNe", or "Collapsars" (by linking GRB 980425 with SN1998bw [55] and GRB 030329 with SN2003dh [67]). The SHB part of the population is believed to come from compact object mergers. This, by no means, means that there is a consensus on how the central engines of GRBs work, nor even that there are no other competing theories that can explain the observational evidence as well (e.g., explosions involving rapidly rotating "magnetars", i.e. pulsars with magnetic fields on the order of 10^{15} G, have also been proposed as LS-GRBs progenitors).

Leaving aside the SHB problem for the moment and concentrating on the LS-GRBs, let us move forward and explain the fundamentals of the Woosley collapsar model. This model is a fundamental, however, incomplete, part of the model I will use in this work in order to explain the central engine and source of energy to power the greatest explosions observed in the Universe since the Big Bang.

In order to explain what a *failed* supernova is, it is important to know first what a *successful* SN is. Hence a brief explanation of stellar evolution of massive stars and their gravitational collapse is in order.

1.2.1 Stellar Evolution

A star can be understood, in a very simplified scheme, as a large self-gravitating mass (from some 10^{32} g to a few times 10^{35} g) in hydrostatic equilibrium. Meaning that the gravitational force of the mass of the star is balanced by the pressure exerted by the energy released in the thermonuclear fusion of atoms in the stellar core. The mass consists mainly of hydrogen and helium plus a small amount of metals ($\sim 71\%$, $\sim 27\%$ and $\sim 2\%$ of the total mass, respectively; and where by *metal*, in astronomy, it is understood any atom other than H or

He).

In order to ignite a reaction fusing ^{1}H into ^{4}He in the core of a star, the temperature (originated by compression of the H-and-He cloud by the gravitational collapse) must reach some 10^{7} K. For stars with mass $m \lesssim 1M_{\odot}$, hydrogen will be fused into helium in the proton-proton ($p-p$) chain, which can be summarized as follows:

$$4p \rightarrow {}^{4}He + 2e^{+} + 2\nu_{e}, \tag{1.1}$$

producing an excess energy (from $(4m_{p} - m_{He})c^{2} = 26.731$ MeV) at a rate that will balance the gravitational attraction and prevent a collapse. However, in more massive stars the CNO cycle will dominate, as the slow $p-p$ reactions are replaced by faster weak-interaction processes where the heavier ^{12}C act as catalyzers. The CNO chain can be summarized by:

$$^{12}C + 4p \rightarrow {}^{12}C + {}^{4}He + 2e^{+} + 2\nu_{e} \tag{1.2}$$

(the energy yield being the same as in the $p-p$ chain given that the end product will only change 4 protons to a He atom, with the C acting only as a catalyst).

It can be calculated [40] that the luminosity L and the mass M of a star are related by

$$L \propto M^{n} \tag{1.3}$$

where n is measured from observations and it varies nonlinearly between 3 and 5.5. And in particular, for stars with $M_{ZAMS} \gtrsim 7M_{\odot}$ where the opacity drops to being proportional to the temperature

$$L \propto M^{3}. \tag{1.4}$$

Now, the lifetime τ of a star can be estimated from

$$\tau = E/L \propto Mc^{2}/L \propto M/L \Rightarrow \tau \approx M^{-2}. \tag{1.5}$$

Empirically however the lifetime of a star is proportional to a power of M between -2 and -4.

Therefore a small star, having a shallow gravitational potential, requires a low rate of H-burning reactions to maintain hydrostatic equilibrium and spends billions of years ($\tau \sim 10^{10}$ years for a $1M_\odot$ star, where $1M_\odot = 1.988 \times 10^{33}$ g$\approx 2 \times 10^{33}$ g) in its Main Sequence stage (H-burning stage). A more massive star, with a deeper gravitational potential, requires more pressure to maintain hydrostatic equilibrium, thus, burning through its larger H fuel reserves much faster and evolving through its Main Sequence in a much smaller amount of time ($\tau \sim 10^7$ years for a $\sim 10M_\odot$ star).

When the core of the star begins to run out of H, the rate of fusion reactions starts to decrease and the He-dense core begins to contract, further increasing its pressure and temperature. The heating of the core will ignite H-shell burning and this will cause a large expansion of the less-bound H envelope, thus the star will become a red giant. This contraction will continue until a pressure build up balances the gravitational collapse and a new hydrostatic equilibrium is established. Such equilibrium can be achieved in two possible ways. First, and the most likely case for the less massive stars, by degeneracy pressure of the electrons, i.e., forming a White Dwarf (WD). Second, by raising the temperature of the collapsing core to the point where a new fusion reaction can be ignited.

When a star originally more massive than a few solar masses finishes its hydrogen-fuel supply, its massive helium-rich core contracts and heats up very rapidly, so much so that electron degeneracy cannot set in before burning of the new fuel is ignited, He fuses into, mainly, carbon and oxygen, and avoids forming a He WD. This process, however produces much less energy than the efficient process of fusing H into He, lasting, therefore, only $\sim 10\%$ of the lifetime of the star, as compared to the $\sim 90\%$ which the star spends in main sequence. Nonetheless, if the star has a Zero-Age-Main-Sequence (ZAMS) mass such that $M_{ZAMS} \gtrsim (8-10)M_\odot$, the core temperature will rise again, at the end of the He-core burning stage, fast enough to avoid forming a $C-O$ WD. It will, in fact, avoid forming any kind of WD all together, always turning one stages ashes into next stage's fuel (passing through ^4He, ^{12}C, ^{20}Ne, ^{16}O and ^{28}Si) all the way up to ^{56}Fe (the most bound nucleus in nature).

1.2.2 Core Collapse

It most be noted that there are several models that attempt to reproduce supernova explosions from the collapse of a massive star. However, to this day, no computer code can successfully produce such a result throughout the range of masses from which supernovae are believed to be generated.

It is agreed (see Shapiro & Teukolsky [125]) that the large temperature gradients in the core cause convection to transport heat and homogenize the composition of the core. The consequence is core convergence, that is, models seem to predict that most massive stars will develop iron cores with masses around $1.5M_\odot$. Now the Chandrasekhar limiting mass (i.e., the maximum mass a WD can have which can be prevented to collapse by electron degeneracy pressure) can be expressed in terms of the electron ratio as:

$$M_{Ch} \sim 5.83 Y_e^2 M_\odot \tag{1.6}$$

(note, however, this formula takes the central temperature $T_c = 0$ and the adiabatic index $\Gamma = 4/3$).

It is clear that Y_e (the number of electrons per baryon) will only diminish as ^{28}Si burns into ^{56}Fe (since $Y_e(^{28}Si) = 0.5$ whereas $Y_e(^{56}Fe) = 0.464$), and, furthermore, as photodissociation of ^{56}Fe and neutronization occur, the limiting mass will diminish more and more and eventually will trigger the collapse of the core.

Hence, when the central part of the core has accumulated a substantial amount of ^{56}Fe, at a density of the order of $\rho_c \sim 4 \times 10^9$ g cm^{-3} [6], it will again begin to collapse further. Nevertheless, this time there will not exist any further source of pressure which might balance the strong self gravitation of the core, since ^{56}Fe has the largest binding energy per nucleon of all nuclei and all further fusion reactions (or fission reactions for that instance) are endothermic. At this stage the central part of the core collapses as a whole.

The collapse will only come to a stop by the degenerate pressure of the nucleons themselves (mainly neutrons and a decreasing amount of protons). This is achieved once the infalling material reaches a density close to nuclear ($\rho_0 \sim 2.7 \times 10^{14}$g cm^{-3}), at which point the equation of state (EOS) of nuclear matter considerably stiffens and stops the collapse as a whole, sending an extremely energetic sound wave in the reverse direction. The initial collapsed core which falls in and rebounds as a whole, launches the original outgoing sound wave. Starting from the center of the star and moving out radially, the surface where the speed of sound drops below the infall velocity is known as the sonic point. The sonic point determines the outer boundary of the homologous core ([7]). The rebound of the homologous core marks the beginning of the transformation of the original core implosion into a supernova explosion.

The sound wave moves outwards at the speed of sound in the collapsing-core medium, somewhere around 1/10 the speed of light (in vacuum, c). However, the propagating medium itself (that is, the rest of the collapsing core) is falling inwards at free fall velocities. These are supersonic from the outer boundary of the homologous core all the way out to a few thousand kilometers. This implies that the sound wave is trapped in the collapsing core in a similar way a runner is in a treadmill. However the energy of the sound wave keeps increasing and once the velocity of the infalling material becomes subsonic the remaining sound wave would be free to increase its radius, i.e. blow the outer layers of the star apart. This will happen if the core is not too massive and ends up draining away all the energy of the sound wave.

In the mean time the homologous core has been contracted beyond nuclear density by the material which has kept piling on it. It has been pushed further in from its equilibrium point, to the point of maximum scrunch. This energy is finally released as the core bounces back outwards and produces an extremely energetic shockwave which may attain a speed between 1/10 and 1/5 the speed of light in vacuum. However this shockwave loses energy (by dissociating the incoming nuclei into nucleons) and seems to stall at around 100 to 200 km from the center of the star.

The original shockwave, produced by the bounce of the homologous core, most likely fails to launch the supernova. However, the large density of the collapsing core and the large temperature produced by the nuclear reactions ($T_c \sim 8 \times 10^9\text{K} = 0.7$ MeV, [6]) produce a large opacity for the neutrinos (neutrinos first become trapped during the collapse when the density grows to $\rho \gtrsim 4 \times 10^{11}\text{gm cm}^{-3}$, [7]), which start accumulating at a radius close to the shockwave called the neutrinosphere. This scenario might be capable of revitalizing the shockwave by depositing energy from the homologous core plus its accumulating material into the region where the shockwave has stalled and finally launch the supernova, however there has been no conclusive evidence from the computer simulations that this mechanism is effective enough to explain the explosion mechanism except for a few very specific cases. However some new ideas have been more recently discussed (see for example [33] and [72]) where acoustic, gravitohydromagnetic, or other mechanisms where spherical symmetry is broken show promise of delivering codes which will produce supernovae explosions.

The main issue for launching a successful supernova explosion, however, is whether the shockwave produced by the bounce of the homologous core plus the energy later deposited on the stalled shockwave by the infalling material (heat from the kinetic energy of the infalling material, neutrinos, etc.) is strong enough to provide it with enough energy to survive the infalling mass still available in the outside of the iron core (a few tents of a solar mass), and to survive the energy sinks of dissociating a large portion of the incoming ^{56}Fe atoms and the neutronization.

1.2.3 Failed Supernovae

When the iron core is too massive, the shockwave fails to reach its outer boundary where density drops dramatically and would allow the shockwave to escape and launch the supernova. Instead, after a few seconds, the shock wave is dragged inwards. The protoneutron star may be too hot to allow a direct collapse but its temperature drops fast as neutrinos escape and it eventually gives in to its own weight. The neutron star limiting mass (equivalent of the Chandrasekhar mass limit for WDs) is easily surpassed by the massive collapsed core, producing a black hole.

The lack of evidence of a pulsar signal in the remnant of SN1987A suggests it may have collapsed into a black hole after a few seconds of launching the supernova explosion. If this were to be later confirmed by observations, it would suggest that stars with ZAMS masses as low as $18 - 20 M_{\odot}$, like the progenitor of SN1987A, may produce low mass black holes instead of neutron stars. This would put a tight constraint on the maximum mass of a neutron star and on the equation of state of nuclear matter, as in all likelihood the mass of the protoneutron star did not exceed the $1.5 - 1.6 M_{\odot}$ range once it collected fallback material.

If the NS mass limit is exceeded, the compact object is "doomed" to become a black hole of a few solar masses. Sometimes, like in the case of SN1987A, a supernova explosion is actually launched and a NS formed, but it later collapses into a low-mass black hole by accretion of fallback material. There exists however another possibility. If the massive collapsing core possesses a great amount of angular momentum it may not all fall directly into a central compact object but it may rather form a compact object and an accretion disk. In this case a gamma-ray burst and a hypernova explosion may be launched by a mechanism which we now briefly sketch.

Start with a core massive enough to produce a rapidly rotating black hole within a few seconds of the collapse of the material in the homologous core plus the material within a few degrees from the rotational axis, as this will not be centrifugally supported. Let us further assume that the homologous core collapses as a rigidly rotating body, i.e., there is no differential rotation. We will justify this assumptions in later chapters by showing that our model has correctly predicted some of the properties of observed black hole binaries. The resulting black hole is born spinning very rapidly as it must conserve angular momentum. The material which has too much angular momentum to fall directly into the black hole will start collapsing into an accretion disk surrounding the black hole in the plane perpendicular to the rotational axis of the star. This scenario will also have the effect of clearing (or drastically lowering the density along) the rotational axis of the Collapsar.

The core of the star is at an extremely large temperature (on the order of 10^9 K) and so the material is extremely ionized, which has the effect of "freezing" or locking the magnetic field to the collapsing core, increasing the magnetic flux as the core shrinks. Eventually, a very large magnetic field which rotates with the central compact object is produced . This field must not necessarily be oriented along the rotational axis of the star, but we know that conservation of magnetic flux must be obeyed at all times, even if a black hole is produced.

By conservation of angular momentum the central black hole is rotating faster than the material in the accretion disk and so the magnetic field dragged along by the black hole will necessarily have to "cut" through the plasma in the accretion disk heating it up. As described in appendix 8.3 the rotating black hole performs the role of a central engine depositing energy along the rotational axis. As long as the disk is not heated beyond the point where it explodes, this mechanism will be in place to produce a gamma-ray burst. However, once enough energy has been accumulated, the pressure inside the disk will grow to the point of blowing it apart, producing a hypernova and dismantling the central engine for the gamma-ray burst.

The general mechanism of a failed supernova collapsing into a black hole and powering up a gamma-ray burst and a hypernova is known as the Collapsar model [145]. The mechanism we use in our model to power up the central engine by using the interaction of the magnetic field with the accreting matterial is known as the Blandford-Znajek mechanism [13].

1.2.4 Massive Star Evolution And Angular Momentum

A rather strong problem facing the collapsar model is the need to retain enough angular momentum up to the time when the core collapses into a black hole.

Assume we have a high-mass strongly-magnetized ZAMS star with nearly breakup rotation, i.e, material in its equator is close to being unbound from the star or $v^2 \lesssim GM/R$, where M and R are the mass and equatorial radius of the star and v is the tangential velocity of a particle on the equatorial surface of the star. As the star evolves beyond main sequence, the star becomes a red giant. At this point material from the outer layer of the star will greatly expand and even if it still were gravitationally bound to the star, strong stellar winds may eventually expel it away. For a (hypothetical) rigidly rotating star, a great amount of the original angular momentum is located in the outer layers and is eventually lost with the mass. Even worse, the core of the star is shrinking and so it spins up as the outer layer expands and spins down, however the strong magnetic field of the star will tend to even out any differential rotation, transferring out even more angular momentum from the core into the outer layer, and, in doing so, draining away most of the angular momentum necessary to power up a gamma-ray burst and a hypernova [52].

This scenario seems to drain away enough angular momentum from the material that will eventually collapse into the compact object, even without a strong magnetic field. Therefore it is highly unlikely that single stars can be progenitors of GRB/HN explosions.

1.3 Evolution In Binaries

Where the collapsar model may fail to be produced from single evolving stars, it may be revived by evolving the collapsar as the massive star (or primary) in a binary system.

As mentioned in 1.2.1, massive stars burn through their available fuel in a shorter time span, therefore, after a binary system forms, and assuming little or no interaction (i.e. mass transfer) the more massive star will always reach the end of the main sequence before the lighter companion.

We need the pre-collapse star to be able to tap into the angular momentum available in the binary system before the black hole forms so the collapsar model can produce the gamma-ray burst/hypernova explosion, it is also required that the binary will not to be disrupted too early in the evolution of the primary because, as mentioned in 1.2.4, most of the angular momentum is lost when the primary becomes a giant.

This means that we need to *trick* the primary into evolving as a single star throughout its main sequence lifetime.

1.3.1 Case C Mass Transfer And Common Envelope Evolution

As Bethe et al. [10] have suggested, the primary not only should evolve as a single star through its main sequence lifetime, but also through its giant or He-core burning phase. In this way, the angular momentum is safely stored in the orbit of the binary until the primary reaches the supergiant or He shell burning stage. At this point the giant ought to fill its Roche Lobe and start transferring mass to close the orbit of the binary and, through common envelope evolution, spin up and power up the collapsar.

Mass transfer at this late stage of evolution is known as Case C mass transfer. When it occurs during the giant stage it is known as Case B, and if the transfer is prior to these stages it is called Case A.

For stars in the $\sim$ 18 to $\sim 40M_{\odot}$ range, Brown et al. [28] modified the evolution of the radius of massive stars [123], by reducing the winds, to the point where the radius of the star during He-core burning does not decrease. They found that the necessary distance for the stars to evolve separately until Case C mass transfer (i.e., prevent Roche lobe overflow, RLOF, until He burning has begun in the core) occurs is $\sim 1,500 R_\odot$.

1.3.2 Common Envelope Efficiency

As mass is transferred from the primary star to the secondary star through Roche lobe overflow, the orbital period of the binary decreases, further shrinking the Roche lobes and thus preventing mass transfer from stopping, this is called unstable mass transfer. Eventually, unstable mass transfer leads to the companion plunging into the envelope of the (supergiant) primary star and the orbit decreases further by expelling the envelope of the latter.

From [140], and using M_p as the total mass of the black hole progenitor before common envelope, M_e as the mass of the H envelope, M_{He} as the mass of its core, a_i and a_f as the separation before and after the common envelope phase, and $r_L \equiv R_L/a$ as the dimensionless Roche lobe radius, [28] and [81] obtain

$$\frac{GM_pM_e}{R} = \frac{GM_pM_e}{r_La_i} = \lambda\alpha_{CE}\left(\frac{GM_{He}M_d}{2a_f} - \frac{GM_pM_d}{2a_i}\right), \tag{1.7}$$

which ties the ejection of the envelope of the black hole progenitor with the orbital energy decrease.

Precisely how the common envelope proceeds is not well known while neither the shape parameter of the density profile, λ, nor the efficiency with which the orbital energy expels the envelope during the common envelope evolution, α_{CE}, are well known. However, [81] constrain their product, $\lambda\alpha_{CE}$ to ~ 0.2, from constraints in the orbital-period modifications (by the common envelope) found in the Galactic SXTs.

1.3.3 Tidal Locking

Case C mass transfer and common envelope evolution are important in order to transfer the angular momentum from the orbit to the spin of the black hole progenitor at the right time. However, for the transfer of angular momentum to be effective, the slow rotation of the massive star must start to synchronize with the orbital period as it shrinks during the common envelope evolution.

As noted in 1.3.2, Lee et al. [81] have developed the evolution of such a binary and constrained the free parameters (i.e., the product $\alpha_{ce}\lambda$ of the common envelope efficiency and the density profile) from observed black hole binaries.

Such synchronization occurs gradually as the orbit shrinks when the companion (or secondary) star *peels* away the outer layers of the primary. As the orbit shrinks and the relative size of the stars to the orbital distance grows, tidal forces become more and more important, up to the point where the spin of primary star becomes synchronized with the orbital period (the spin of the secondary star may also become synchronized with the orbit, however this is not be relevant for the collapsar). Throughout this work, this process is also referred to as tidal locking. As the orbital period decreases more and more, the remaining core of the primary gets spun up. We estimate that even the interior of the core should become synchronized for stars with strong magnetic fields [129]. This point will become crucial in estimating Kerr parameters of black holes in binaries where the orbital period is known.

1.3.4 Black Hole Binaries

In our model a binary with a massive star and a smaller companion evolving in a somewhat large orbit interact through Case C mass transfer when the giant finally fills its Roche Lobe. The orbit shrinks as the masses tend to equalize, and the secondary plunges into a common envelope evolution that removes the envelope of the primary at the expense of shrinking the binary orbit. As this takes place, the spin of the primary synchronizes with the orbital period, spinning up. Finally, the stellar wind of the almost bare He core removes the remaining envelope, breaking the Roche Lobe Overflow (RLOF), and the orbital period briefly stabilizes.

However, the core is almost done burning through its fuel reserves and the collapse occurs soon after. The recently spun-up core collapses forming a Kerr black hole which powers a gamma-ray burst and hypernova explosion up. Material is lost from the binary asymmetrically related to the center of mass; therefore, a Blaauw-Boersma kick [12, 17] is produced and the nearly circular orbits turn into elliptical ones. The size of the kick and the ellipticity depends on the amount of mass loss from the primary.

If the system does not lose more than half its pre-explosion mass, the binary survives and in a few 10^5 years the orbit recircularizes because of tidal interactions. However, if more than half the mass of the system is lost, the system becomes unbound.

Single black holes are very difficult to locate as they have no observable signature other than being accidentally detected by their gravitational lenssing or a direct interaction with another object. However, black holes in a binary, and especially in a tight binary, will eventually interact with the secondary star as this one evolves beyond the main sequence. The accretion from the companion will form a hot accretion disk with abundant X-ray emission. Signatures are plentiful and around 18 black hole binaries have been identified and confirmed in our Galaxy (see e.g. [81] and [118]); several more have been found in other, nearby, galaxies.

We identify many of these black hole binaries as relics of gamma-ray burst and hypernova explosions. Nevertheless, the size of the explosions produced by these varies greatly depending on the binary that was produced, as we will explain in the next chapters.

1.4 Kerr Parameter

The Kerr parameter of a black hole can be defined by

$$a_\star = \frac{a}{M} = \frac{Jc}{GM^2} \tag{1.8}$$

where M is the mass of the black hole, J is the angular momentum ($J = I\Omega$) and a is the specific angular momentum ($a = Jc/GM$) of the black hole. $a_\star$ is the ratio between the speed of a point along the equator of the event horizon and the speed of light, so

$$|a_\star| \leq 1. \tag{1.9}$$

$a_\star = 0$ for a Schwarzschild black hole (non-rotating) and $a_\star = 1$ for a maximally spinning Kerr Black hole.

1.4.1 Inner Radius Of The Accretion Disk

Measuring the Kerr parameter of a black hole is a complicated task. Since black holes have an event horizon instead of a surface, no information can be retrieved directly from them, but it rather has to be obtained from their interaction with their surroundings, i.e. with the accretion disk. To complicate matters further, stable orbits do not exist all the way down to the event horizon, there are rather three important orbits that have to be considered close to the inner rim of an accretion disk. The marginally bound (mb) and the marginally stable (ms) orbits (see Figs. 1.1 and 1.2) can be obtained from the energy and the specific angular momentum

$$\tilde{E} \equiv \frac{E}{\delta m} = \frac{r^2 - 2Mr \pm a\sqrt{Mr}}{r\left(r_3^2 Mr \pm 2a\sqrt{Mr}\right)^{1/2}} \tag{1.10}$$

$$\tilde{l} \equiv \frac{l}{\delta m} = \pm \frac{\sqrt{Mr}\left(r^2 \mp 2a\sqrt{Mr} + a^2\right)}{r\left(r^2 - 3Mr \pm 2a\sqrt{Mr}\right)^{1/2}} \tag{1.11}$$

and are defined as follows

$$r_{mb} = 2 \mp a_\star + 2\sqrt{1 \mp a_\star}, \tag{1.12}$$

$$\begin{aligned} r_{ms} &= M\left(3 + Z_2 \mp \sqrt{(3 - Z_1)(3 + Z_1 + 2Z_2)}\right), \\ Z_1 &\equiv 1 + \left(1 - a_\star^2\right)^{1/3}\left[(1 + a_\star)^{1/3} + (1 - a_\star)^{1/3}\right], \\ Z_2 &\equiv \sqrt{3a_\star^2 + Z_1^2}, \end{aligned} \tag{1.13}$$

and an Innermost Stable Circular Orbit (ISCO), which is somewhere in between, depending on the accretion rate (see e.g. Abramowicz et al. [1], Brown et al. [26]). Figure 1.1 is obtained from

$$\tilde{V}^2(\tilde{l}, r) = \left(1 - \frac{2M}{r}\right)\left(1 + \frac{\tilde{l}^2}{r^2}\right) \tag{1.14}$$

the effective potential for a Schwarzschild black hole. Finding the value of R_{ISCO} is highly non trivial, and a great amount of research has been focused

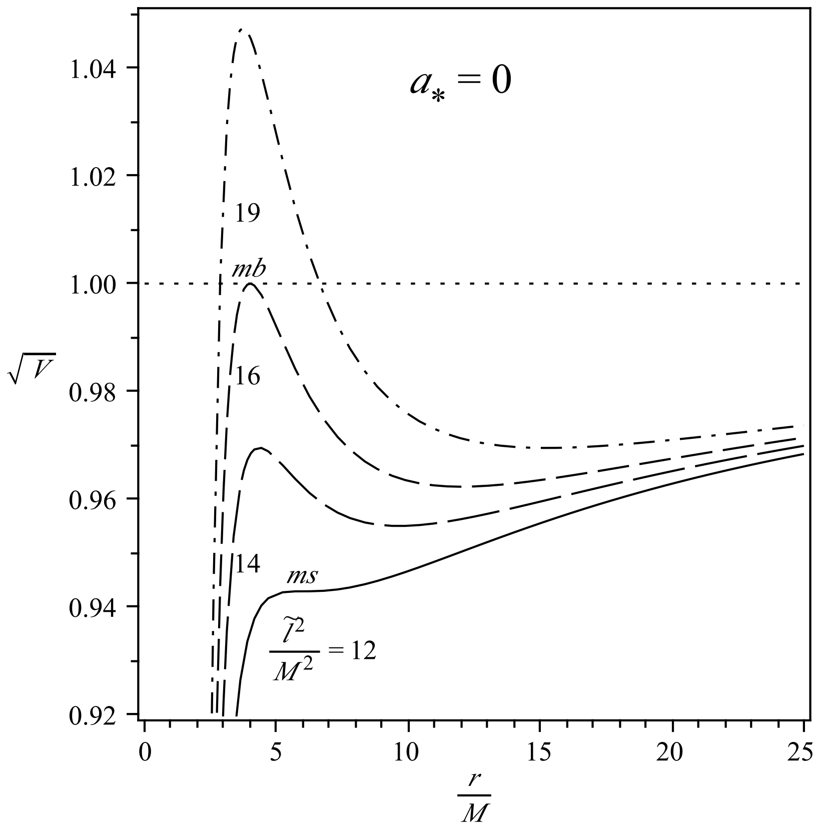

Figure 1.1: The square root of the gravitational effective potential of a non-spinning black hole ($a_\star = 0$) vs the radius for 4 different values of the specific angular momentum. The marginally stable and marginally bound orbits are indicated by ms and mb.

on this particular subject. Nevertheless, for mass accretion rates well below the Eddington limit, R_{ISCO} and Rms are treated as having the same value (this will be the case in the measurements of Kerr parameters). However, this does not hold true for supercritical or hypercritical accretion [1].

As can be seen on Fig. 1.2 the radius of r_{ms} varies with the value of the Kerr parameter of the black hole: For $a_\star = 0$, $R_{ms} = 6M$ (M, the dimensionless mass, is used for GM/c^2 by taking $G = c = 1$), whereas for $a_\star = 1$ $R_{ms} = 1M$ for a direct or co-rotating orbit, and, $R_{ms} = 9M$ for a retrograde or counter-rotating orbit. Further in, the marginally bound orbit: $a_\star = 0$, $R_{mb} = 4M$, whereas for $a_\star = 1$ $R_{mb} = 1M$ for a direct orbit, and, $R_{mb} = 5.83M$ for a retrograde orbit.

Accretion disk material falling closer to the black hole than the marginally

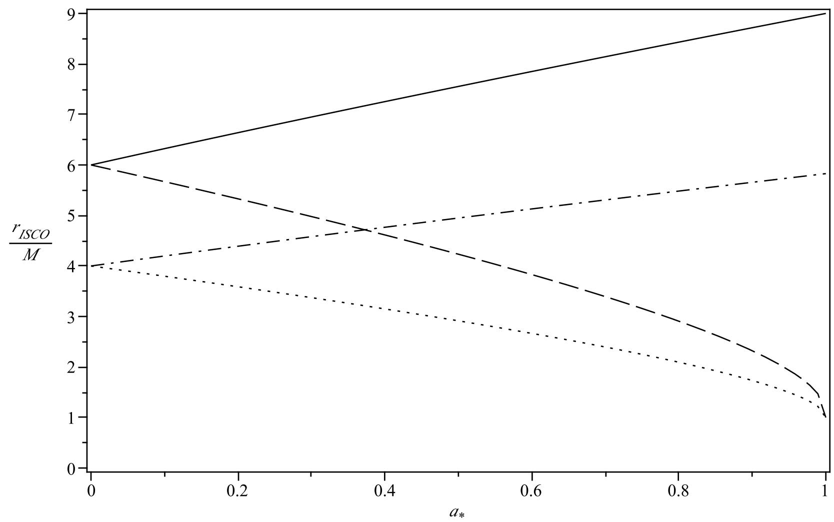

Figure 1.2: The marginally stable (ms) and the marginally bound (mb) radii in the equatorial plane of a rotating black hole. The dashed line represents the co-rotating ms orbit, the solid line represents the counter-rotating ms orbit, the dotted line represents the co-rotating mb orbit and the dash-dotted line represents the counter-rotating mb orbit.

stable will not remain in the accretion disk for too long (it will either fall across the event horizon or be ejected if it does not lie within the Roche lobe [1]), but accreted material inside the marginally bound orbit will fall directly into the black hole. When a slim or thick (not a thin one) accretion disk is present and supercritical or hypercritical accretion is occurring, its ISCO can be pushed closer to the marginally bound orbit.

1.4.2 Observational Methods

In quantum physics, objects can be described by a small number of parameters, e.g., quarks and leptons can be described by their mass, spin, magnetic moment, electric dipole moment, charge (electric, strong and/or weak), etc. Black holes are the only astrophysical objects which can be described by three parameters: mass, spin (Kerr parameter) and charge. At astrophysically relevant scales, for all practical purposes, charge is expected to be zero at all times, given that any deviation from this value would be shorted out by local plasma. That leaves black holes with two parameters to be measured in order to have all the relevant information to describe them. Mass can usually be measured by studying the orbits of binary companions from which mass can be estimated (of course, distance and angle of the orbital plane must also be known). However, only recently have any reliable measurements of spin begun to emerge.

When attempting to measure the rate of rotation on the inner part of the accretion disk, an accurate estimate has to be made of its inner radius if one is to distinguish between an inner boundary with a radius at $1M$, $4M$, $6M$ or $9M$.

There are presently two different methods by which Kerr parameters of stellar mass black holes are being measured. The first one is the fitting of the thermal X-ray continuum. The second one is by modeling the Fe K line profile. However, the use of high-frequency Quasi Periodic Oscillation (QPO, $100 - 400$ Hz) and X-ray polarimetry has been proposed (see [95]) as different alternatives which would allow one to cross-examine the results of observed black holes. However, there is no satisfactory model by which QPO frequencies can be interpreted. All the observations referred to in this work have been been the result of fitting of the X-ray continuum or the Fe K line profile.

The stellar-mass black hole Kerr parameter measurements done by fitting the X-ray continuum used by Gou et al. [60], Liu et al. [86], McClintock et al. [93], Shafee et al. [124] are done using emission only from the thermal state, i.e, black body radiation from the internal region of the accretion disk, during periods with low rms variability and avoiding those with QPOs. One must take into account the Doppler shift, the transverse Doppler shift and beaming from special relativistic effects, plus the gravitational redshift and light bending while interpreting an observation (see, e.g., Fig.3 in [47]). This has the purpose of improving the odds that the inner part of the accretion disk is being observed (and not any other excited material), and that the accretion disk inner radius is that of the ISCO (this is an assumption based on decades of observations).

1.4.3 Observational Results

Observations utilizing the fitting of thermal X-ray continuum of accretion disks around stellar mass black holes in binaries from Gou et al. [60], Liu et al. [86], McClintock et al. [93], Shafee et al. [124] have so far provided us with

Name	M_{BH} $[M_\odot]$	M_d $[M_\odot]$	Measured $a_\star$	P_{Orbit} [days]	Refs
Galactic					
GRO J1655−40	$5.1 - 5.7$	$1.1 - 1.8$	$0.65 - 0.75$	2.6127(8)	(1)
4U 1543−47	$2.0 - 9.7$	$1.3 - 2.6$	$0.75 - 0.85$	1.1164	(2)
GRS 1915+105	14(4)	1.2(2)	> 0.98	33.5(15)	(3)
Extragalactic					
LMC X−1	$8.96 - 11.64$	30.62 ± 3.22	$0.81 - 0.94$	3.91	(4)
LMC X−3	$5 - 11$	6 ± 2	< 0.26	1.70	(5)
M33 X−7	$14.20 - 17.10$	70.0 ± 6.9	$0.72 - 0.82$	3.45	(6)

Table 1.1: Observed parameters of the 6 black hole binaries where $a_\star$ has been measured. REFERENCES: (1)[3], [4], [124]; (2)[106], [108], [124]; (3)[61], [93]; (4)[60], see also chapter 6; (5)[42]; (6)[86].

Kerr parameter measurements on 3 galactic and 3 extragalactic systems as summarized in the table 1.1.

1.5 Objectives

In this work the question of whether theories of stellar and binary evolution can be linked to the values of mass and spin obtained from X-ray observations of black hole binaries is considered.

Lee et al. [81] estimated the Kerr parameters of GRO J1655−40 and 4U 1543−47 to be $a_\star \cong 0.8$ based on binary evolution with Case C mass transfer, followed by a common envelope phase which shrunk the orbital period and synchronized the spin of the primary (the secondary might also have had its period synchronized). This was followed by the collapse of the core of the primary star into a black hole, where they assumed that most of the available angular momentum of the collapsing core was acquired by the black hole.

I will reconstruct the evolution of 15 galactic and 3 extragalactic black hole binaries (including the 6 on table 1.1). From the most current measurements of masses and orbital period for each of the 18 black hole binary systems I will obtain the orbital period at the time of the formation of the black hole. With this I will provide estimates, for each black hole, of the natal values of the Kerr parameter and the available rotational energy (in the Blandford-Znajek mechanism) to power a GRB/HN up. These values need not be the same as the currently measured ones, given that powering up a GRB/HN may partially deplete the energy available in the rotation of the black hole, and/or mass transfer from the companion star may increase the Kerr parameter.

However, the GRB/HN explosion may leave one more signature, the Blaauw-Boersma kick can provide us with further constraints on the mass loss during the explosion. If measurements of peculiar velocities for these binaries are carried out, their evolutionary paths can be narrowed down even further and with it our estimates of the current Kerr parameter and available energies can be more accurate.

Chapter 2

Kerr Parameters for Stellar Mass Black Holes and Their Consequences for GRBs and Hypernovae[1].

2.1 Abstract

Recent measurements of the Kerr parameters $a_\star$ for two black-hole binaries in our Galaxy [124], GRO J1655−40 and 4U 1543−47 of $a_\star = 0.65 - 0.75$ and $a_\star = 0.75 - 0.85$, respectively, fitted well the predictions of Lee et al. [81], of $a_\star \cong 0.8$. In this report we also note that Lee et al. [81] predicted $a_\star > 0.5$ for 80% of the Soft X-ray Transient Sources. The maximum available energy in the Blandford-Znajek formalism for $a_\star > 0.5$ gives $E > 3 \times 10^{53}$ergs, orders of magnitude larger than the energy needed for the GRB and hypernova explosion. We interpret the Soft X-ray Transients to be relics of GRBs and Hypernovae, but most of them were subluminous ones which could use only a small part of the available rotational energy.

[1]A version of this chapter has been submitted by Moreno-Méndez, E., Brown, G.E. Lee, C.-H. and Walter, F.M. to *The Astrophysical Journal.*

2.2 Introduction

Recent estimates have been made for the Kerr parameters ($a_\star \sim 0.75$) for two Soft X-ray Transients (SXTs) [124], GRO J1655−40 (X-Ray Nova Sco) and 4U 1543−47 (Il Lupi). These results facilitate a test of stellar evolution, in that the spins of the black holes in these binaries should be produced in common envelope evolution which begins with the evolving massive giant and companion donor, and ends up in helium-star–donor binary, the hydrogen envelope of the massive star having been stripped off and the helium in the core having been burned.

Lee, Brown and Wijers [81] (hereafter denoted as LBW) assumed common envelope evolution to begin only after He core burning has been completed; i.e., Case C mass transfer [28]. Otherwise the He envelope, if laid bare, would blow away to such an extent that the remaining core would not be sufficiently massive to evolve into a black hole [27]. The black-hole-progenitor star, in which the helium core burning has been completed, is tidally locked with the donor (secondary star) so the spin period of the helium star is equal to the orbital period of the binary. In this tidal locking, LBW assumed uniform rotation of He star by assuming that the inner and outer parts of He are strongly connected due to the presence of a strong internal magnetic field. The C-O core of the helium star drops into a rapidly spinning black hole due to angular momentum conservation. In this process, the spin of the black hole depends chiefly on the mass of the donor because the orbital period chiefly depends on the donor mass as we explain later (in section 2.4). LBW calculated this Kerr parameter ($a_\star$) as a function of binary orbital period. The results are given in their Fig. 12 which we reproduce as Fig. 2.1.

LBW reconstructed the pre-explosion orbital period (see Fig. 2.2) for the Galactic Soft X-ray Transient Sources (SXTs) based on the observed masses and orbital period for the binaries. They assumed conservative mass transfer from evolved companions filling their Roche Lobes onto the black holes. For binaries where the companion was not evolved they estimated that angular momentum could be lost from the system through magnetic braking and gravitational wave radiation. With these reconstructed pre-explosion orbital periods they were able to estimate the original Kerr parameters of the black holes.

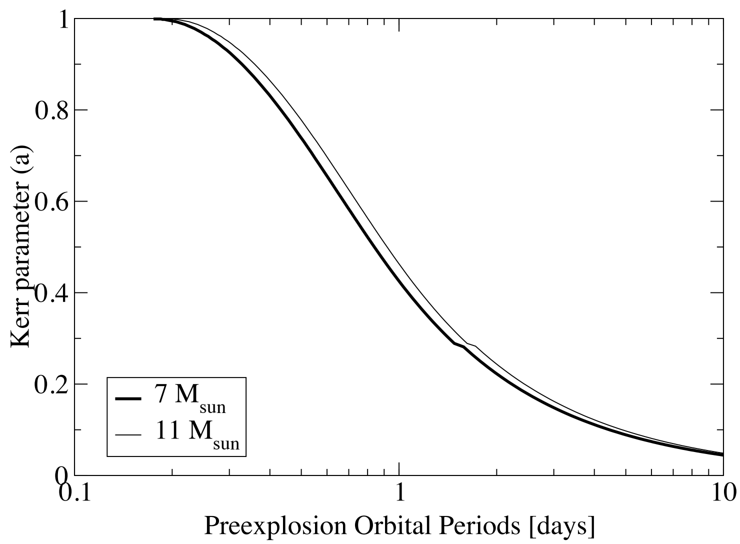

Figure 2.1: The Kerr parameter of the black hole (here represented by a, but in the text referred as $a_\star$) resulting from the collapse of a helium star synchronous with the orbit, as a function of orbital period (LBW). Thick (thin) solid line corresponds to initial $7M_\odot$ ($11M_\odot$) He star. Note that the result depends very little on the mass of the helium star.

In the LBW calculation, the $a_\star$ for GRO J1655$-$40 was slightly greater than for 4U 1543$-$47. Combining this with observation, $0.75 < a_\star < 0.85$ would be our best estimate for both binaries. We also note that LBW predicted $a_\star > 0.6$ for 7 Soft X-ray Transient Sources with main sequence companions and $a_\star \cong 0.5$ for XTE J1550$-$564 (V381 Normae) and GS 2023+338 (V404Cygni) with evolved companions.

The agreement of the natal Kerr parameters with those measured by Shafee et al. [124] means that only a small amount of angular momentum energy could have been lost after the formation of the black hole. The good agreement in this comparison supports the assumption of Case C mass transfer and the tidal locking at the donor-He star stage assumed in the LBW calculations.

The maximum available energy in the Blandford-Znajek [13, 134] formalism for $a_\star > 0.5$ gives $E > 3\times10^{53}$ergs, orders of magnitude larger than observed in any GRB and hypernova explosion. Based on this consideration, we interpret the Soft X-ray Transients to be relics of GRBs and Hypernovae. It should be noted that the way in which the hypernova explodes can be similar to the Woosley Collapsar model [145]. The main advantage in our scenario is that the H envelope in our binary is removed by the donor and the rotational energy is naturally produced in the common envelope evolution. The necessity for Case C mass transfer, given Galactic metallicity, and the measured system velocity (provided by the Blaauw-Boersma kick [12, 17] at the time of the formation of the black hole, [25]) lock us into the Kerr-parameter values we find.

In section 2.3 we show how to determine the Kerr parameters of soft X-ray transient black-hole binaries and the connection of tidal locking to Case C mass transfer. We also discuss the energetics for GRBs and Hypernovae based on the black-hole spin.

In section 2.4 we show that the angular-momentum energy of the black-hole binary is determined mainly by the mass of the donor. We discuss 12 Galactic transient sources with angular-momentum energies $\geq 10^{53}$ergs, all of which are likely relics of GRBs and Hypernovae. The energies of the GRB and Hypernova explosion powered by these, as we shall develop, should be subtracted from the natal rotational energies, to give the current available energy.

2.3 Case C Mass Transfer and Tidal Locking

Case C mass transfer is defined to cover that the mass transfer from the primary to the donor takes place late in the evolution of the former, after the He in the core of the giant progenitor of the black hole has been burned. Aside from the fact that we use Case C mass transfer to achieve the tidal coupling of the donor to the black hole progenitor, the rest of our scenario, especially the collapse, is the same as in the Woosley Collapsar model. Nevertheless Case C mass transfer and tidal locking are essential in order to acquire the necessary spin during the collapse of the He star into the black hole. The evidence that validates our scenario was given in the measured Kerr parameters for GRO J1655−40 and 4U 1543−47 [124] which agreed with the predictions of LBW.

The GRB and Hypernova explosions are all of type I_c, so that He lines do not appear. In the Woosley Collapsar model, He burning is not necessarily finished before explosion, but a) the interacting He may fall into the black hole or, b) the He may not mix with the ^{56}Ni, so that in either case He lines would not be seen.

In Brown et al. [25] the black hole formation was described by a Blaauw-Boersma explosion, which should be sufficient for calculating the system velocity of the binary because conservation laws are respected. However, we believe the Woosley Collapsar model to give a more detailed description of the black hole formation and the hypernova explosion. The MacFadyen & Woosley [88] description includes magnetohydrodynamics in the form of the Blandford-Znajek mechanism. However, mass loss in the explosion and conservation laws are those of Blaauw-Boersma.

2.3.1 Soft X-ray Transients as relics of GRBs and Hypernovae

LBW found that there are two classes of soft X-ray transients (SXTs), those with main sequence companions[2] (denoted as AML, for Angular Momentum Loss), and others with evolved companions (denoted as Nu, for Nuclear evolved). Within the Nus we can further divide into two different groups

[2]Although they are called main sequence, the companions are mostly highly evolved K-dwarfs.

among the observed Galactic SXTs, those with small mass companions ($< 2M_\odot$), and those with companions whose mass was larger than that of the black hole at the time it was formed. We believe the evolution of the three black hole binaries with massive companion is somewhat similar (although Cyg X−1 seems to have a considerably larger companion than XTE J1819−254 or GRS 1915+105), however they are at different stages of mass transfer into the black hole.

There is a gap between the masses of the two groups of SXTs with Nus. Interestingly enough the most energetic GRBs seem to be generated from binaries where the companion has a mass within this gap. This topic will be further discussed below, however it will be pursued in more detail in Chapter 5.

Table 2.1 lists the current best estimates for the masses an periods of the known Galactic black hole binaries separated into AMLs and Nus.

AMLs

Due to the angular momentum loss via gravitational wave radiation and magnetic braking the orbits of AMLs are shortened after black-hole formation. Based on this argument and the current observations, LBW traced back the orbital period at the time of black-hole formation in their Fig. 10. The estimated Kerr parameters for AMLs are $a_\star > 0.6$ (about half of them are $a_\star > 0.8$). The maximum available energies for these systems via the Blandford-Znajek formalism are $E > 3 \times 10^{53}$ ergs, more than necessary to power a GRB and a Hypernova. So we believe that the AMLs are the relics of GRBs and Hypernovae.

Nus

The evolution of Nus after black-hole formation is mainly controlled by the donor which is evolving beyond the main-sequence stage, and the orbit is widened due to the conservative mass transfer from the originally less massive donor to the black hole. We will discuss the possibility that Nu's are also relics of GRBs and Hypernovae.

We will be brief in reconstructing GRO J1655−40 as a relic of GRB and Hypernova because this was constructed in considerable detail with assumed Kerr parameter of $a_\star = 0.8$ in Brown et al. [25]. With the Shafee et al. [124]

Name	$M_{BH,now}$ $[M_\odot]$	$M_{d,now}$ $[M_\odot]$	$P_{Orbit,now}$ [days]	Refs
AML: with main sequence companion				
J1118+480	$6.0-7.7$	$0.09-0.5$	0.169930(4)	(1)
Vel 93	$3.64-4.74$	$0.50-0.65$	0.2852	(2)
J0422+32	$3.4-14.0$	$0.10-0.97$	0.2127(7)	(3)
1859+226	$7.6-12$		0.380(3)	(4)
GS1124−683	6.95(6)	$0.56-0.90$	0.4326	(5)
H1705−250	$5.2-8.6$	$0.3-0.6$	0.5213	(6)
A0620−003	11.0(19)	0.68(18)	0.3230	(7)
GS2000+251	$6.04-13.9$	$0.26-0.59$	0.3441	(8)
Nu: with evolved companion				
GRO J1655−40	$5.1-5.7$	$1.1-1.8$	2.6127(8)	(9)
4U 1543−47	$2.0-9.7$	$1.3-2.6$	1.1164	(10)
XTE J1550−564	$9.68-11.58$	$0.96-1.64$	1.552(10)	(11)
GS 2023+338	$10.3-14.2$	$0.57-0.92$	6.4714	(12)
XTE J1819−254	$8.73-11.69$	$5.50-8.13$	2.817	(13)
GRS 1915+105	14(4)	1.2(2)	33.5(15)	(14)
Cyg X−1	~ 10.1	17.8	5.5996	(15)

Table 2.1: Parameters of the Galactic black hole binaries at present time. Subindex *now* stands for recently measured values. REFERENCES: (1)[90], [139]; (2)[49]; (3)[3]; (4)[50]; (5)[3], [56]; (6)[3]; (7)[3], [57]; (8)[3], [62]; (9)[3], [4], [124]; (10)[106], [108], [124]; (11)[107]; (12)[3], [121], [122]; (13)[107]; (14)[61], [93]; (15)[66].

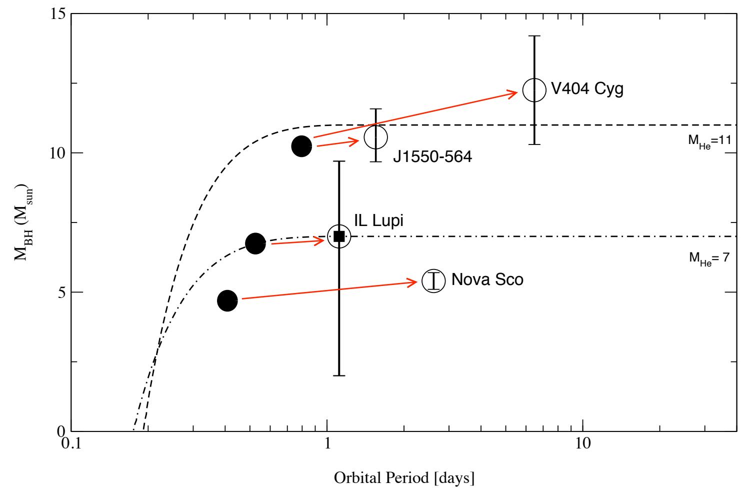

Figure 2.2: Reconstructed pre-explosion orbital period vs. black hole masses of SXTs with evolved companions. The reconstructed pre-explosion orbital periods and black hole masses are marked by filled circles, and the current locations of binaries with evolved companions are marked by open circles. The solid lines are ideal polytropic He stars, but both GRO J1655−40 and 4U 1543−47 were evolved from $11M_\odot$ He stars. This figure is obtained from Fig. 11 of LBW.

measurement of $a_\star = 0.65 - 0.75$ we don't need to change the Brown et al. [25] discussion by much. Indeed, the Shafee et al. [124] $a_\star$ gives its value after powering the GRB and Hypernova explosion, whereas the Brown et al. [25] give the preexplosion value, so the two are not significantly different, since the explosion can be powered by the energy from an $\sim 5\%$ change in $a_\star$ when $a_\star$ is large.

The hypernova aspect of the explosion in GRO J1655−40 was clear by the accretion of α-particle nuclei onto the donor as a result of the hypernova explosion. Israelian et al. [70] found that the F-star donor has O, Mg, Si and S abundances $6 - 10$ times solar. These nuclei were presumably absorbed by the donor, which acts as witness to the explosion.

Due to the similarity in the orbital period between GRO J1655−40 and 4U 1543−47 (Fig. 2.2; see also LBW), we argue that 4U 1543−47 is also a relic of a GRB and Hypernova. Although the black hole masses are not known as well in XTE J1550−564 and GS 2023+338 (V404 Cygni), it can be seen from Fig. 2.1 and Fig. 2.2 that using their reconstructed preexplosion periods they have a Kerr parameter of $a_\star \cong 0.5$, possibly somewhat less definite than the prediction of the $a_\star$ for GRO J1655−40. The latter two binaries have black holes with nearly double the mass as the first two, and, therefore, larger accretion disks. From our arguments in Appendix B, we believe that they may be able to accept more rotational energy, which could be checked by subtracting the measured Kerr parameters from the natal ones.

Brown et al. [25] remarked that for GRO J1655−40 *"After the first second the newly evolved black hole has* $\sim 10^{53}$*erg of rotational energy available to power these (*GRB and hypernova explosion*). The time scale for delivery of this energy depends (inversely quadratically) on the magnitude of the magnetic field in the neighborhood of the black hole, essentially that on the inner accretion disk. The developing supernova explosion disrupts the accretion disk; this removes the magnetic fields anchored in the disk, and self-limits the energy the Blandford-Znajek mechanism can deliver."* This, together with the total rate of creations of binaries of our type of 3×10^{-4}galaxy^{-1}yr^{-1} estimated by Brown et al. [25] will be shown in section 2.4.3 to reproduce the population of subluminous bursts in nearby galaxies. It is considered likely that they are subluminous, at least in the cases of GRO J1655−40 and 4U 1543−47, because their Kerr parameters measured by Shafee et al. [124] are indistinguishable from the natal predicted $a_\star = 0.8$ within observational errors. The disruption of the black hole disk is discussed in detail in Appendix B.

Our evolution of black hole binaries in our Galaxy might appear to be irrelevant for the long (high luminosity) γ-ray bursts because Fruchter et al. [51] show that these come chiefly from low metallicity, very massive stars in galaxies of limited chemical evolution, quite unlike our Milky Way. However, we can construct a quantitative theory of the rotational energies of the black holes which power the central engine for the GRBs and Hypernovae in our Galaxy because we can calculate the black hole Kerr parameters. Having this quantitative theory it is straightforward to apply it to explosions in low metallicity galaxies. There the high luminosities in the GRBs result because the donors are more massive than those in our Galaxy, therefore being able to accept more rotational energy from the binary before disrupting the accretion disks as compared to those in high metallicity binaries with low mass donors. However, there may be some overlap of the high metallicity stars with the same mass donor as the low metallicity stars, as we shall discuss. We predict that this will be the case of XTE J1550-564 (V383 Normae), in which the Kerr parameter will be measured (J. McClintock, private communication). Our calculations will be able to obtain the energy used up in the explosion (the difference between our estimated energy and the measured one), which should be nearly that of cosmological (high luminosity) GRBs.

Recently the eclipsing massive black hole binary X−7 has been discovered in the nearby Spiral Galaxy Messier 33[3] [109]. Since the metallicity is ~ 0.1 solar, we believe it to mimic low metallicity stars which are more massive than the Galactic ones. The donor has mass $68.5 M_\odot$ now, possibly $\sim 80 M_\odot$ earlier. We expect that this system may have gone through a dark explosion due to the high donor mass, which implies a low rotational energy at the time of formation of the black hole, as we discuss in Chapter 6 (see also Moreno Méndez et al. [100]).

Evolution of Cyg X−1, XTE J1819−254 and GRS 1915+105

A highly relevant discussion for the XTE J1819−254 (V4641 Sgr) evolution, which will be a template of an earlier GRS 1915+105 development, was given by Podsiadlowski et al. [114], for Cyg X−1 (1956+350). They discussed the latter as if the donor and black hole were very nearly equal in mass which we shall show will happen in the future for Cyg X−1, although its donor is now $\sim 18 M_\odot$ and black hole $\sim 10 M_\odot$. The donor could have been substantially more massive when the black hole was born after common envelope evolution and lost mass through wind.

The donor in the Podsiadlowski et al. [114] scenario had a stellar wind of $3 \times 10^{-6} M_\odot$ yr^{-1} [66] throughout the evolution. Once the mass of the donor became reduced to a mass comparable to the mass of the black hole, the donor established thermal equilibrium and filled its Roche Lobe transferring mass at the rate of $4 \times 10^{-3} M_\odot \mathrm{yr}^{-1}$. Because of the continuing wind loss the donor shrank significantly within its Roche Lobe and the system widened. The donor started to expand again after it had exhausted all of the hydrogen in the core

[3] See also Bulik [32].

and filled its Roche Lobe a second time. In this phase the mass transfer reached a second peak of $\sim 4 \times 10^{-4} M_\odot$ yr^{-1}, where mass transfer was driven by the evolution of the H-burning shell.

The most interesting feature of this calculation was that the system became detached after the first initial time scale because of the stellar wind from the donor. Since the donor is close to filling its Roche Lobe, such a wind may be focussed towards the accreting black hole, as is inferred from the tomographic analysis of the mass flow in Cyg X-1 by Sowers et al. [128].

Sowers et al. [128] decompose the stellar wind of the supergiant into two moments, one representing the approximately spherically symmetrical part of the wind and the second representing the focussed enhancement of wind density in the direction of the black hole. The latter component of the wind transferred mass in an essentially conservative way (although the former would bring about mass loss). We shall use a similar wind in both, XTE J1819−254 and GRS 1915+105, to accrete matter from the donor to the black hole later on. Note that the wind, transferring matter at hypercritical (much greater than Eddington) rate basically shuts off the initial Roche Lobe overflow, because it can transfer mass sufficiently by itself. The second period of Roche Lobe overflow transfer is driven by the evolution of the H-burning shell; i.e., by the secondary star becoming a red giant.

Podsiadlowski et al. [114] say that "irrespective of whether this particular model is applicable to Cyg X−1, the calculation... illustrates that it is generally more likely to observe a high-mass black-hole X-ray binary in the relatively long-lived wind mass-transfer phase following the initial thermal timescale phase which only lasts a few 10^4 yr. In this example, the wind phase lasts a few 10^5 yr, but it could last as long as a few 10^6 yr if the secondary were initially less evolved."

We believe the above scenario to apply not only to Cyg X−1, XTE J1819−254 and GRS 1915+105, but also to LMC X−1 and M33 X−7, binaries with donors initially more massive than the black-hole companion. We find that these all had dark explosions, as we show below, basically because the donor had too high a mass at the time of the explosion to give an energetic GRB (although the companion stars are no longer more massive than the black hole in XTE J1819−254 nor in GRS 1915+105).

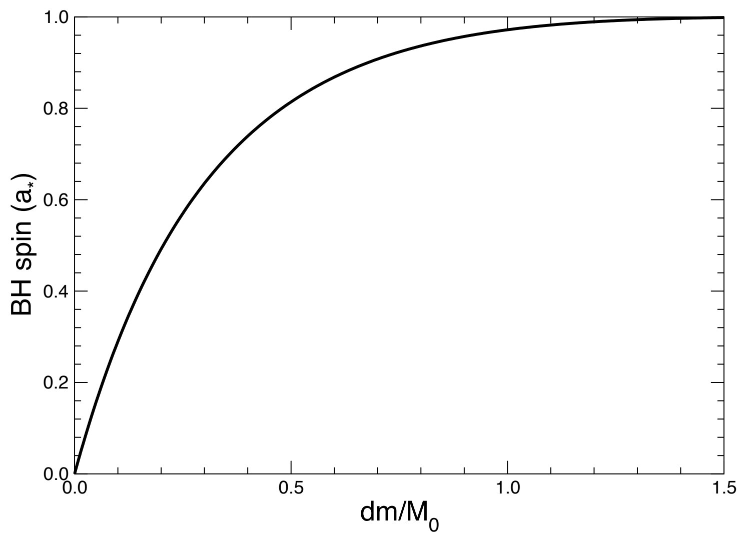

Figure 2.3: Spinning up of black holes. Black hole spin $a_\star$ is given in units of $[GM/c^2]$ and dm is the total rest mass of the accreted material. Note that M_0 is the mass of the non-rotating initial black hole. Here we assumed that the last stable orbit corresponds to the marginally stable radius [25].

We suggest that the wind in these cases may resemble the tidal stream of Blondin et al. [16]. In the two-dimension system studied by these authors, they show that when $D/R_\star$ becomes less than ~ 2, where D is the distance between O-star and black hole and $R_\star$ is the radius of the primary (or accreting) star, the tidally-enhanced-wind accretion exceeds Bondi-Hoyle accretion (steady-state, spherical accretion), a factor that increases to several as $D/R_\star$ decreases.

We are thus now able, given the evolution of Podsiadlowski et al. [114], to describe in detail the crude mass transfer used in the evolution of GRS 1915+105 by Lee et al. [81], as conservative mass transfer in a wind. We disagree, however, with the procedure of Podsiadlowski et al. [114] to limit the accretion to Eddington. In Appendix 8.3 we show hypercritical accretion to be able to occur under these circumstances. In the case of 1915+105, Lee et al. [81] found the average mass transfer rate to be $\dot{M} \sim 10^{-5} M_\odot \mathrm{yr}^{-1}$, about 200 times Eddington (see Bethe et al. [10] p.355).

None of the three binaries we consider had appreciable natal $a_\star$'s because of the high masses of the donors. Thus, we can see that the high Kerr parameter must come from mass accretion (see Fig. 2.3) and that the accreted mass must be slightly greater than the natal mass in order to have such a high Kerr parameter for GRS 1915+105. Taking the final black hole mass to be $14M_\odot$ we need a natal mass of $\sim 6M_\odot$.

If XTE J1819−254 is to be a template for GRS 1915+105, then its black hole should have had a mass of $\sim 6M_\odot$ when born, which would have evolved to its present value of $9.6M_\odot$ by accretion by wind from the donor. Thus, the present $6.5M_\odot$ donor and $9.6M_\odot$ black hole will have essentially interchanged masses from the time of the birth of the black hole until present, through mass exchange from the donor to the black hole. As outlined in LBW, XTE J1819−254 will transfer by wind another $4.6M_\odot$ from donor to black hole to reach the present GRS 1915+105 with black hole mass $14M_\odot$. Thus, the total mass accreted by GRS 1915+105 is estimated to be $\sim 8M_\odot$. The companion mass evolved in this way is $1.9M_\odot$, somewhat larger than the measured [63] mass of $0.81 \pm 0.51M_\odot$. Some mass is, however, lost in jets and winds, which is not taken into account in our approximation of conservative mass transfer.

Our evolution of the black hole birth is similar to the second evolution version of Sadakane et al. [120] which would give the right chemical abundances to the secondary star in XTE J1819−254. These investigators suggested that the black hole mass at birth was $7.2M_\odot$, and to obtain the right surface abundances they proposed that the explosion was a dark one; i.e., one of low energy. Mirabel & Rodrigues [97] suggest that the explosion in Cyg X−1 was a dark explosion, much less energetic than the one in GRO J1655−40. We shall generalize that the explosions are dark because the donors have high masses.

2.3.2 Energetics for Gamma Ray Bursters and Hypernovae

We will construct estimates of the energies available in the spin of the black holes of the transient sources in the next section, but we want to make some general comments here. The principal question with GRBs is whether there is enough angular momentum to power the GRB and Hypernova explosion. In the case of the widely accepted theory, Woosley's Collapsars, this question is unanswered, although one may take the point of view that we observe GRBs and hypernova explosions, so there must be enough angular momentum, which is an integral part of the mechanism. Nonetheless, Heger et al. [65] say "when recent estimates of magnetic torques [130] are added, however, the evolved cores spin an order of magnitude slower. This is still more angular momentum than observed in young pulsars, but too slow for the collapsar model for gamma-ray bursts." Furthermore, the usual scenarios for the interactions of Wolf-Rayets with other stars is that they slow down the rotation.

The above arguments in support of the binary model which was used in Brown et al. [25], were made in Brown et al. [29], and applied to Galactic Transient sources.

The Hypernova formed in 1998bw had $\sim 3 \times 10^{52}$ ergs in kinetic energy [71]. In addition, the jet formation in the GRB requires lifting all of the matter out of the way of the jet. MacFadyen [89] estimates that this costs $\sim 10^{52}$ergs in kinetic energy. At early times the thermal and kinetic energies in supernova explosions are roughly equal, satisfying equipartition. We believe this to be at least roughly true in our explosions here, so that $\sim 6 \times 10^{52}$ergs would be needed for GRB 980425/SN1998bw and possibly more[4], because MacFadyen [89] describes GRB 980425 as a "smothered" explosion.

As noted in Appendix of Brown et al. [29], the Blandford-Znajek efficiency drops substantially as the Kerr parameter decreases below $a_\star \sim 0.5$. Thus, the available rotational energy will decrease rapidly with increasing donor mass. In Appendix 8.3 we will estimate that the highest rotational energy that can be accepted by a binary with a $5M_\odot$ donor is of $\sim 6 \times 10^{52}$ ergs.

[4]In fact, literally, his estimate of $(0.01 - 0.1)M_\odot c^2$ would be $(10^{52} - 10^{53})$ergs. Note that this energy is "invisible", and is not normally included in estimates of GRB energies.

2.4 Donor Mass and Black Hole Spin Anti-Correlation

2.4.1 Mass - Period Relation

Using the relation between the mass of the He core to the star,

$$M_{He} = 0.08(M_{Giant}/M_\odot)^{1.45} M_\odot, \tag{2.1}$$

LBW found that following common envelope evolution

$$a_f = \left(\frac{M_d}{M_\odot}\right)\left(\frac{M_{Giant}}{M_\odot}\right)^{-0.55} a_i. \tag{2.2}$$

Here a_f is the final separation of the He star from the giant following the stripping of its H envelope, and a_i is its initial separation. The He core has inherited the angular momentum of the He star and is tidally locked with the donor. The giant masses found by LBW were all about $30M_\odot$. From Kepler we have the preexplosion period

$$\frac{\text{days}}{P_b} = \left(\frac{4.2R_\odot}{a_f}\right)^{3/2}\left(\frac{M_d + M_{He}}{M_\odot}\right)^{1/2}. \tag{2.3}$$

From Fig. 2.1 one sees that the Kerr parameter increases sharply as the period of the binary P_b decreases. From eq. (2.2) we see that a_f is proportional to M_d, the donor mass, and from eq. (2.3) that P_b is proportional to $a_f^{3/2}$. Estimated Kerr parameters for Galactic sources are summarized in Table 2.2, which makes the dependence of $a_\star$ on the donor clear for the galactic sources.

The increased gravitational binding between donor and the He-star left from the giant after being stripped of hydrogen must furnish the energy to remove the hydrogen envelope, modulo the product of a shape parameter for the density profile (λ) and the efficiency with which orbital energy is used to expel the envelope (α_{ce}): $\lambda\alpha_{ce}$ (see LBW for a discussion on these). The latter decreases inversely with the radius of the giant, so that when the radius is large, the envelope can be removed by a low-mass companion. Combining

Name	M_{BH} $[M_\odot]$	M_d $[M_\odot]$	Initial $a_\star$	$E_{\rm BZ}$ $[10^{51}$ ergs]
AML: with main sequence companion				
J1118+480	~ 5	< 1	0.8	~ 430
Vel 93	~ 5	< 1	0.8	~ 430
J0422+32	6–7	< 1	0.8	$500 \sim 600$
1859+226	6–7	< 1	0.8	$500 \sim 600$
GS1124	$6 - 7$	< 1	0.8	$500 \sim 600$
H1705	$6 - 7$	< 1	0.8	$500 \sim 600$
A0620−003	~ 10	< 1	0.6	~ 440
GS2000+251	~ 10	< 1	0.6	~ 440
Nu: with evolved companion				
GRO J1655−40	~ 5	1–2	0.8	~ 430
4U 1543−47	~ 5	1–2	0.8	~ 430
XTE J1550−564	~ 10	1–2	0.5	~ 300
GS 2023+338	~ 10	1–2	0.5	~ 300
XTE J1819−254	6–7	~ 10	0.2	$10 \sim 12$
GRS 1915+105	6–7	~ 10	0.2 †	$10 \sim 12$
Cyg X−1	6–7	$\gtrsim 30$	0.15	$5 \sim 6$

Table 2.2: Parameters at the time of black hole formation. $E_{\rm BZ}$ is the rotational energy which can be extracted via Blandford-Znajek mechanism with optimal efficiency $\epsilon_\Omega = 1/2$ (see Appendix 8.3). The AML (Angular Momentum Loss) binaries lose energy by gravitational waves, shortening the orbital period whereas the Nu (Nuclear Evolution) binaries will experience mass loss from the donor star to the higher mass black hole and, therefore, move to longer orbital periods. † $a_\star > 0.98$ is the measured Kerr parameter, the difference comes after the accretion of $\sim 8 M_\odot$.

eqs. (2.2) and (2.3) we have

$$\frac{\text{days}}{P_b} = \left(\frac{4.2R_\odot/a_i}{M_d/M_\odot}\right)^{3/2} \left(\frac{M_d + M_{He}}{M_\odot}\right)^{1/2} \left(\frac{M_{giant}}{M_\odot}\right)^{0.83} \tag{2.4}$$

so that $P_b \propto M_d^{3/2}$ for $M_d \ll M_{He}$ and $P_b \propto M_d$ for higher donor masses. As we develop later, the distances a_i at which mass transfer begins depend very little on donor mass so we can use eq. (2.4) in order to scale from one donor mass to another. LBW found all the giants they needed for the transient sources had mass $\sim 30M_\odot$, and, therefore, $M_{He} \sim 11M_\odot$, substantially larger than the donor masses for the most energetic GRBs.

For higher mass donors mass loss in the explosion can be neglected, giving pre and post-explosion periods which are nearly the same.

2.4.2 Explosion Energies of Galactic Black Hole Binaries

In supernova explosions at short times the kinetic and thermal energies are equal, following equipartition. The kinetic energy, which is an order of magnitude or more greater than the GRB energy [89], results from the ram pressure which is needed to clear the way for the jet which initiates the GRB, as just described. We shall assume the kinetic energy to be the same as the thermal energy, which is more easily measured from the observations. Thus, we assume that twice the hypernova energy is needed to power the explosion. On the other hand, the efficiency ϵ_Ω in depositing the energy in the perturbative region is usually taken to be 1/2, the optimum value. This optimum value is obtained by impedance matching as in ordinary electric circuits.

The Blandford-Znajek energy is deposited in a fireball in the perturbative region. Paczyński [111] and Goodman [59] have shown that in order to power a GRB all that is necessary is for the fireball to have enough energy so that the temperature is well above the pair production temperature (assuming the path is clear for the fireball); i.e., $T > 1\text{MeV}$. Then the GRB and the afterglow will follow, just from having as source the localized hot fireball.

The estimated explosion energies of Galactic sources are summarized in Table 2.2. We claim that ours is the first quantitative calculation of the explosions giving rise to GRBs and Hypernovae, in the sense that we calculate the energy that is supplied in the form of angular momentum energy. What the central engine does with this energy is another matter. We cannot calculate how much of the energy is accepted by the accretion disk in detail, because we cannot calculate analytically the Rayleigh-Taylor instability in the magnetic field of the black hole coupling to the accretion disk (but see Appendix 8.3).

We do know from observations that accretion disks have been formed; the Kerr parameters of the black holes have been measured by observing and modeling radiation produced on the inner edge of the accretion disks.

The rotational energy that is not accepted by the black hole does, however, remain in the binary and appears later in the Kerr parameter of the black hole. Thus far, in GRO J1655−40 and 4U 1543−47 a tiny part of the available rotation energy was accepted by the central engine, so little that the final rotational energy could not be discriminated from our calculated initial energy, because the uncertainty in the measurement of Kerr parameters was the same order of magnitude as the explosion energy.

We can understand why AMLs and some Nus (GRO J1655−40, 4U 1543−47, XTE J1550−564 and GS 2023+338) have such rotational energies from Case C mass transfer. The lower the companion mass, the greater the radius R_{sg} that the supergiant must reach before its Roche lobe meets the companion. Given giant radii such as shown in Fig. 2.4, the typical separation distance between giant and companion is $\sim (1200-1300)R_\odot$, higher (because of the Roche lobe of the companion star) than the $1000R_\odot$ the giant radius reaches. The binding energy of the supergiant envelope goes as R_{sg}^{-1} and at such a large distance it can be removed by the change in binding energy of an $\sim 1M_\odot$ donor, spiralling in from $\sim 1500R_\odot$ to $\sim 5R_\odot$. A higher mass donor would end up further out, since it would not have to spiral in so close in order to release enough binding energy. Thus low mass $\sim (1-2)M_\odot$ companions can naturally deposit their increase in gravitational binding energy in removing the high R_{sg} envelopes. A detailed discussion of these matters is given in Brown et al. [24]. Thus Case C mass transfer naturally gives the ultra-high rotational energies of the binaries with low-mass donors discussed as relics of GRB and hypernova explosions in our Galaxy.

Cyg X-1: One can see that there was essentially no mass loss in the Cyg

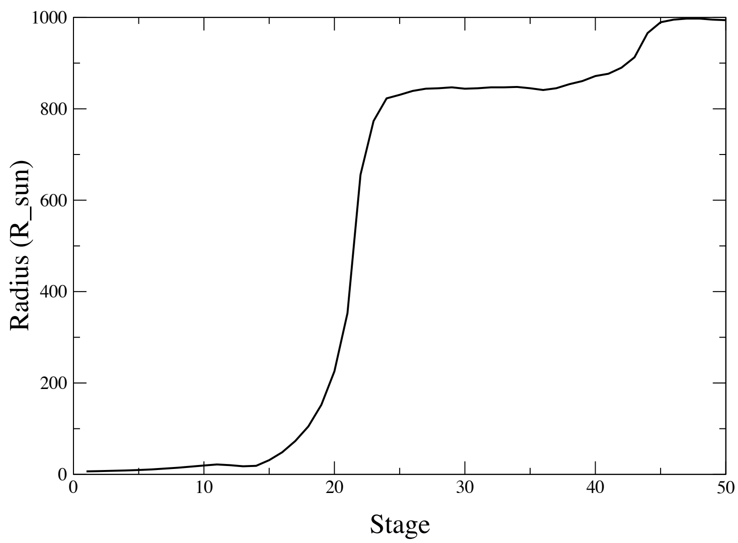

Figure 2.4: Radial expansion of the 25 ZAMS mass star, obtained by Brown et al. [28] from modifying Schaller et al. [123]. Helium core burning ends at stage 43.

X−1 explosion because the space velocity of Cyg X−1 relative to the Cyg OB3 cluster of O-stars of (9 ± 2)km s^{-1} is typical of the random velocities of stars in expanding associations [97]. In comparison, GRO J1655−40 had a very strong explosion from the fact that its space velocity after explosion is (112 ± 18)km s^{-1}, although only a small part of the available energy is used up in the system velocity. In our reconstruction of the Cyg X−1 evolution the explosion takes place when the black hole mass was about $7M_\odot$, $3M_\odot$ less today. The mass transfer from donor to black hole is nonconservative because of the higher mass of the donor. The initial donor would have been substantially more massive than it is today, at least $\sim 30M_\odot$. For a $30M_\odot$ donor we obtain the maximum available energy is $< 10^{52}$ergs.

Mirabel & Rodrigues [97] argue that Cyg X-1 had a dark explosion. As discussed in Appendix 8.3, the efficiency ϵ_Ω can be taken to be 0.5 for the higher $a_\star$, say $a_\star > 0.5$, but it decreases for small $a_\star$, so that for $a_\star = 0.15$ we calculate it to be 0.15 and for $a_\star = 0.2$ we calculate $\epsilon_\Omega = 0.2$. The $(5-6) \times 10^{51}$ergs is clearly not enough for an explosion in Cyg X−1.

Binaries in low metallicity galaxies: As discussed in a previous section, Orosz et al. [109] have recently measured the extra Galactic M33 X−7 in a neighborhood where the metallicity is ~ 0.1 solar. Following our prediction that the donors in low-metallicity galaxies are generally more massive than those of our Galaxy, the now $68M_\odot$ star ($\sim 80M_\odot$ at the time of common envelope evolution) is much more massive. In fact, it is so massive that we estimate $a_\star$ to be < 0.12 at the time of the explosion (but see Moreno Méndez et al. [100] for a discussion on this system); it probably went through a dark explosion.

Far from being irrelevant, as Fruchter et al. [51] imply (because of their lack of dynamics), the measurements of Shafee et al. [124] teach us how to calculate the energies of GRB and Hypernova explosions. Having a dynamical theory, we can easily extend it to low metallicity galaxies, by increasing the donor masses.

2.4.3 Subluminous Bursts

All of the GRBs in our binary model come about from the same mechanism, but their angular momentum energy will be decided by the mass of their donor.

Our mechanism suggests why the binaries are usually left spinning with the measured Kerr parameter. There must be a “Goldilocks” scenario for the energy needed to power a high-luminosity GRB, neither too big nor too small. For the black holes in the Milky Way, with the $(1-2.5)M_\odot$ low mass donors, the available rotational energy is clearly too large. We know this because the calculated initial Kerr parameters were essentially the same as those found by Shafee et al. [124]; thus, very little of the energy had been used up in the explosion. On the other hand we have M33 X$-$7, which we will discuss in more detail later, which had a donor of ZAMS mass $\sim 80M_\odot$, with Kerr parameter $a_\star \leq 0.12$ which probably went into a dark explosion, like Cyg X$-$1. So we have bracketed (but rather widely) the luminous explosions.

Initially in supernova explosions the kinetic energy, the main part of which results from clearing out the matter in the way of the jet that accompanies the GRB, is about equal to the thermal energy of the hypernova. MacFadyen [89] finds this to be at least approximately true and it would follow from equipartition of energy. We should be able to connect GRS 980425 with our galactic GRBs because of its high metallicity, nearly solar [127]. (The galaxy of GRB 980425 is incorrectly put in the class of low-metallicity by [132] and by [137].)

Note that the high luminosity GRBs turn out to be only a small fraction of the total number, even if we use a beaming factor of 100 for them. Thus, they must be formed in very special circumstances (see Appendix 8.3).

The question is whether there are enough transient sources to supply subluminous GRBs in nearby galaxies. Brown et al. [25] estimated that in our Galaxy the total rate of creation of the transient source binaries was $\sim 3 \times 10^{-4}\mathrm{galaxy}^{-1}\mathrm{yr}^{-1}$. Given 10^5 galaxies within 200Mpc this number translates into $3,750\mathrm{Gpc}^{-3}\mathrm{yr}^{-1}$. Liang et al. [85] find a beaming factor typically less than 14; such a beaming factor would reduce our number to $268\mathrm{Gpc}^{-3}\mathrm{yr}^{-1}$, in agreement with that of Liang et al. [85] of $\sim 325^{+352}_{-177}\mathrm{Gpc}^{-3}\mathrm{yr}^{-1}$. This is much higher than their estimated rate of high-luminosity GRBs of $1.12^{+0.43}_{-0.20}\mathrm{Gpc}^{-1}\mathrm{yr}^{-1}$. The usual beaming factor for the high luminosity bursts is ~ 100. Even with such a large factor, the high luminosity GRBs are estimated to be much less in number, by a factor of ~ 40, than the subluminous ones. Although one should add the Woosley Collapsar rate to our binary rate, we have enough binaries to account for all of the bursts. Woosley & Heger [147] estimate that $\sim 1\%$ of the stars above $10 M_\odot$ can, under certain circumstances, retain enough angular momentum to make GRBs.

The effect of cutting down the wind losses in Galactic stars gave a hint about how the rotational energy in the binaries could be decreased so as to be in the ballpark needed for high luminosity GRBs. The winds are particularly high because of Galactic metallicity. Low metallicity stars have lower $\dot{m}$. In general the low Z stars are more massive than Galactic ones, which we believe has the effect of scaling up all of the Galactic masses. We pursue the question of cosmological GRBs and their abundances in Appendix 8.3.

2.4.4 A General Discussion of Black Hole masses

If one accepts the Schaller et al. [123] numbers literally, then Case C mass transfer is actually limited to a narrow interval of ZAMS masses about $20 M_\odot$, $\sim (19-22) M_\odot$ as found by Portegies Zwart et al. [116]. This is because the binary orbit widens with mass loss of the supergiant so that in order to initiate mass transfer only after helium burning the supergiant has to expand sufficiently that this widening of the orbit is compensated for. A graphic display of this is shown in Fig. 1 of Brown et al. [28].

LBW realized that in order to reproduce black holes from the interval of ZAMS masses $(18-35) M_\odot$, necessary for their evolution of transient sources, they had to cut down the wind losses in the red giant stage (by hand -see LBW Fig. 3). This was clearly necessary because Brown et al. [27] had shown that high mass X-ray black hole binaries could be evolved with the black holes coming from ZAMS masses $(18-35) M_\odot$ provided Case C mass transfer was used. It may be that the donor mass for high luminosity GRBs has to be higher than $5 M_\odot$. We do not yet know how rapidly the binaries are left rotating after the explosion. Measurements of black hole binaries with donor masses $\sim (10-20) M_\odot$ would be very helpful.

In LBW it seemed strange that the giant progenitors of the black-hole binaries GRO J1655−40, 4U 1543−47 and GRS 1915+105 all came from $(30-33) M_\odot$ giants, whereas black holes were formed from ZAMS mass $(18-35) M_\odot$ in Case C mass transfer in our Galaxy, and the lower mass black holes are certainly more copius than $(30-33) M_\odot$ ones.

In the case of GRO J1655−40 and 4U 1543−47 the explosion was so energetic that the black hole of $5.5M_\odot$ was only about half of the progenitor He star mass, i.e. the explosion was so violent that nearly half of the mass of the system was lost in the explosion; a loss of half or more resulting in system breakup. Presumably most binaries with lower-mass black holes did lose more than half of their system mass and did not survive the explosion.

GRS 1915+105, XTE J1819−254 and Cyg X−1 did have black holes of $(6-7)M_\odot$, with little mass loss in the explosion, which came from $(20-22)M_\odot$ progenitors (our evolution of GRS 1915+105 in the present paper is an improvement over that in LBW). Thus, the black holes in Galactic soft X-ray transient sources do really come from a wide range of ZAMS mass progenitors.

2.5 Conclusions

Our theory of GRBs and hypernova explosions was developed in Brown et al. [25] and is essentially unchanged. In the meantime we learned in Lee et al. [81] how to calculate the Kerr parameters of the black holes, essentially through an understanding of the tidal locking. Our Kerr parameters have been checked in our Galaxy by the measurement of the Kerr parameters of GRO J1655−40 and 4U 1543−47 by Shafee et al. [124]. Both Paczyński [111] and Goodman [59] have shown that when sufficient energy has been delivered to the fireball (so that the temperature is above the pair-production threshold) the GRB and Hypernova explosions follow and the afterglow is that as observed. In this sense the production of energy and the explosion decouple, but the latter follows from the former once sufficient energy is furnished. In this sense we have a complete and calculable theory of GRBs and Hypernovae.

The GRB and Hypernova explosions are just those of the Collapsar model of Woosley, but with an important improvement; namely, any required amount of rotational energy is obtained from the tidal spin up of the black-hole progenitor by the donor. The donor then, after furnishing the angular momentum, acts as a passive witness to the explosion, but can show some detail of the latter in the chiefly alpha-particle nuclei which it accretes. The magnetic field lines threading the disk of the black hole are well placed to power the central engine in the Blandford-Znajek mechanism. The jet formation and hypernova explosion are powered just as in the MacFadyen & Woosley [88] Collapsar. We check by population synthesis that our binaries are sufficient in number to reproduce all GRBs.

The great advantage that the soft X-ray transient sources in our Galaxy have is that their properties can be studied in detail. They are, however, a special class because of the high metallicity in our Galaxy. Nonetheless, it is easy to extend our galactic description to one of low-metallicity galaxies, because the angular momentum energy is determined by the mass of the donor. Donors in low-metallicity galaxies tend to be more massive than in high-metallicity ones, furnishing a lower rotational energy.

We find that the subluminous GRBs come from two sources: 1) Galactic metallicity systems with low-mass donors, where the magnetic field coupling to the black hole disk is so high that it dismantles the central engine before much angular momentum energy can be delivered. GRO J1655−40 and 4U 1543−47 are excellent examples of these, in that only a tiny part of the angular momentum energy was used up in the explosion. 2) Binaries with massive, low-metallicity donors, going up to the $80 M_\odot$ donor in M33 X−7. Somewhere in between these extremes the binaries will have the rotational energies of the cosmological GRBs.

We estimate the high-luminosity, cosmological GRBs result from a "Goldilocks" phenomenon, being produced only in binaries with donor masses $\sim 5 M_\odot$, but this is uncertain until the Kerr parameters or the system velocity of binaries such as XTE J1550−564 are measured. With such a Kerr parameter (or system velocity) in hand, we can subtract the rotational energy left in the binary from the preexplosion energy which we calculate (see Table 2.2). This will tell us the energy of the explosion. In the cases of GRO J1655−40 and 4U 1543−47, the energy used up in the explosion was tiny compared with the initial rotational energy, but this must change as the initial rotational energy decreases, and black hole mass increases.

We have shown that there are 12 relics of GRBs and hypernova explosions in the Galaxy, and that 3 (XTE J1819−254, GRS 1915+105, Cyg X−1) might have gone through a low-energy dark explosion[5], although the first two of these may have gone through GRBs and hypernova explosions. So we believe that

[5]Called "smothered" explosion by MacFadyen [89].

the soft X-ray black hole binaries are the major sources for the subluminous GRBs.

Chapter 3

GRBs and Hypernova Explosions of Some Galactic Sources[1].

3.1 Abstract

Knowing the Kerr parameters we can make quantitative calculations of the rotational energy of black holes. We show that Nova Sco (GRO J1655−40), Il Lupi (4U 1543−47), XTE J1550−564 and GS 2023+338 are relics of gamma-ray burst (GRB) and Hypernova explosions. They had more than enough rotational energy to power themselves. In fact, they had so much energy that they would have disrupted the accretion disk of the black hole that powered them by the communicated rotational energy, so that the energy delivery was self limiting. The most important feature in producing high rotational energy in the binary is low donor (secondary star)mass.

We suggest that V4641 Sgr (XTE J1819−254) and GRS 1915+105 underwent less energetic explosions; because of their large donor masses. These explosions were one or two orders of magnitude lower in energy than that of Nova Sco. Cyg X−1 (1956+350) had an even less energetic explosion, because of an even larger donor mass.

We find that in the evolution of the soft X-ray transient sources the donor

[1]A version of this chapter was published in Brown, G.E., Lee, C.-H. and Moreno Méndez, E. 2007, *The Astrophysical Journal*, 671, L41.

(secondary star) is tidally locked with the helium star, which evolved from the giant, as the hydrogen envelope is stripped off in common envelope evolution. The tidal locking is transferred from the helium star to the black hole into which it falls. Depending on the mass of the donor, the black hole can be spun up to the angular momentum necessary to power the GRB and Hypernova explosion. The donor decouples, acting as a passive witness to the explosion which, for the given angular momentum, then proceeds as in the Woosley Collapsar model.

High mass donors which tend to follow from low metallicity give long GRBs because their lower energy can be accepted by the central engine.

3.2 Introduction

The hypernova explosions accompanying GRBs are Type I_{bc}; i.e., they show no hydrogen lines and no helium lines. Arguments have been given that the helium lines would not be seen even if the helium were present, that helium would have to mix with ^{56}Ni if the lines were to be seen, etc. Thus, hydrogen is not present at the time of the explosion. As we shall outline, this is the situation in common envelope evolution in Case C mass transfer. Case C mass transfer means mass transfer after the helium burning of the giant is finished. For such a case the GRB and hypernova explosion for Nova Sco (GRO J1655−40) was described by Brown et al. [25]. We can now reconstruct the explosion for this case, since the Kerr parameters ($a_\star$) of Nova Sco and Il Lupi have been measured in the Smithsonian-Harvard-MIT observations [124], with $a_\star = 0.65 - 0.75$ and $a_\star = 0.75 - 0.85$ respectively. They check against the prediction of Lee et al. [81] (denoted as LBW) who found $a_\star = 0.8$ for both. From the $a_\star$ we can construct available energies in the Blandford-Znajek mechanism. We have a simple guiding principle for the sources considered; namely, that the explosion energy depends chiefly on the mass of the donor (secondary star), and this is easily seen if the binaries are evolved in Case C mass transfer, as we shall show.

In similar vein, the GRBs and hypernova explosions can be constructed for XTE J1550−564 and GS 2023+338 (V404 Cygni), the available rotational energy being nearly the same as in Nova Sco.

Moreno Méndez et al. [99] also reconstructed the explosions of GRS 1915+105 and V4641 Sgr. They found the explosion of Cyg X-1, in agreement with Mirabel & Rodrigues [97], to be a dark explosion; i.e., orders of magnitude less explosive than Nova Sco.

3.3 Role of Donor Star in Common Envelope Evolution

Using the relation between the He core mass ($M_{\rm He}$) of a giant after finishing H-core burning and the initial giant mass ($M_{\rm giant}$),

$$M_{\rm He} = 0.08(M_{\rm giant}/M_\odot)^{1.45} M_\odot, \tag{3.1}$$

LBW found that following common envelope evolution,

$$a_f \simeq \frac{M_d}{M_\odot} \left(\frac{M_{\rm giant}}{M_\odot}\right)^{-0.55} a_i. \tag{3.2}$$

Here a_f is the final separation of the He star which remains from the giant following the strip off of its H envelope, and a_i is its initial separation, M_d is the mass of the donor (secondary star). Noteworthy about eq. (3.2) is that the main dependence of the final separation a_f is on the donor mass M_d, only roughly as the square root of $M_{\rm giant}$.

The He star remainder of the giant and the donor are tidally locked at the end of common envelope evolution (LBW). The tidal locking ends here, the Kerr parameter of the black hole being determined by its angular momentum at formation, minus the decrease from angular momentum spent in powering explosions.

From Kepler we have for the preexplosion period

$$\frac{\rm days}{P_b} = \left(\frac{4.2 R_\odot}{a_f}\right)^{3/2} \left(\frac{M_d + M_{\rm He}}{M_\odot}\right)^{1/2} \tag{3.3}$$

where M_d and $M_{\rm He}$ are the masses at the time of common envelope evolution. Given P_b we can easily find the Kerr parameter $a_\star$ from Fig. 12 of LBW, reproduced as Figure 2.1 here. In the case of Nova Sco, $P_b = 1/4$day, $a_f =$

$5.33R_\odot$, $M_{\rm He} = 11M_\odot$ and $M_d = 1.91M_\odot$ (LBW).

The big advantage that Case C mass transfer has is that it not only produces an explosion with no hydrogen envelope, but it produces a great deal of angular momentum, as quantified in the Kerr parameter of the black hole, to power the GRB and Hypernova. The angular momentum results from the tidal locking of the donor and the He star, the latter falling into the black hole. In the core, the helium is burned before common envelope evolution into carbon and, rather quickly, oxygen. The strong $\vec{B}$-field lines, which at one end thread the disk of the black hole as it is formed from the collapse inwards of the ionized metals, are frozen at the other end in the metals and lock the disk tidally with these metals which constitute what is left of the original helium star. If we replace the helium star in the MacFadyen & Woosley [88] paper by our He star then the formation and spin up of the black hole is as these authors described. Thus we basically have a collapsar with high angular momentum that has been spun up by tidal locking with the donor. Note that there is no hydrogen envelope of the giant left, the hydrogen having been expelled in common envelope evolution.

3.4 GRBs and Hypernovae from Soft X-ray Transients With Evolved Companions

In Table 2.2 we list the black hole masses, and our estimates of donor masses, all at the time of the end of common envelope evolution when the tidal locking was established between donor and helium star. These came from LBW and from Moreno Méndez et al. [99]. The Kerr parameters are changed into energies using the Blandford-Znajek formulas [80]

$$E_{BZ} = 1.8 \times 10^{54} \epsilon_\Omega f(a_\star) \frac{M_{BH}}{M_\odot} \text{ergs} \tag{3.4}$$

where the efficiency $\epsilon_\Omega = \Omega_F/\Omega_H$ for energy deposition in the (perturbative) fireball is taken to be 1/2 (for optimum impedance matching) and

$$f(a_\star) = 1 - \sqrt{\frac{1}{2}(1 + \sqrt{1 - a_\star^2})}. \tag{3.5}$$

We note that Cyg X−1 (1956+350y) probably went through a dark explosion [97] meaning that at most, a very low energy, one or two magnitudes less than in the case of Nova Sco. The high Kerr parameter ($a_\star > 0.98$) for GRS 1915−105 [93] came chiefly from mass accretion following the explosion in which the black hole was born, and, therefore, had no influence on the GRB [99]. The measured Kerr parameters are the present ones, and the energies to produce the GRB and Hypernova should be subtracted from our calculated ones.

What we see from Table 2.2 is that the transient sources Nova Sco, Il Lupi, XTE J1550−564, and GS 2023+338 clearly had enough rotational energy to power both a GRB and Hypernova explosion. Brown et al. [25] in discussing these for Nova Sco suggested that the energy was so great that the explosion disrupts the accretion disk; this removes the magnetic fields anchored in the disk and self-limits the energy the Blandford-Znajek mechanism can deliver (see the appendix). In addition to the 7 sources in Table 2.2, LBW worked out the Kerr parameters of the 8 Galactic X-ray Transient sources with main sequence companions, all of which had $a_\star$'s of $0.6 - 0.8$ which correspond to spin energies of $430 - 600 \times 10^{51}$ ergs.

In Brown et al. [25] the GRB and hypernova explosion were reconstructed in all detail. The F-star donor in Nova Sco bore witness to the hypernova explosion through the α-particle nuclei deposited on it. In particular, a large amount of Sulfur, which Nomoto et al. [105] found typical of differentiating hypernovae from the more usual supernovae, was found. The Kerr parameter of 0.8 found by LBW for the preexplosion spin was, within uncertainties, the same as the post explosion Kerr parameters measured by Shafee et al. [124]. The GRB was, of course, not recorded, but the rotational energy was tremendous so that the GRB was either just begun or the accretion disk was smashed immediately. The system velocity was worked out. Almost all of the natal angular momentum energy is still in the system, as measured by Shafee et al. [124], meaning that very little was accepted for the explosion.

3.5 Subluminous GRBs

In Brown et al. [25] the population synthesis suggested a soft X-ray transient birth rate of 3×10^{-4} sources per year per galaxy, which with 10^5 galaxies within 200 Mpc translates into 3750 $\text{Gpc}^{-3}\text{yr}^{-1}$. If we consider the beaming factor of $\sim$ 10%, this is the same rate as the rate of subluminous sources investigated by Liang et al. [85], estimated at 325^{+352}_{-177} $\text{Gpc}^{-3}\text{yr}^{-1}$. The latter are thought to have come from low-metallicity galaxies, but it is none the less interesting that the rate of hypernovae from soft X-ray transient sources is the same as that of the subluminous bursts, especially because we have shown that only a small part of the black hole spin energy in soft X-ray transient sources went into the explosion, so that they would tend to be subluminous.

The question of central engine for GRB060218 was tackled by the 119 astronomers who signed the 5 papers in Nature [34, 91, 112, 126, 148]. From the Supplementary Information of Mazzali et al. [91] one finds that the explosion 2006aj was Type I_{bcd} in nature; i.e., in addition to no hydrogen lines, no helium nor carbon. The only place where this could occur was in a black hole in which convective carbon burning ceases because the carbon abundance drops below 15%: see Fig. 1 of Brown et al. [27]. This leaves no doubt but that the central engine was powered by a black hole, one of low mass.

Galactic GRBs (GRBs from galaxies with solar metallicity) must be subluminous, relatively little of their tremendous rotational energy being used up in the explosion. For the population of metal poor subluminous GRBs one would expect their donors to be more massive because of their low metallicity. Because of the more massive donors they will have less rotational energy, which may be all utilized in the explosions or, at least, will take larger to dismantle the disk. Thus they would be of relatively long duration, but subluminous in the integrated energy in the explosion.

Recently an eclipsing binary M33 X-7 was discovered in a metal poor neighborhood ($\sim$ 10%solar) by Orosz (2007). This can be evolved like a more massive Cyg X-1, but with the advantage that one knows the donor to be $\sim 80M_\odot$ at the time of explosion. The Kerr parameter was $a_\star = 0.12$ and the angular momentum energy of $\sim 10^{52}$ ergs was too little to both power the jet for a GRB and the hypernova, so the explosion was probably "dark." Had the donor been less massive, according to our arguments, then with more energy the GRB and hypernova could have been powered. We agree that the subluminous bursts come chiefly from metal poor galaxies [132], giving the dynamical reason that they have low angular momentum energies because of larger donor masses.

3.6 Discussion

We show that the rotational energy of black holes in soft X-ray transient sources is greatest when the donor in the binary is of low mass. In the case of large donor masses, the rotational energy in the black hole binary is lower.

One can see that Nova Sco had a very high explosion energy from the fact that its space velocity after the explosion is 112 ± 18 km s^{-1} as to compare with Cyg X$-$1 relative to Cyg OB3 in the cluster of O-stars of 9 ± 2 km s^{-1}, which is typical of the random velocities of stars in expanding associations [97].

The explanation of why the angular momentum energy is so high in Nova Sco was given on p.176 of Bethe et al. [10]: *"The massive star will have evolved through its supergiant (He core burning stage) before matter overflows its Roche lobe. Then, by that time, a main sequence companion must be at just the right distance to receive the overflow; this means* $a \sim 1500R_\odot$*, the Roche lobe of the massive star being at* $\sim \frac{2}{3}a$*. Since the binding energy of the envelope of the massive star goes as* $1/a$*, this binding energy is very small, so that the envelope can be removed by the drop in gravitational energy of an* $\sim 1M_\odot$ *main sequence star as it moves inwards in common envelope evolution with the massive star from* $\sim 1500R_\odot$ *to the much smaller Roche lobe of the He star which results when the H envelope is removed from the massive star. In this way one could understand why all of the main sequence companions of the black holes in the transient sources were of nearly the same low masses,* $(0.5-1)M_\odot$*."* For the companions with masses $10M_\odot$, the necessary drop in gravitational energy is only 1/10 that of the $1M_\odot$ companion, so the final a_f can be an order of magnitude greater. The result is an order of magnitude lower rotational energy.

From the above explanation we see that the ultrahigh rotational energies in soft X-ray transient sources are a result of the low donor masses. The rotational energy drops roughly inversely with mass so we would expect it to be an order of magnitude less for stars of low metallicity whose masses are roughly an order of magnitude greater than stars in our Galaxy. Thus, cosmological GRBs will not have so much rotational energy as to dismantle the disk, and may be able to furnish their rotational energy to the GRB and Hypernova. At least, now that we understand why Galactic GRBs are so energetic, we can offer reasons why the cosmological GRBs have lower energy, but may be able to use up more of it in the explosion.

Measurement of the Kerr parameter for XTE J1550-564 (J. McClintock et al., Smithsonian-Harvard coalition, in progress) will enable us to say how much of the natal ~ 300Bethes was used up in the explosion.

3.7 Summary

In summary, the essential points of our paper are that the Woosley Collapsar model can be obtained from our Case C mass transfer, but with the black hole having any desired angular momentum, by making choice of the donor mass. Because the helium is burned preceding the explosion in Case C mass transfer, the ashes of the central helium, carbon and oxygen, fall first into the black hole and ensure the tidal locking through the strong B-field lines which are frozen in the ionized metals.

Our results for the LBW calculation of Kerr parameters have been confirmed by the Smithsonian-Harvard group. Given the Kerr parameters, we can make quantitative calculation of the spin energy of black hole. We give predictions for the Kerr parameters of 12 Galactic black hole sources which have not yet been measured.

We note that the rotational energy of M33 X-7 was lower than that of cosmological GRBs and suggest that these originate from low metallicity donors of somewhat less mass than that of M33 X-7. Our suggestion that XTE J1550-564 should have the angular momentum energy in its explosion, which is as large as that of cosmological GRBs, should soon be tested by the measurement of the Kerr parameter.

Chapter 4

LMC X-3 May Be a Relic of a GRB Similar to Cosmological GRBs[1].

4.1 Abstract

The present scenario for high-luminosity long γ-ray bursts is strongly influenced by the paper of Fruchter et al. (2006). Whereas the main contention of this paper that these GRBs occur in low-metallicity irregular galaxies is based on a considerable collection of observational results and although the main thesis is doubtless correct, the paper does not explain the dynamics that produces such GRBs and much of the discussion not directly concerning the main thesis is incorrect. We propose a dynamics and elucidate how the Fruchter et al. [51] results may be tested, in our neighborhood in the LMC, suggesting that LMC X-3 is a relic of a high luminosity explosion, probably accompanied by a GRB and hypernova explosion. The way to test our suggestion is to measure the system velocity of the present black hole. We correct errors of the Fruchter et al. paper in stellar evolution, so that the study of GRBs is consistent with it. We show that the subluminous GRB 060218 had a low-mass black hole as central engine.

[1] A version of this chapter was published in Brown, G.E., Lee, C.-H. and Moreno Méndez, E. 2008, *The Astrophysical Journal*, 685, 1063.

4.2 Introduction

Fruchter et al. [51] collect an important amount of data on long γ-ray bursts. They propose that the long-duration γ-ray bursts are associated with the most extremely massive stars and that they may be restricted to galaxies of limited chemical evolution. Also that long γ-ray bursts are relatively rare in galaxies such as our own Milky Way.

In recent papers [29, 99] we have developed the Blandford-Znajek model of GRBs into quantitative calculations of the angular momentum energy that can be delivered for GRBs and hypernova explosions. This was possible because Lee et al. [81] showed how to work out the Kerr parameters of the rotating black holes. We showed that the Galactic soft X-ray transient sources were relics of Galactic explosions and constructed the energy of all 15 of the known Galactic sources.

It turns out that most of the Galactic sources underwent subluminous GRBs, not because they did not possess sufficient rotational energy, but mostly because the rotational energy was so large that it destroyed the accretion disk so quickly that the central engine was dismantled before the GRB could properly develop. We give as examples the transient sources Nova Sco (GRO J1655−40) and Il Lupi (4U 1543−47) for which Lee et al. [81] had predicted Kerr parameters of $a_\star = 0.8$ and which were measured, by Shafee et al. [124], at the present time, to be $a_\star = 0.65 - 0.75$ and $a_\star = 0.75 - 0.85$, respectively. The natal rotational energy was 430 bethes (one bethe= 10^{51} ergs) and the final (measured) energy is indistinguishable, within errors, from the natal energy. Brown et al. [25] had reconstructed the GRB and hypernova explosion for Nova Sco. The nature of the explosion could be reconstructed from the donor, which accepted a number of α-particle nuclei, especially ^{32}S which is special for hypernova explosions, but rare in supernova explosions.

Thus, from the near equality of the natal and present rotational energies, only a few percent of the available rotational energy could have been accepted in the GRB and hypernova explosion.

Moreno Méndez et al. [99] and Brown et al. [29] showed from population syntheses that the soft X-ray transient sources were sufficient in number to account for all of the subluminous bursts in our neighborhood.

What about the cosmological bursts which are the long γ-ray bursts considered by Fruchter et al. [51]? The Woosley Collapsar model was invented to describe these. In terms of numbers these are only a few percent of the subluminous bursts [85]. We will not take issue with the Woosley Collapsar model describing these, because our binary evolution begins by the donor spinning the He star, progenitor of the black hole, to whatever angular momentum is needed to power GRBs and Hypernovae. Then the donor decouples, acting only as a witness to the explosion, with the He star collapsing into a black hole in the same way as in the Woosley Collapsar model, the tidal locking between the donor and the He star being transferred to synchronization between the donor and the black hole. From the latter, Lee et al. [81] predicted the Kerr parameter of the black hole.

The main point developed by Brown et al. [29] was that the rotational energy of the binary is roughly inversely proportional to the mass of the donor. This follows from Kepler's law and from the fact that in Case C mass transfer (following the He burning), the initial a_is (distance between the giant and the companion) of the binaries are roughly equal. The fact that, according to Fruchter et al. [51], the long γ-ray bursts are in low-metallicity galaxies, does not elucidate the dynamics which produce the long bursts. The dynamics result from the fact that low-metallicity galaxies tend to have stars of higher mass than Galactic. The higher mass of the donors slows the binary down sufficiently that the rotational energy can be accepted by the central engine. In other words, the question of energy is a "Goldilocks" one. It must be not too much, because in that case the central engine will be dismantled, and not too little, because that would only be sufficient for a subluminous burst, but for a long high luminosity γ-ray burst it must be just right.

We have learned enough about Kerr parameters from our calculations and from the Smithsonian-Harvard measurements to construct a "guesstimate". Namely, we believe that LMC X$-$3 underwent the closest (in energy) explosion in our Galaxy, to a cosmological GRB. (Of course, one can say that LMC X$-$3 is not in our Galaxy, but in the LMC.)

We believe that in its $\sim 1/3$ solar metallicity [119], it tends towards the low-metallicity stars considered by Fruchter et al. [51], so that at the least it has somewhere between Galactic and the low metallicity favored there. There is uncertainty in the masses, but Davis et al. [42] have a value of $a_\star \simeq 0.26$ for the present Kerr parameter of LMC X-3. They took $7M_\odot$ as the mass of the black hole, which would imply a donor of $\sim 4M_\odot$ from the lower end of the measurements of Cowley et al. [41]. We take these to be representative; other investigators have found other masses, so we suggest our evolution as only a possible one.

All of the binaries in Brown et al. [29] and Moreno Méndez et al. [99] which had so much energy that they dismantled the black hole accretion disk had donor masses of $(1-2)M_\odot$, so the donor in LMC X$-$3 is at least double those masses. The donor mass is close to the $\sim 5M_\odot$ that Moreno Méndez et al. [99] estimated would give the energy of a cosmological GRB. In any case, LMC X$-$3 is the closest "nearby" relic of binaries similar to the progenitors of cosmological GRBs. We can, therefore, use it as an example to try to reconstruct the explosion and model the energy of the explosion. We estimate the mass loss in the explosion, finding it to be substantial. It is likely that the estimated system velocity can be measured, at least the radial component of it, which should test our prediction.

In this paper we wish to also summarize results of earlier calculations which have a direct bearing on the 5 papers in the 31 August 2006 Nature [34, 91, 112, 126, 148], in order to show that useful evolutions of black holes have been carried out in the past, of which the 119 astronomers who signed these articles were unaware. We show that these previous calculations have a direct bearing on the measurements of GRB 060218/SN 2006aj; namely that the central engine was a black hole, not the magnetar conjectured by most of the authors, and that the black hole was one of minimum black hole mass with an $\sim (18-20)M_\odot$ ZAMS (Zero Age Main Sequence) progenitor. We also correct a number of errors in Galactic black hole evolution in the Fruchter et al. [51] article.

4.3 Evolution of Black Holes in Our Galaxy

We begin by expanding on the evolutionary discussions in Brown et al. [27]. The history of black hole evolution in binaries (which is the only place where black holes could be studied in detail) was that whatever mass, within reasonable limits, one proposed, the binary would turn out to have had a neutron star rather than black hole as compact object. The first clear explanation of this was given by Brown et al. [27]; namely, the evolution of black holes in our Galaxy depends upon binarity. Namely, in Case A or B mass transfer (mass transfer while the giant star is in main sequence or red giant stage) the strong winds in our Galaxy blow off sufficient of the "naked" He envelope so that the remaining core of metals was too light to evolve into a black hole; rather, it would end up as a neutron star.

Only in Case C mass transfer, if mass transfer following He burning were carried out, would the remaining core have the possibility of evolving into a black hole. Now, just what the limit is for the lowest ZAMS mass star that will evolve into a black hole is determined by what we call the Woosley Ansatz. In our opinion this is one of the most powerful developments in stellar evolution. We will combine this "Ansatz" with the Bethe et al. [6] considerations of entropy in the Fe core.

The Woosley Ansatz basically divides the burning of ^{12}C into low energy, $T \sim$ 20keV burning through the $^{12}\mathrm{C}+\alpha \rightarrow {}^{16}\mathrm{O}$ process and the $T \sim$ 80keV $^{12}\mathrm{C}+{}^{12}\mathrm{C}\rightarrow{}^{24}\mathrm{Mg}$ etc. process. The ^{12}C is produced by $\alpha+{}^{8}\mathrm{Be}^{\star} \rightarrow {}^{12}\mathrm{C}$; i.e., essentially through $\alpha + \alpha + \alpha \rightarrow {}^{12}\mathrm{C}$. This is a three body process, going as the square of the density, ρ^2. The ^{12}C is burned into ^{16}O by the two body process $^{12}\mathrm{C}+\alpha \rightarrow {}^{16}\mathrm{O}+\gamma$, which goes as ρ. As long as ^{12}C is present, the latter reaction will take place. However, with increasing M_{ZAMS}, the density decreases. The entropy, which goes inversely with the density, is known to increase with M_{ZAMS}. Therefore, there will be a value of M_{ZAMS} where the ^{12}C is removed by the $^{12}\mathrm{C}+\alpha \rightarrow {}^{16}\mathrm{O}$ as rapidly as it is formed by the $\alpha + \alpha + \alpha$ process. At this value of M_{ZAMS} the $^{12}\mathrm{C}+{}^{12}\mathrm{C}\rightarrow{}^{24}\mathrm{Mg}$ etc. shuts off, because there is no ^{12}C; actually, it shuts off once the central carbon abundance is less than $\sim 15\%$ because there is not enough carbon for convective (steady) burning. At this point the burning processes are all of low temperature, $T \sim$ 20keV. At this point the metallicity is close to zero, independent of the average metallicity of the star.

Now, this temperature is too low for neutrino-pairs to carry off appreciable energy and entropy (the relativistic ν, $\bar{\nu}$ pair cross section goes as T^{11} power), whereas copious amounts of entropy were carried off by $^{12}\mathrm{C}+{}^{12}\mathrm{C}\rightarrow{}^{24}\mathrm{Mg}$ etc. What happens to the entropy which increases with increasing M_{ZAMS}?

Bethe et al. [6] showed that the entropy per nucleon in the Fe core of a star in advanced burning was ~ 1 in units of k_B. The only place the increase in the entropy can go is into an increase in the number of nucleons in the Fe core, once the burning is confined to the low-temperature ($T \sim 20$keV) region. Thus at the ZAMS mass at which the two-body process takes over completely from the three-body process, central abundance of ^{12}C has decreased below 15% and the increase in entropy comes from the Fe cores increasing rapidly with $M_{\rm ZAMS}$.

Given the Woosley value for the ^{12}C+$\alpha \rightarrow^{16}$O process of 170keV barns (at $E = 100$keV), to be correct, the best experiments [77] obtaining 165 ± 50keV barns, the ^{12}C abundance drops below 15% just around $18 M_\odot$, the mass of the progenitor of SN1987A, Sanduleak 69°202. We reproduce Fig.1 of Brown et al. [27] as our Fig 4.1. Thus with Case C mass transfer the threshold for black hole production is ZAMS $18 M_\odot$. We show the calculated compact object masses from Brown et al. [27] as Fig 4.2. This is independent of metallicity, depending only on the low-energy burning. In fact, the metallicity is essentially zero at the threshold in ZAMS masses for black hole formation.

Fryer [53] and Fryer [54] independently obtained a ZAMS mass of $\sim 20 M_\odot$ to give the lowest mass black hole. Our use of the Woosley Ansatz is not only a simple, elegant argument, based on the behavior of entropy, but also a practical one. For example, it settles the questions of central engine in GRB060218/SN2006aj, which was debated by the 119 observers who wrote 5 papers: Campana et al. [34], Mazzali et al. [91], Pian et al. [112], Soderberg et al. [126] and Young [148], mostly speculating that the central engine was a magnetar, although they gave no mechanism by which it could unwind its magnetic field to power the GRB. If one looks at the additional information of Mazzali et al. [91], one sees that there were no carbon lines. (Actually, these lines need not be zero, but some weak ones can come from carbon shell burning.) Therefore, the explosion was I_{bcd} in nature. This occurs only in one venue in astronomy; namely, in a black hole formed at the lowest ZAMS mass possible, $\sim 18 M_\odot$. Thus, clearly the central engine was a rotating black hole. The hypernova had a low energy of 2 bethes [91] and the GRB was highly subluminous. However, from equipartition of energy we expect the kinetic energy to be roughly equal to the thermal energy, so the kinetic energy must

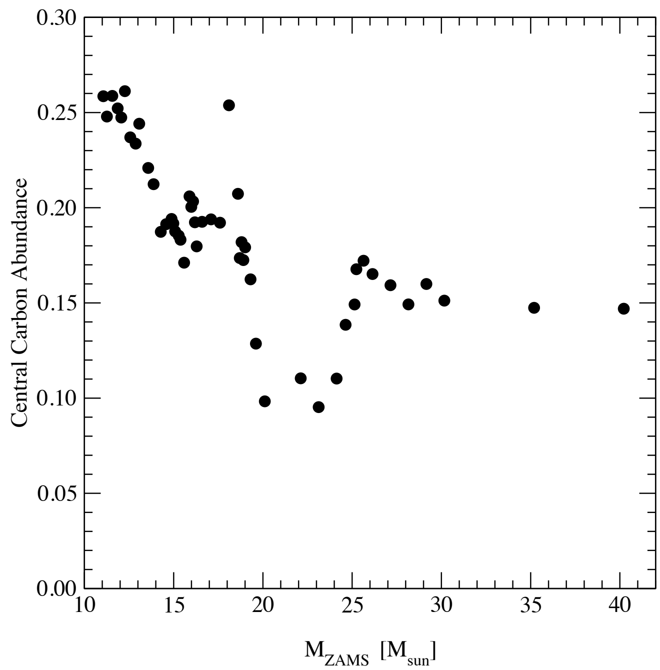

Figure 4.1: Central carbon abundance at the end of He-core burning for "clothed" (single) stars as function of ZAMS mass. The rapid drop in the central carbon abundance at ZAMS mass $M_{\rm ZAMS} \sim 20 M_\odot$ signals the disappearance of convective carbon burning. The resulting iron core masses are summarized in Fig. 4.2.

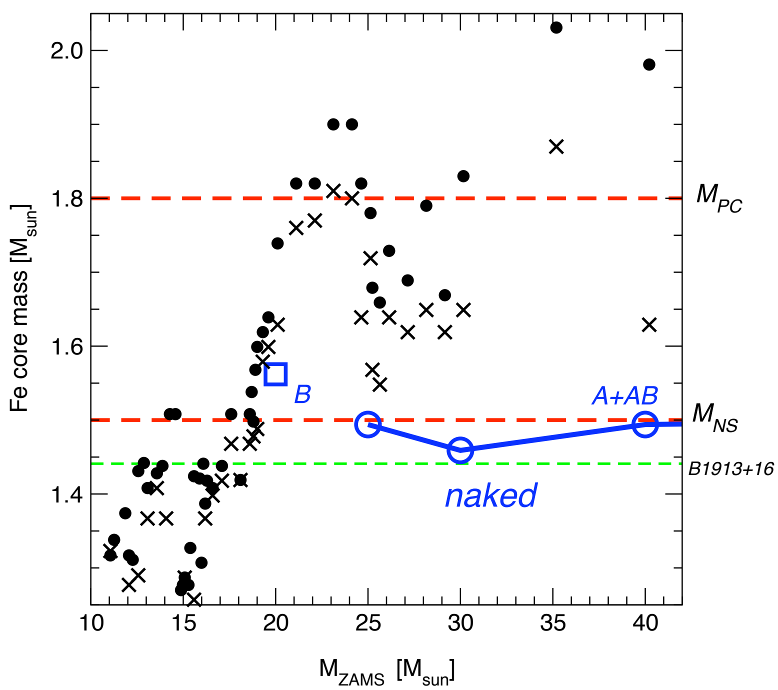

Figure 4.2: Comparison of the iron core masses at the time of iron core implosion for a finely spaced grid of stellar masses [64]. The circular black dots were calculated with the Woosley & Weaver [146] code, whereas the crosses employ the vastly improved Langanke & Martinez-Pinedo [78] rates for electron capture and beta decay. If the assembled core mass is greater than $M_{PC} = 1.8M_{\odot}$, where M_{PC} is the proto-compact star mass as defined by Brown & Bethe [21], there is no stability and no bounce; the core collapses into a high mass BH. $M_{NS} = 1.5M_{\odot}$ denotes the maximum mass of NS [21]. The mass of the heaviest known well-measured pulsar, PSR 1913+16, is also indicated with dashed horizontal line [135].

have been mostly used up by the ram pressure which cleared the way for the GRB (For GRB 980425, MacFadyen [89] estimates this to be $\sim 10^{52}$ergs). Clearly the "unwinding" of the lowest possible rotational energy will continue until the GRB falls below observational threshold. The T_{90} of GRB 060218 is 2100 ± 100s [34], extremely long. It is clear that if all of the rotational energy can be accepted, which should certainly be true for a total energy of only $E_{\rm rot} \sim 4 \times 10^{51}$ergs, then the time will be very long because the kinetic energy is mostly used up in clearing the way through the star for the GRB. We understand in this way that the subluminous bursts from low rotational energy have a long lifetime. Kaneko et al. [73] find the energies, kinetic and thermal, to be the smallest in GRB 060218 among the subluminous bursts, GRB 980425, GRB 030329, GRB 031203 and GRB 060218, they investigate. It should be noted that the very low GRB energy comes from a near cancellation of the kinetic energy by the ram pressure work. Note that with such an energy this would be a nearly "dark" explosion like Cyg X$-$1 [29], which was estimated to have ~ 6 bethes of rotational energy, except that the black hole must be $\gtrsim 1.5 M_\odot$ at the time of explosion rather than the $\sim 7 M_\odot$ black hole calculated by Brown et al. [29] for Cyg X$-$1. In the case of Cyg X$-$1, extensive mass transfer from the donor to the black hole and mass loss from the donor complicated the simple interpretation using Kepler's law. None the less, the ~ 6 bethes obtained from Cyg X$-$1 should not be so different from the ~ 4 bethes for GRS 060218 because the donor in the latter must be a factor of ~ 3 lower in mass than the estimated $\sim 30 M_\odot$ in Cyg X$-$1, in order to be less massive than the black hole progenitor. This would increase the rotational energy by a factor of ~ 3 whereas the black hole in GRS 060218 is $\gtrsim 1.5 M_\odot$, a factor of $\sim 14/3$ less than that of the $7 M_\odot$ black hole in Cyg X$-$1 at the time of common envelope evolution. Scaling 6 bethes by $3/(14/3)$ gives ~ 4 bethes. **?**] found the mass loss in the explosion of Cyg X$-$1 to be very small, the space velocity of Cyg X$-$1 relative to the Cyg OB3 association to be typical of the velocities of stars in expanding O-star associations and gave other arguments supporting a "dark" explosion. The truly remarkable property of GRB 060218 is the long T_{90} time of the GRB of 2100 ± 100 seconds. This is, however, the typical time of the Blandford-Znajek engine (eq.(8) of Lee et al. [80])

$$\tau_{\rm BZ} = \frac{M_{\rm BH} c^2}{B^2 R^2 c^2} \sim 2.7 \times 10^3 \left(\frac{10^{15}{\rm G}}{B}\right)^2 \left(\frac{M_\odot}{M}\right) {\rm sec.} \tag{4.1}$$

The low magnetic field, appropriate for low-mass black holes and the low $M_{\rm BH}$ tend to make $\tau_{\rm BZ}$ longer, such that the measured $T_{90} = 2000 \pm 100$ sec. is just the right time for the GRB in 060218. That of GRB 060218 is certainly no more than "dusky".

We feel confident that the subluminous GRBs can be described by our binary scenario. As noted, we do not claim that the cosmological ones can be; they may require the original Woosley model. However, our model is not really different, in the sense that we use the Woosley collapsar model after the donor has spun up the black hole progenitor to the necessary amount of rotation, and we feel that understanding the dynamics of LMC X−3 will help us make a connection between the subluminous and cosmological GRBs.

4.4 Why Case C Mass Transfer?

It is clear that Case C, mass transfer following He burning, is useful in the binary evolution of black holes, especially in a Galaxy with solar metallicity. In this case the large winds off the stellar surfaces, referred to by Fruchter et al. [51], do not blow away the He because it is covered by hydrogen -just like in single stars, and, therefore, protected from the strong He-star winds. Were this not so, the winds in our solar-metallicity Galaxy, *would* blow off the helium, lowering the stellar masses sufficiently that they would go into neutron stars, the fate described by Fruchter et al. [51]. The idea that life in the Milky Way is protected by the metals keeping the GRBs away [132] obtained from Fruchter et al. [51] is correct in that they do keep the high-luminosity long GRBs away. The idea that the black hole progenitor does not tidally lock with the donor, until the explosion, found in the literature of the soft X-ray transient sources is incorrect [137].

Aside from the help in tidal locking, the Case C mass transfer accomplishes two other mechanisms. Firstly, the remaining helium is originally on the outer part of the He star, later supported by the angular momentum that cannot be transferred into the black hole, so that it can leave in the Blaauw-Boersma explosion. As can be seen from Lee et al. [81], Case C mass transfer also makes the initial binary separation a_i insensitive to donor mass, making our use of Kepler's law easy.

This latter argument results from the nature of the He supergiant at the time of mass transfer [81]. In order for the mass transfer to be delayed sufficiently for Case C mass transfer to take place, winds must be sufficiently low that the metals obtained at the end of the burning do not blow away, but extend in space beyond the helium. Otherwise Case B mass transfer would take place. For a ZAMS $20M_\odot$ giant, this means that the He supergiant must expand to $\sim 1000R_\odot$. Adding the $\sim (200-300)R_\odot$ Roche lobe of the donor gives mass transfer beginning at $(1200-1300)R_\odot$.

In either case, the mass transfer is well localized in radius because the donors will have sizes which are small compared with the radii of the supergiants, so that their Roche lobes are typically about 0.2 of that radius. In Lee et al. [81] we found all supergiants to have $\sim 30M_\odot$. This was because in the Galactic evolution, binaries with lesser mass black holes would lose more than half their mass in the explosion and not be stable [99], but with the lower rotational energies of the low metallicity binaries, the lower-mass giants of $\sim 20M_\odot$ are likely to be more copious. These would allow an $a_i \sim (1200-1300)R_\odot$ for Case C mass transfer, without strong variation, so it is reasonable to forget the binary dependence of a_i. If we can do this we can easily use eq. (4) of Moreno Méndez et al. [99]

$$\frac{\text{days}}{P_b} = \left(\frac{4.2R_\odot/a_i}{M_d/M_\odot}\right)^{3/2} \left(\frac{M_d+M_{He}}{M_\odot}\right)^{1/2} \left(\frac{M_{giant}}{M_\odot}\right)^{0.83}, \tag{4.2}$$

where it is seen that the approximation

$$\frac{\text{days}}{P_b} \propto \left(\frac{4.2R_\odot/a_i}{M_d/M_\odot}\right)^{3/2} \left(\frac{M_d+M_{He}}{M_\odot}\right)^{1/2} \tag{4.3}$$

gives the main dependence as

$$\frac{P_b}{days} \propto \left(\frac{M_d}{M_\odot}\right)^{3/2} \tag{4.4}$$

for $M_d \ll M_{He}$ and

$$\frac{P_b}{days} \propto \frac{M_d}{M_\odot} \tag{4.5}$$

for higher M_d. These can be normalized to Galactic binaries, say Nova Sco or Il Lupi, where the theoretical and observational Kerr parameters agree. We reproduce Fig.1 of Brown et al. [29] as our Fig. 2.1 to show how it is possible to obtain the Kerr parameter from P_b.

4.5 LMC X−3

The LMC has $\sim 1/3$ solar metallicity, so it should go a long way towards having low metallicity stars. The donor of LMC X−1 is more massive than the black hole and it seems to be somewhere between Cyg X−1 and M33 X−7 in mass and their relative sizes. In the case of LMC X-3 Davis et al. [42] find $a_\star \simeq 0.26$ now; using the Cowley et al. [41] $7M_\odot$ for the black hole would imply using $4M_\odot$ for the donor. The present period is 1.7days, with present rotational energy of 54 bethes, obtained from the period from Fig. 2.1.

Note that the much lower energy than those of most Galactic binaries result from the $2-4$ times larger donor mass here. It would be natural for the natal rotational energy to be about double this, with energy ~ 108 bethes. The hypernova energy of SN1998bw is ~ 30 bethes, and, as noted earlier, the kinetic energy should be about equal to this because of equipartition of kinetic and thermal energies. The GRB 980425 is a "smothered" one [89], but, in general, the visible GRB is only ~ 1 bethe, so in most cases the energy used up by the ram pressure in order to clear out material from the path of the jet must be cancelled in first order of magnitude by the kinetic energy. We believe this near cancellation between kinetic energy and energy needed for the ram pressure work to produce the various observed low-energy GRBs (see our later discussion).

We believe that adding the ~ 30 bethes kinetic energy to the 30 bethes measured hypernova explosion energy, assuming equipartition between the two energies, gives us an energy of ~ 60 bethes, possibly that of cosmological GRBs. At least, this is the largest energy that we have found recently in the literature. We argue that GRB 980425 is subluminous because of the near cancellation between work of the ram pressure and the kinetic energy. (But it could be subluminous because of the high, nearly solar metallicity [127] in the background galaxy.)

The present rotational energy of LMC X−3 with Davis et al. [42] $a_\star = 0.26$ can be obtained from the Blandford-Znajek formula

$$E_{BZ} = 1.8 \times 10^{54} \epsilon_\Omega f(a_\star) \frac{M_{BH}}{M_\odot} \text{ergs} \tag{4.6}$$

where we take the efficiency $\epsilon_\Omega = \Omega_F/\Omega_H$ (where Ω_F is the rate of rotation of the field and Ω_H is that of the black hole) to be 1/2 (for optimum impedance matching). Here

$$f(a_\star) = 1 - \sqrt{\frac{1}{2}(1 + \sqrt{1 - a_\star^2})}. \tag{4.7}$$

For the present $a_\star = 0.26$, this gives $E_{BZ} = 54$ bethes. It would be natural if the natal energy would be about double this, 108 bethes. In any case, this can be checked, in that the mass loss in the Blaauw-Boersma explosion (see the Appendix of Brown et al. [27]) has (after recircularization) the present period P_{now}

$$P_{now} = \left(1 + \frac{\Delta M}{M_{BH} + m_D}\right)^2 P_{natal} \tag{4.8}$$

where m_D is the donor mass and we have used the relation between $a_\star$ and P of fig. 2.1 leads to a natal period of little over 1 day, compared with $P_{now} = 1.7$ days. From eq. 4.8 we find

$$\Delta M = 3.25 M_\odot, \tag{4.9}$$

which gives a systems velocity of 43.5 km/s. This velocity is in the process of being measured by the Smithsonian-Harvard coalition [94]. If even one component of the system velocity is measured, it will give a lower limit on the explosion energy. From our estimates this should be substantially larger than the $\sim$ 1 bethe energy usually discussed for cosmological GRBs. We have arrived at our estimate from the measured Kerr parameter and what is essentially dimensional scaling; so we believe our prediction will be fulfilled.

We note that LMC X−1 has a donor mass more massive than its black hole [108], $M_{BH}/M_{donor} \sim (0.3 - 0.7)$ with black hole mass $\sim (7 - 13) M_\odot$. Thus LMC X−1 is similar to Cyg X−1, and would have been expected to undergo a “dark” explosion.

4.6 Conclusions

We suggest that LMC X-3 may be similar to relics of cosmological GRBs, to the extent that some or most of these latter result from binaries. The number of the latter is certainly sufficient to produce the GRBs and the binary nature takes care of the necessary angular momentum, which can be achieved by choosing the donor mass.

Fruchter et al. [51] have made a case that high-luminosity, long GRBs came from irregular low-metal galaxies. We suggest LMC X-3 as the closest nearby binary from a region of 1/3 solar metallicity in an irregular galaxy. The lower metallicity environment has more massive stars, and the donor in LMC X-3 is probably at least twice as massive as the donors of the relics of the Galactic GRBs, which should slow the binary down to a rotational energy that can be accepted. We believe that LMC X-3 brings us in the metallicity towards the irregular low-metallicity binaries considered by Fruchter et al. [51].

Working in this region of energies we feel that we can make predictions, because our calculation of Kerr parameters makes it possible to make quantitative calculations.

Our predictions are, however, at best, order of magnitude, because the necessary properties of the binary LMC X-3 have not been accurately measured. As they are measured we will probably have to readjust our numbers. In particular, it would be most valuable to obtain a reliable value for the explosion energy. Whereas the hypernova energy is probably measured with reasonable accuracy, the kinetic energy required by the work performed by the ram pressure to clear the way for the GRB is mostly hidden, and the net GRB energy is a small difference between large kinetic energy and a large amount of work done by the ram pressure. We believe that there will be great variations in GRBs depending upon these small differences which must be sensitive to stellar properties.

Our present estimate of ~ 60 bethes for the total cosmological GRB energy is the highest recent estimate that we have seen. For this estimate we invoke equipartition of kinetic and thermal energies.

Finally, we wish to point out that our ability to calculate Kerr parameters and the confirmation of these by the Smithsonian-Harvard measurements makes it possible to calculate in a reliable way the amount of angular momentum carried by the black-hole binary. This angular momentum is conserved, but the questions to be answered concern how it is to be distributed: how much remains in the binary, how much energy goes into the explosion and what is the fraction of the latter which goes into kinetic and into heat energy? Progress in answering these questions is necessary to make a quantitative study out of the GRBs, and to put order into their classification.

Chapter 5

The Role of Metallicity in Long GRBs.

5.1 Abstract

The role of metallicity in determining long gamma-ray bursts has been expanded greatly recently, beginning with the Fruchter et al. [51] Nature paper. Long γ-ray bursts have been shown to come predominantly from irregular dwarf galaxies with low metallicity. In addition to the role of metallicity, Modjaz et al. [98] and others have shown that empirically the GRB in the usual GRB plus Hypernova combination tends to be absent, or "smothered", as the metallicity increases. Essentially no GRBs observed in galaxies with the metallicity of our Galaxy. We have a dynamical model in terms of the Blandford-Znajek mechanism we developed to explain the soft X-ray transient sources in our Galaxy. We changed the metallicity in this model, so as to fit the subluminous GRBs [29]. The change to lower metallicity increases the mass of the donor in the transient sources, and this slows down the GRB so that it can accept the rotational energy of the binary.

We have, in fact, a Goldilocks scenario where the rotational energy of most of the binaries in our Galaxy is several hundred bethes (1 bethe = 10^{51} ergs) so that the impact of such high energy temporarily destroys the accretion disk and dismantles the central engine of the GRB. The mass of the donor must be "just right", a few $M_\odot$ in order for the central engine to accept most of the rotational energy of the binary. If the donor is too massive, the rotational

energy of the binary will not be enough to power a GRB.

We show that the GRB is much more fragile than the hypernova, because of the work that the ram pressure must do in order to clear the way through the helium star for the jet that makes up the GRB. This net ram pressure work is nearly equal to the kinetic energy, so it is easy for the GRB to be "smothered". On the other hand the rotational energy of the black hole is fed into the hypernova explosion over a viscous timescale of days and an order-of-magnitude more energy is transferred than the net energy in the GRB.

5.2 Introduction

Papers on the role of metallicity have received a great deal of observational support, the long (cosmological) GRBs coming chiefly from star-forming, irregular dwarf galaxies [51, 79]. As metallicity is increased, up towards solar, the GRBs disappear leaving only broad SN I_c as relics of the explosion [75, 79, 98].

In a number of papers [25, 29, 81] we have shown that the Blandford-Znajek mechanism can provide a quantitative description of the GRB (when visible), why the GRBs are subluminous, or even absent, in our high metallicity Galaxy, and makes it clear that the role of metallicity is dynamical in the sense that donor masses tend to increase as the metallicity decreases. Brown et al. [29] showed that as the donor masses increase from Galactic (low-mass companion) they reach a value like that of LMC X−3 such that the BH spin energy is equivalent to the GRB and hypernova explosion energy, and this determines the region of donor masses that the cosmological GRBs come from. As the donor mass continues to grow, the spin of black hole progenitor will decrease, the angular momentum energy decreasing until the explosions are low-energy "dark" ones, like Cyg X−1 [29] or in M33 X−7 [109].

Binaries have a few great advantages in powering GRB and Hypernova explosions, as brought out in Brown et al. [25]:

(i) There is no problem in having sufficient angular momentum. The donor spins up the black hole progenitor, with which it is tidally locked, to whatever angular momentum is required, the amount determined by the mass of the donor. The donor then contracts out, remaining only as a witness to the explosion. The α-particle nuclei implanted in the donor help to reconstruct the nature of the explosion.

(ii) Provided Case C mass transfer (mass transfer taking place only after helium core burning is finished) takes place, no hydrogen envelope is left before the explosion. This clearly explains why there are no hydrogen lines in the GRB associated supernovae.

Of course, the He burning transforms the helium to metals, ^{12}C and ^{16}O, in the center of the He star. These metals fall first into the pre-forming black hole. The outer part of the He star which is centrifugally supported, and contains mainly helium, is ejected from the star without interaction in the Blaauw-Boersma explosion, in which the black hole is born. No helium lines appear in the supernovae spectra, all GRB-associated supernovae being Type I_c.

(iii) A further advantage of binaries is that the population predicted by Brown et al. [25] agrees with the number of observed explosions [29].

In this paper we wish to focus on the hypernova explosion, showing that it is much less sensitive to the central engine than the GRB. The latter is powered through the accretion disk onto the black hole. The presence of the accretion disk is important for the Blandford-Znajek process because it is the supporting system of the strong magnetic field on the black hole, which would disperse without the presence from the fields anchored in the accretion disk. In most of the Galactic explosions, the rotational energy thrown onto the accretion disk basically in a time of several Kerr times

$$t_{Kerr} = \frac{2\pi R_{BH}}{v} = \frac{2\pi R_{BH}}{c a_\star} \tag{5.1}$$

where $a_\star$ is the Kerr parameter, is $> 4 \times 10^{53}$ergs, $\sim 20\%$ of $M_\odot c^2$. This is a massive amount of energy and can immediately destroy the disk. The fact that the transfer of rotational energy to the disk is self limiting was realized in 2000 [25], where the estimate of Kerr parameter $a_\star = 0.8$ was used for Nova Sco (the Kerr parameter measured later by Shafee et al. [124]), "The developing supernova explosion disrupts the accretion disk; this removes the magnetic fields anchored in the disk, and self limits the energy the Blandford-Znajek mechanism can deliver". Thus, the reason GRBs are absent in the high metallicity region results from *too much* energy being supplied so that the central engine is dismantled. See Brown et al. [29] for more details.

The magnetic fields will, however, resume coupling to the accretion disk, as the latter is spun up by the field lines from the black hole embedded in it (and threading the disk of the black hole at the other end). Estimates of this will be given in section 5.4 of this chapter. The energy is supplied to the hypernova over a viscous time scale, $\sim$ days, orders of magnitude longer than the dynamical time scale. Since the energy is supplied by the angular momentum, the explosion is cylindrical in nature. A rather complete description of it is given in chapter 17 of Bethe et al. [10].

Whereas the deposition of heat energy into the hypernova explosion is straightforward and observable, the kinetic energy is more problematic. In the description of the usual supernova explosions, the dependence of the explosion can be factored into a function of time times a function of energy. The constant involved is usually set so as to ensure equipartition of energy. In GRB 980425/SN 1998bw the (kinetic) GRB energy is a fraction of a bethe, whereas the hypernova energy is ~ 30 bethes [71]. The point is that there is a large (invisible[1]) ram pressure work, on the kinetic energy side, clearing the way through the He star so as to make the GRB possible. In many papers, the inclusion of this energy is taken into account by giving a low efficiency to the GRB. We cover this matter in chapter 5.5.

We can move from the Galactic binaries into the LMC with reduced metallicity, $\sim 1/3$ Galactic, and find LMC X$-$3 in a star forming irregular dwarf galaxy, a nearby local example of the Fruchter et al. [51] environment for long and luminous GRBs. Reconstructing the explosion of which LMC X$-$3 is a relic, we estimate that it had ~ 4 bethes energy. This value is a small difference between an estimated ~ 30 bethes kinetic energy and the energy expenditure of ~ 26 bethes in the ram pressure work necessary to clear the way for the jet in the helium star. Therefore it may not be precise, but it does have nearly the $\sim 10\%$ efficiency usually found to produce γ-rays.

In chapter 5.7 we show that the GRB and Hypernova explosions can be modelled in terms of Galactic binaries, so that much of what is necessary was carried out by Lee et al. [81]. They can be converted into low-metallicity

[1]So the total kinetic energy may be roughly equal to the potential energy.

binaries by changing the masses of the donors. In particular, giants up to ZAMS mass $\sim 30M_\odot$ are needed, corresponding to He stars up to $\sim 11M_\odot$. These are the size arrived at by Woosley and collaborators in his papers on the Collapsar model.

As developed by Brown et al. [29], the energy of the binary depends chiefly on the mass of the donor. The smaller the donor mass, the higher the spin energy of black hole progenitor.

One can therefore, slow the high-metallicity, Galactic binaries down by introducing more massive donors, which happens naturally with transition to low-metallicity galaxies. Once slowed down, the accretion disk is no longer destroyed and more of the rotational energy can be accepted, perhaps $\sim 50\%$. The donor should not, however, be made much more massive, as is the case of the $\sim 80M_\odot$ donor in M33 X$-$7 because then the explosion energy is too low and the explosion will be a dark one.

5.3 The Role of Metallicity

Larsson et al. [79] provide a strong constraint for gamma-ray burst progenitor masses. They show that long-duration gamma-ray bursts (LGRBs) are much more concentrated on their host galaxy light than supernova explosions. From this they say "GRBs are likely to arise from stars with initial masses $> 20M_\odot$. This difference can naturally be explained by the requirement that stars which create a LGRB must also create a black hole". Of course, the Blandford-Znajek mechanism which we use as central engine to power the GRBs and Hypernova explosions gets the energy from the rotating black hole, as we shall review in detail. Brown et al. [29] have been able to reconstruct, with the help of the Smithsonian-Harvard-MIT-measured Kerr parameters [124] the Galactic GRBs and Hypernovae. In Brown et al. [29] we showed the Galactic population to be just right for the population of subluminous sources in neighboring galaxies.

Fig. 5 of Modjaz et al. [98] divides the high metallicity galaxies (oxygen abundances $12 + \log(\mathrm{O/H}) > 8.5$ broad SN I_c; the GRB being absent) from the low metallicity ones (oxygen abundances$12 + \log(\mathrm{O/H}) < 8.5$) which accompany both broad SN I_c and GRBs. The Modjaz et al. [98] results suggest that the GRBs in high-metallicity environments never materialize. The fact that GRB 980425, just below solar metallicity[2] [127], has a "smothered explosion" [89] and such a low luminosity that it probably is seen only because it is so close.[3] The problem of high metallicity shutting off the GRBs is more complicated than usually discussed. In Brown et al. [25] it was already realized that Nova Sco had a natal rotational energy of ~ 400 bethes$\sim 0.2 M_{\odot}c^2$. With the Shafee et al. [124] measured Kerr parameter of 0.8 (calculated by Lee et al. [81]) this rotational energy is tremendous and that with this much energy "the developing explosion disrupts the accretion disk; this removes the magnetic fields anchored in the disk, and self limits the energy the Blandford-Znajek mechanism can deliver" [25]. In Brown et al. [29] this was described in more detail, in terms of a Rayleigh-Taylor instability. The Shafee et al. [124] measurements of the present Kerr parameter give, within experimental error, the same Kerr parameter as predicted for the natal angular momentum, so it is clear that not much of the rotational energy was used up in the explosion.

Brown et al. [29] showed that the tremendous rotational energies of most of the Galactic black-hole-binary explosions resulted from low, $(1-2)M_{\odot}$ donors. They pointed out that the rotational energy could be decreased to that of cosmological GRBs, by increasing the donor mass and, therefore, reducing the Kerr paramater and the rotational energy down to where it could be accepted. Brown et al. [30] suggested that locally we have LMC X$-$3 in an irregular, low-metallicity star-forming dwarf galaxy, the conditions established by Fruchter et al. [51] for a long (and luminous) GRB, the appropriate donor mass being $\sim 4M_{\odot}$, cutting down the rotational energy supplied in Nova Sco by an order of magnitude. Moreno Méndez et al. [99] showed that the Galactic GRBs with higher donor masses Cyg X$-$1, V4641 Sgr and GRS 1915+105 went through subluminous or dark explosions. We must admit that at first sight two of the Galactic binaries (out of the 15 considered) did have donor masses of $\sim 6M_{\odot}$, V4641 Sgr (XTE J1819-254) and GRS 1915+105. In Moreno Méndez et al.

[2] In fig. 5 of Modjaz et al. [98] it is the subluminous sources with lowest metallicity in which a GRB was seen; it is clearly on the dividing line between "smothered" and "nonsmothered" GRBs.

[3] The Modjaz et al. [98] objects are all in neighboring galaxies and subluminous. The poster child GRB 980425/SN 1998bw, which has nearly Galactic metallicity and a subluminous GRB in addition to an energetic hypernova with ~ 30bethes [71], is the closest to our Galaxy and might well not be seen if it were farther away.

[99] we worked through the binary evolution in detail and showed that these binaries would have had subluminous or dark explosions. In other words, we had to enhance our Galaxy by the LMC in order to find (in LMC X−3) a binary which would have had a long and high-luminosity gamma-ray burst.

So far we have shown that the situation with respect to GRBs is a sensitive one. They can easily be partially or totally smothered and we shall return later to a more quantitative, although model-dependent, discussion of this. A good example being GRB 980425/SN 1998bw the GRB has a very low luminosity, but the hypernova energy of ~ 30 bethes [71] is about the most energetic one measured.

Whereas this has been understood in the Blandford-Znajek mechanism - how the GRB can be smothered but the hypernova explosion can be strong- and has been described in Bethe et al. [10], it is not generally known in the GRB community, so we develop in detail the hypernova explosion in the next section.

5.4 Powering of the Hypernova Explosion by the Rotating Black Hole

Exactly how the black hole powers the heating of the accretion disk, formed by what is left from the outer part of the helium star (the inner part having fallen into the black hole) may not be known, but the general mechanism must be that of Li [83, 84]. The black hole is formed rotating very rapidly.

In Fig. 5.1 we show the schematic view of Blandford-Znajeck mechanism. It can be considered that a virtual wire is run down a magnetic field line to the black hole, around from the north pole of the black hole to the equator and from the equator into part of the accretion disk and then joined into a magnetic field line to complete the loop. The surface resistance of the black hole is 377 ohms, the impedance of a wave guide going into a vacuum. The resistance of the accretion disk may be neglected, because the ionized magnetic fields are inside it and frozen into the ionized matter of the disk. Thus, the disk is spun up.

Now, first, energy is deposited by viscosity into the disk, but later, the viscosity turns the energy into heat. The time scale for this is (p.359 in Bethe

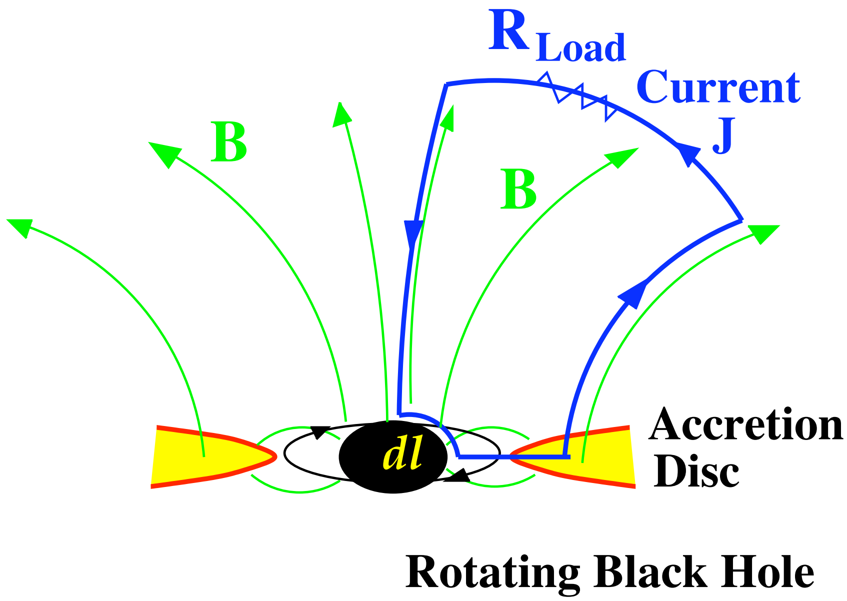

Figure 5.1: Schematic representation of the Blandford-Znajek mechanism. The left hand (red) loop powers the GRB, sending Poynting Vector energy vertically up the rotation axis to the fireball (not shown) in the “loading area”. For simplicity the threading of the magnetic field lines $\vec{B}$ into the disk of the black hole are not shown. In fact, when sufficient rotational energy is fed suddenly, the coupling of the magnetic field is destroyed and the magnetic field lines are pushed out of the event horizon [29]. The coupling of the blue line which goes into the accretion disk is zero to begin with, but then the magnetic field begins to thread the disk of the black hole as the disk begins rotating. The viscous energy deposited into the accretion disk is converted into thermal energy and powers the hypernova explosions.

et al. [10])

$$\tau_{vis} \sim \frac{1}{\alpha}\left(\frac{r}{h}\right)^2 \tau_{dyn} \tag{5.2}$$

where α is the parameterized viscosity, r the radius of the accretion disk and τ_{dyn} the dynamical time, ~ 10 minutes for an $8M_\odot$ He star; $h \sim 0.2r$ [88] for a thin disk, growing τ_{vis} ∼days for $\alpha \sim 1$. Of course, the disk widens out with time, but α drops. The deposition of energy drops once the hypernova takes place.

In the case of SN 1998bw the hypernova energy is $E_{HN} \simeq 30$ bethes. Our point here is that although the coupling of the GRB through its accretion disk to the black hole might be dismantled, the hypernova energy is transferred over days, so the coupling of the accretion disk to the black hole can easily be established. In fact, the delivery of energy to the hypernova is expected to be irregular and somewhat chaotic, not smooth [14]. One can view the black hole as the little ball in a roulette wheel which bounces back and forth, out from the center after it hits it, the roulette wheel being the accretion disk. None the less, there will be plenty of time for attachment of the rapidly rotating black hole to the disk because of the ∼day long τ_{vis}, and it will attach itself. Then angular momentum will be delivered by the black hole to the disk. Indeed, Li [84] shows that the maximum in energy delivery to the disk exceeds the maximum energy that can be delivered to the GRB explosion. However, we believe that a reasonable estimate is that the thermal and kinetic energies are more or less the same as would be suggested by equipartition of energy.

Already in 2000 Brown et al. [25] had realized that the rotational energy in Nova Sco was sufficiently energetic so that the Blandford-Znajek central engine would be dismantled "we shall show that up to 10^{53}erg is available to be delivered into the GRB and into the accretion disk... we suggest that injection of energy into the disk shuts off the central engine by blowing up the disk and thus removing the magnetic field needed for the energy extraction from the black hole. If the magnetic field is high enough the energy will be delivered in a short time, and the quick removal of the disk will leave the black hole still spinning quite rapidly".

In Moreno Méndez et al. [99] we developed the concept of Rayleigh-Taylor instability to keep the magnetic field lines, which otherwise would power the GRB, away from the horizon, were the couplings sufficiently strong. The field lines couple the black hole to the accretion disk. However, the time for transferring energy from the black hole rotation to the accretion disk and heating it up is the thermal time scale, $\tau_{viscous}$ ∼days, not minutes, as for the GRB. Thus there is a long time for the field lines to reattach themselves to the black hole and for it to spin the disk up. In fact. from Faraday's law the

$$EMF = \oint_c \alpha \vec{E} \cdot d\vec{l} = \oint_c -\alpha[\vec{v} \times \vec{B}] \cdot d\vec{l} \tag{5.3}$$

where the lapse function $\alpha = d\tau/dt$ along the closed circuit going through the accretion disk in Fig 5.1, and $\vec{v}$ is the velocity of the accretion disk so that at time zero (chosen as the time when energy is delivered to the GRB) the EMF to the accretion disk is zero. The couplings of the magnetic field between black hole and accretion disk will slowly reestablish themselves and build up with time, delivering energy from the black hole, which slows down.

The viscous energy turns into heat and it is delivered over the viscous time scale until enough energy has been delivered to power the hypernova. For GRB 060218/SN 2006sj, with the least massive possible black hole [30] the hypernova energy is "only" $E_{HN} \sim 2$ bethes [91] but for more massive black holes like GRB 980425/SN 1998bw the energy is ~ 30 bethes. To date these seem to span the spectrum.

5.5 Towards a Schematic Model of the Cosmological GRB

We distinguish the high luminosity cosmological GRBs from the long γ-ray bursts, some of which can be subluminous. About the highest hypernova energy, however, is that of the poster child subluminous burst GRB 980425/SN 1998bw, $E_{\rm HN} \simeq 30$ bethes. The GRB, although often classified as long, is essentially a smothered one [89]. As with ordinary supernovae it is reasonable to assume, at short times, equipartition of energy, so that the kinetic energy is roughly equal to the thermal energy.

In fact, it is amazing that we see GRB 980425 so clearly with E_γ only $\sim 7 \times 10^{47}$ ergs, some 4 orders of magnitude less energy than that associated with typical GRBs [73]. Of course, at 35.6 Mpc it is the closest GRB to Earth.

We believe that GRB 980425 is doubly held down in luminosity:

(i) Because the environment has the highest metallicity, nearly solar, of the subluminous bursts. We do not, of course, know the companion -assuming that GRB 980425 came from a binary- but from our forego-ing arguments it is reasonable to assume that like in our Galaxy, the companion is low in mass and the GRB is turned off almost immedi-ately after it begins because the vast amount of energy presented to it destroys the central engine. For want of a more quantitative argument, we might attribute an approximate two orders of magnitude decrease in luminosity due to this.

(ii) There are indications that GRB 980425 is a "smothered explosion" [89, 90]. This means that, although the kinetic energy supplied may be roughly equal to the heat energy (the Hypernova energy), the jet must use up most of this kinetic energy to remove the He star matter in the way in order to get out. This work that the net ram pressure must furnish is estimated by MacFadyen [89] to be $\sim 10^{52}$ergs.

Since we can make some rough estimates of (ii), we shall assume that the donor in GRB 980425 was several $M_\odot$ in magnitude, so that (i) does not come into play, although GRB 980425 has such a low luminosity that both (i) and (ii) may be needed. Because of the low metallicity in almost all of the observed subluminous sources, (i) would not be applicable to them and we have only (ii) available. Of course, it may be naïve to model simply a situation that is as complicated as jet quenching, but doing so will make us think about relevant effects. The kinetic energy, assumed to be $\sim$ 30bethes for a cosmological GRB, must be nearly completely spent in removing the matter in the way of the jet giving rise to the GRB, freeing the funnel of matter. MacFadyen [89] estimates this energy to be $(0.01 - 0.1)M_\odot c^2$ between 10^{52} and 10^{53} bethes. We clearly could take it to be only slightly less than 30 bethes, so that the explosion is smothered, only $\lesssim 10^{-2}$bethes coming out in the GRB [73].

We make a schematic model about the quenching of the jets by the He star. The helium star radius can be approximated to be

$$R_{He} = 0.2 \left(\frac{M_{He}}{M_\odot}\right)^{0.6} R_\odot. \tag{5.4}$$

The jet must push out all of the matter initially in the open cone along the rotational axis. Since SN 1998bw had the most energetic hypernova, we take the ZAMS mass of the giant to be $30M_\odot$, with He star mass of $11M_\odot$ (essentially the giant and helium star masses chosen by Lee et al. [81] in evolving Nova Sco, Il Lupi and V4641 Sgr); this was the largest giant mass considered. Nomoto et al. [105] chose a $16M_\odot$ He star, which corresponds to a ZAMS mass $\sim 40M_\odot$ for the hypernova 1998bw. Remarkable in the spectra was the large amount of ^{32}S ejected, much greater than in a supernova explosion (the much greater hypernova explosion digs more deeply into the more bound elements). We believe, however, that a $40M_\odot$ progenitor is too massive and that with the metallicity of nearly the Galactic value, such a massive star would "blow away". We believe that the hypernova could be evolved from an $\sim 30M_\odot$ ZAMS star, such as Lee et al. [81] used.

Now, the GRB energy in GRB 980425 is tiny. Kaneko et al. [73] put it at $E_{GRB} = 0.0015$bethes; GRB 980425 may have been seen only because it is the closest (35.6 Mpc) GRB to our Galaxy. A ZAMS $30M_\odot$ giant would have an $11M_\odot$ helium star. So we assume for purposes of comparison with LMC X$-$3 that GRB 980425 was completely "smothered", neglecting its small luminosity. On the other hand the He star in LMC X$-$3 lost $\sim 3.25M_\odot$ from an initial $10.25M_\odot$ helium star in the Blaauw-Boersma explosion in which the black hole was born [30]. Now the ram pressure work must first pay the binding energy of the matter in the space that must be cleared from the jet and then accelerate the matter to some fraction of the velocity of light. The radius of a He star of mass M goes as $M^{0.6}$, so that the binding energy goes as $M^2/R \propto M^{1.4}$. Thus, the difference between kinetic energy and ram pressure work is

$$\mathrm{KE} - \mathrm{RPW} \simeq \left[1 - \left(\frac{10.25}{11}\right)^{0.6}\right] 30\text{bethes} \simeq 2.8\text{bethes}, \tag{5.5}$$

which is just the usual energy for a cosmological GRB.

In other words, we get from a nearly "smothered" GRB to the energy of a cosmological GRB by removing a near cancellation due to the greater mass loss in the latter. This certainly has the aspect of being contrived, but it illustrates that very special circumstances may be required to form a cosmological GRB and basically supports our "Goldilocks" scenario that the energy supplied must be neither too great, so as to dismantle the central engine, nor too little, but just right. It is likely that in a better, more detailed calculation than our rough schematic one, there are unforeseen feedbacks which put together the high-luminosity GRBs.

The high explosion energy which results in $\sim 3.25M_\odot$ being lost from the He star in LMC X$-$3 results from the $\sim 4M_\odot$ donor being close to optimum for rotational energy that can be accepted. So these effects bootstrap each other resulting in an extremely strong GRB, just as the effects suppressing the GRB tend to combine to make it less luminous in GRB 980425.

Our goal here is to bring out the "invisible" energies so that one can reach a semiquantitative picture. The Universe has a massive number of black hole binaries from which to get luminous GRBs, so there are large selection effects in those we do observe.

However, we believe that the near cancellation between kinetic energy from the angular momentum energy and the net ram pressure (NRP) work will continue as the mass of the He star is decreased. The Blandford-Znajek black-hole energy goes linearly with black-hole mass, although much of it is determined by the Kerr parameter. Nevertheless, as seen in GRB 060218 the energy for the lowest mass binary is also the smallest GRB energy. Furthermore, the NRP work decreases as the helium star mass decreases. Thus, we can postulate that order-of-magnitude the NRP work cancels the kinetic energy, so that only bethes are left in the difference, not the tens of bethes in the original kinetic energy and NRP work.

Obviously we cannot calculate the ~ 1 bethe often discussed as a possible standard candle for cosmological GRBs. The calculation of MacFadyen & Woosley [88] of the GRB jet was quantitative and included the possible energy from magnetohydrodynamic effects (Blandford-Znajek mechanism) as needed. "However, our jet is produced over a longer time than 1 second, so its duration will not be governed solely by light propagation effects, but by the time the engine operates after the polar regions have cleared, about $(10-20)$s. Moreover, ours is an unsteady jet modulated both by accretion disk instabilities and the dynamics of the stellar nozzle through which the jet flows. Thus the GRB will have time structure given not only by the circumstellar interaction, but also by any observable residuals of the unstable flow" [88].

We have tried to give a schematic model of this complicated situation, showing the near cancellation between kinetic energy and net ram pressure work. At least we have shown that for situations encountered in the cosmological GRBs there must be this cancellation. Many of the GRBs will be suppressed; some like GRB 980425 will be nearly smothered and some obviously get out with $\sim$ 1bethe of energy.

We should remark that in GRB 980425 and GRB 060218 we were able to see very subluminous effects. It is no wonder that these were the two closest GRBs to our Galaxy.

In the last section we showed that the situation with respect to the hypernova explosion is much simpler, and the best bet on overall necessary energy input is $\sim 2E_{HN}$.

5.6 The "Goldilocks" Scenario.

None of the observational papers in the literature, except those dealing with our Galaxy, discuss GRBs as arising from binaries. In Brown et al. [29] we showed that the Galactic soft-X-ray transient sources could be extended to reproduce the subluminous GRBs in the neighboring Galaxies. As occurring in binaries, and to show that although the binaries like Nova Sco and Il Lupi had too much rotational energy, so that the central engine was quickly dismantled, and M33 X$-$7 had too low a rotational energy, so that it probably underwent a "dark explosion", a binary like LMC X$-$3 probably had just the right rotational energy so as to mimic the cosmological long GRBs. LMC is as good example of the star-forming, irregular-dwarf galaxies with low metallicity favored by the Fruchter et al. [51] arguments for long and highly luminous GRBs.

We now proceed to map Galactic soft-X-ray transients onto intensively studied GRBs, including "Goldilocks" rotational energies that are too high, just right and too low for γ-ray bursts.

In our Galaxy, with high metallicity, 12 of the 15 black hole binaries had low donor masses $M_D \sim (1-2)M_\odot$ and 3 had higher masses $M_D \gtrsim 10M_\odot$ [99]. In the first case the GRB explosions were subluminous because there was so much energy that the black hole accretion disk was temporarily destroyed

Reconstruction of Galactic Soft-X-Ray Transients, Subluminous and Cosmological GRBs.				
	Giant Mass	Donor Mass	Example	Fate
Too much Energy	Galactic Transient Source $\sim 30 M_\odot$	$(1-2)M_\odot$ ($E_{GRB} \simeq$ 400 bethe)*	Nova Sco Il Lupi	Dismantle Accretion Disk
		($E_{GRB} \simeq$ 0.0015 bethes)	GRB 980425	
Just Right	Galactic Transient Source $\sim 22 M_\odot$	$\sim 4 M_\odot$ ($E_{GRB} \simeq$ 4 bethes)	LMC X$-$3	Cosmological type. Long luminous GRBs.
Too little Energy	$(18-20)M_\odot$	$\sim 10 M_\odot$ ($E_{GRB} \simeq$ 0.000021 bethes)	GRB 060218	Long, but subluminous GRBs.
	$85 M_\odot$	$\sim 80 M_\odot$ ($E_{GRB} \simeq$ 10 bethes)	M33 X$-$7	Probably "Dark"** Explosion.

Table 5.1: The GRB energies of GRB 980425 and GRB 060218 are the GRB E_{min} from Table 8 of Kaneko et al. [73], showing essentially that the GRBs were subluminous. The GRB energy of LMC X$-$3 is calculated, by subtracting the ram pressure work needed to clear the way for the GRB to blast out of the He star; alternatively, the GRB energy multiplied by the γ-ray efficiency. Comparison of the sets of energies is unwarranted; all that is meant by showing them is that the subluminous energies are small fractions of a bethe; energies of long, high-luminosity GRBs are bethes. (* Only a tiny fraction of this energy can be accepted. Almost all of this energy remains in the binary, as found by Shafee et al. [124]. ** Although the energy is that of cosmological GRBs, the Kerr parameter is only $a_\star \sim 0.12$ because of the high mass donor. See Mirabel & Rodrigues [97] who characterize the explosion of Cyg X$-$1 as "dark".)

almost immediately. In the second case, the explosions were essentially "dark".

The relatively high GRB energy in LMC X−3 results because the net ram pressure work required to clear the way for the jet is less than in GRB 980425 because the mass loss in the Blaauw-Boersma explosion is $\sim 3.25 M_\odot$ (as discussed in the previous section and from [30]). The He star progenitor of the black hole is substantially less massive than in the "smothered" case (where the progenitor was $\sim 11 M_\odot$, [31], or Chapter 5). Because of the high mass donor in the latter case there is very little mass loss in the explosion. Thus, we find a kind of self-consistency. A lot of mass loss in the explosion will mean a lower-mass He star, and therefore less ram pressure work to clear matter from the path of the jet, and less ram pressure work to be subtracted from the kinetic energy leaves more kinetic energy to power the GRB.

Once again, with the Blaauw-Boersma explosion, Goldilocks enters. In the case of the Galactic Nova Scorpi, the explosion is so energetic that only slightly more than half of the original mass survives, leading to the very large mass loss from an initial $11 M_\odot$-He star down to an $\sim 5.7 M_\odot$ black hole. This deposits the $\sim$ 400-bethes rotational energy of the binary which destroys the accretion disk of the black hole.

The more modest $3.25 M_\odot$ mass loss in LMC X−3 leaves a $7 M_\odot$ black hole. The estimated system velocity for LMC X−3 is 43 km s^{-1}, about 1/3 of the $\sim$ 114 km s^{-1} for Nova Sco.

Whereas Brown et al. [29] showed that the rotational energy furnished to the black hole in Nova Sco was so great that it dismantled the disk, the lower rotational energy in the case of LMC X−3, only $\sim$ 60 bethes (see Chapter 4 or [30]), should be able to be accepted. Most of the $\sim$ 30 bethes kinetic energy is used up in the work of the net ram pressure to clear the way for the GRB. Since $3.25 M_\odot$ of the $\sim 10.25 M_\odot$ helium star has been lost in the explosion, it has less ram pressure work to pay than in GRB 980425 where the explosion is very weak, so the GRB must penetrate the full $11 M_\odot$ He star.

For GRB 060218, as noted in Brown et al. [29], the lack of carbon lines (see Mazzali et al. [91], additional information) determined that the black hole came from the giant threshold for black hole production with ZAMS mass $(18-20) M_\odot$ giving a black hole[4] $M_{BH} \lesssim 2 M_\odot$. In Lee et al. [80] the

[4] Note that this was just the ZAMS mass for the Sanduleak sk-69^{o}202a, the progenitor

Blandford-Znajek GRB lifetime is

$$T_{BZ} = 2.7 \times 10^3 (M_\odot/M_{BH}) B_{15}^2 = 1,700\text{s} \tag{5.6}$$

for $B_{15} = 1$ and $M_{BH} = 1.6 M_\odot$. This is to be compared with the (very long)

$$T_{90} = (2,100 \pm 200)\text{s} \tag{5.7}$$

measured by Campana et al. [34]. The tiny fraction of a bethe in the GRB energy must result from a very small efficiency for γ-rays, essentially a near cancellation between kinetic energy and ram pressure work [31].

5.7 We Can Model GRBs With Galactic Black Hole Biunaries

It may have occurred to the reader, had he read Lee et al. [81], that our remodelling of GRBs and Hypernova explosions followed the modelling of black hole binaries in our Galaxy, beginning with Brown et al. [25] with one great exception. Namely, the donors in the low-metallicity galaxies favorable for long γ-ray bursts, had to be more massive than donors in Galactic black hole binaries, having masses up to $5M_\odot$.[5] The only extragalactic binary in which the masses were measured is M33 X−7 in which the $\sim 70M_\odot$ donor had a mass of $\sim 80M_\odot$ at the time of common envelope evolution and the $15.65M_\odot$ black hole is the most massive measured stellar black hole. Brown et al. [29] showed that with such a massive donor, the explosion would have been a dark one.

In Lee et al. [81] Nova Sco and Il Lupi were evolved from binaries with $\sim 30M_\odot$ giants with $\sim 11M_\odot$ helium stars. The 8 AML (angular momentum loss) binaries could not be evolved quantitatively in Lee et al. [81], but they suggest that all of their black holes could have come from $\sim 11M_\odot$ or less massive He stars. Woosley and collaborators in a number of papers have

in SN 1987a

[5]In our Galaxy, we have the binaries XTE J1819−254, GRS 1915+105 and Cyg X−1 which had to be evolved with a large amount of mass exchange and which were shown to be subluminous or dark. The other 12 binaries had donors less than $\sim 2M_\odot$.

estimated $\sim 10M_{\odot}$ to be the mass of their Wolf-Rayet Collapsar.

In other words, as far as we know, even though with low metallicity there are large numbers of very massive stars, $M \gg 30M_{\odot}$, only stars of ZAMS mass $\sim 30M_{\odot}$ and less, giving He stars of $\sim 11M_{\odot}$ and less, are relevant for the cosmological GRBs. Thus, essentially all of the quantitative background work was done in Lee et al. [81].

This background work can be immediately applied to the other galaxies, with the low-metallicity in most of them simply changing the masses of the donor (the mass of the giant plays a much less important role).

In the 12 Galactic binaries other than the 3 we mentioned before (which gave subluminous explosions) the natal rotational energies are all > 430 bethes ($\sim 1/5M_{\odot}c^2$), which would be nearly unbelievably large, except that for Nova Sco and Il Lupi the Smithsonian-Harvard-MIT collaboration measurements confirmed that the binary still has this much rotational energy. In Moreno Méndez et al. [99] and Brown et al. [29] we showed that with so much energy supplied to the accretion disk the magnetic coupling would temporarily go into Rayleigh-Taylor instability with stochastic properties that would turn the Blandford-Znajek central engine off. Thus, the fact that GRBs are not seen at high metallicity, although the hypernovae are, is explained by the quantitative calculations. In fact the dismantling of the Blandford-Znajek mechanism in Nova Sco was already mentioned by Brown et al. [25] where a Kerr parameter of 0.8 was used.

We showed in Brown et al. [29] and Moreno Méndez et al. [99] by population synthesis (again, using the population of soft X-ray transient sources in our Galaxy given in Brown et al. [25]) that there are enough of our binaries to reproduce all GRBs. This does not exclude GRBs from Woosley's Collapsar model which could easily be included in the population and, other than the fact that we use a donor to spin the black hole up to the necessary angular momentum, our model then proceeds precisely in the way of Woosley's (except that we may have an advantage in producing a I_c explosion because of our Case C mass transfer).

5.8 Conclusions

We saw many years ago, after Woosley had proposed his Collapsar model [145] that using a rotating black hole to power the central engine for a GRB had great advantages. The chief one was the nearly baryon-free region necessary for the explosion, because it would take only a few nucleons to slow it down.

Although the Blandford and Znajek had been published in 1977 [13], there had been many papers criticizing it, modifying it, etc., all collected in Lee et al. [80]. They were put into a form that makes them easy to apply.

The missing links were the Kerr parameters of the black hole, which are necessary in Blandford-Znajek to give quantitative answers. In Lee et al. [81] the Kerr parameters of some Galactic binaries were published, they also showed how the others could be calculated. Without knowing this paper, the Harvard-Smithsonian-MIT collaboration in 2006 [124] published the Kerr parameters of Nova Sco and Il Lupi, which confirmed the Lee et al. [81] calculations within observational errors ($\sim 10\%$).

After the publication of Lee et al. [81] a number of papers were published saying that the black hole in the soft X-ray transient sources would not be spun up as if tidally locked with the donor. These papers did not take into account the strong magnetic fields (progenitors of the $\sim 10^{15}$ gauss B-fields attached to the black hole) which, when Case C mass transfer is used, lock the preforming black-hole progenitor core with the outer layers of the He star, the latter being tidally locked with the donor.

In 2006 the subluminous GRBs were highly publicized in 5 papers in Nature [34, 91, 112, 126, 148], and population syntheses were carried through. Moreno Méndez et al. [99] and Brown et al. [29] realized that the population fit that predicted by Brown et al. [25] for soft X-ray transient sources, furthermore that in going over to low metallicity the donor masses of the binaries would grow, slowing them down to where most of them would become subluminous. By choosing a donor mass estimated to be $\sim 5M_\odot$, the energy would be sufficient for a high luminosity cosmological GRB. We have proposed that LMC X$-$3, which is in the environment of an irregular, fairly low-metallicity ($\sim 1/3 - 1/4$ solar), dwarf galaxy, coming close to fulfilling the condition of Fruchter et al. [51] for long, high-luminosity GRBs, is as close as we can come in our neighborhood to testing our theory. We predict [30] that it is a relic of an explosion of cosmological energy. The black hole mass in LMC X$-$3 is probably relatively low, $\sim 7M_\odot$, so the net ram pressure work is not so large as to cancel the kinetic energy, and there is an estimated $\sim 3.25M_\odot$ mass loss in the explosion. Therefore, there should be a large system velocity which should be able to be at least partially checked (in the radial direction). The Harvard-Smithsonian-MIT coalition has this on their program [94].

We have given models for the most energetic "smothered" explosion, GRB 980425/SN 1998bw with thermal energy of $\sim$ 30bethes, and the least energetic explosion GRB 060218/ SN 2006sj, with thermal energy $\sim$ 2bethes and Brown et al. [30] have suggested that LMC X$-$3 is the relic of a GRB of cosmological energy, $\sim$ 2bethes. The latter energy results from the difference of large kinetic energy and large ram pressure work, so it is only order-of-magnitude. None of our estimates have taken into account specific properties, other than estimates of the masses, of the various stars.

A lot of observational material has accumulated since the early work in Brown et al. [25]. This has enabled us to be much more specific in the application of the Blandford-Znajek mechanism, especially since the Kerr parameters can now be calculated. We must stress the usefulness of the observational data from the Smithsonian-Harvard-MIT group, which discovered that in Nova Sco and Il Lupi very little energy had been used up in the explosions and that the binaries were still rotating very rapidly. Within observational uncertainties, observations checked the Lee et al. [81] theory predictions. There are many links in the chain of arguments leading to these predictions, the most important one being the Case C common envelope evolution, which we believe helps substantially in synchronizing the spin of the black hole with the orbit of the donor and in bringing about the supernova Type I_c nature of the explosion.

In this paper various types of explosions have been modelled: (1) Smothered, (2) Subluminous, (3) Cosmological, and stuck our necks out substantially as to semiquantitative predictions. We believe that this may help to organize the study of GRBs and Hypernova explosions in case our predictions are not too far off. We have shown that using the Blandford-Znajek mechanism we can power GRB and Hypernova explosions by the rotational energy of black holes. We show by calculations of the soft X-ray transient sources in our Galaxy that we can reproduce the Kerr parameters measured by the Smithsonian-Harvard-

MIT coalition; in fact, we predicted some of them.

The long and luminous GRBs found cosmologically are not produced in our Galaxy, not because the energy of the transient sources is not high enough, but because it is too high. However the GRB from which LMC X−3 is a relic can be reconstructed to have the environment of a star-forming, irregular dwarf, low-metallicity galaxy where the long, high-luminosity GRBs are formed by Fruchter et al. [51]. These GRBs are numerically a tiny fraction, at most a few percent, of the total number.

Chapter 6

The Cases for Hypercritical Accretion in M33 X−7 and LMC X−1[1].

6.1 Abstract

The Kerr parameters of the black holes in M33 X−7 and LMC X−1 have been measured to be $a_\star = 0.77 \pm 0.05$ [86] and $a_\star = 0.90^{+0.04}_{-0.09}$ [60], respectively. It has been proposed that the spins of the $15.65 M_\odot$ and 10.91 black hole are natal. We show that these are not viable evolutionary paths given the observed binaries orbital periods of 3.45 and 3.91 days, respectively, since the explosions that would produce the black holes with the cited spin parameters and orbital periods would disrupt the binaries. Furthermore, we show that these systems have to be evolved through the hypercritical mass transfer of $\sim 5 M_\odot$ for the former and $\sim 4 M_\odot$ for the latter from the secondary stars to the black holes.

6.2 Introduction

The black-hole spin in M33 X−7 has been very accurately measured to be

$$a_\star = 0.77 \pm 0.05$$

[1]A version of this chapter, only including M33 X−7, was published in Moreno Méndez, E., Brown, G.E., Lee, C.-H. & Park, I.H. 2008, *The Astrophysical Journal*, 689, L9

in dimensionless spin parameter [86]. The authors of this paper show that *"In order to achieve a spin of* $a_\star = 0.77$ *via disk accretion , an initially non-spinning black hole must accrete* $4.9M_\odot$ *from its donor in becoming the* $M_{BH} = 15.65M_\odot$ *that we observe today. However to transfer this much mass even in the case of Eddington limited accretion (*$\dot{M} \sim 4 \times 10^{-8} M_\odot/yr$*) requires* ~ 120 *million years whereas the age of the system is only* $2-3$ *million years. Thus the spin of M33 X−7 must be natal, which is the same conclusion that has been reached for two other stellar black holes"* (Shafee et al. [124] and McClintock et al. [93]). However, Liu et al. [86] noted that their spin derivation is model-dependent and subject to possible systematic errors.

Similarly, the black-hole spin of LMC X−1 has been determined to be

$$a_\star = 0.90^{+0.04}_{-0.09}$$

[60]. And for similar reasons it has been argued that its spin is also natal.

Lee et al. [81] predicted the spin parameters of Nova Sco (X-ray Nova Scorpii 1994) and Il Lupi (4U 1543−47) to be ~ 0.8, with small effects after they were born in the explosion from mass accretion; i.e., they predicted them as natal. However Brown et al. [29] showed that the rotational energy in such binaries scaled inversely with the donor mass at the time of common-envelope evolution preceding the explosion in which the black hole was born. The donor masses of Nova Sco and Il Lupi are $\sim 2M_\odot$ whereas those of M33 X−7 and LMC X−1 were $\sim 80M_\odot$ and $\sim 25M_\odot$ at the time of the explosions, and so Brown et al. [29] and Moreno Méndez et al. [99] suggested that the long orbital periods in M33 X−7 as well as LMC X−1 resulted from dark explosions; the high spin parameters having resulted from mass accretion. This mass accretion would have had to take place at hypercritical rate for both systems as discussed in this letter.

We briefly comment on the hypercritical accretion for $\dot{M}/\dot{M}_{Edd} \sim 10^3$ or greater [23]. The scenario begins with Bondi accretion through the sonic point, which is often greatly larger than accretion at the Eddington limit. Because it had been worked out for a value of $0.31 \times 10^4 \dot{M}_{Edd}$ by Brown & Weingartner [23] and because this value is in the middle of those we shall use in stellar evolution, we use this value, although it could be much greater. The Brown & Weingartner [23] work had been carried out earlier in all detail by Houck &

Chevalier [68] and checked by Chevalier [38]. Let it be noted, however, that there is still considerable controversy in the astrophysical community about whether hypercritical accretion can take place or not. Nevertheless, the general point we address is that if $\dot{M}$ exceeds $\dot{M}_{Edd}$, then some of the accretion energy must be removed by means other than photons. In the case of hypercritical accretion, this excess energy can be carried off by neutrino pairs [23]. In the case of a neutron star, neutrino losses allow the matter flow to join smoothly onto the neutron star surface. In the case of a black hole, the neutrino losses let the matter flow smoothly over the event horizon and disappear into the black hole.

In the work of Podsiadlowski et al. [114], we note two possible stages where hypercritical ($\dot{M}/\dot{M}_{Edd} \gtrsim 10^3$) or supercritical accretion may take place. Podsiadlowski et al. [114] evolve a binary with $M_{BH} = 12M_\odot$, $m_{secondary} = 25M_\odot$ and orbital period of 6.8 days.

a) Hypercritical accretion stage: They show that after the formation of the black hole, (first scenario) the secondary star will overflow its Roche Lobe and will transfer mass at a rate which peaks at $3 \times 10^{-3} M_\odot \mathrm{yr}^{-1}$. The binary detaches after about 10^4 years, once the secondary has been reduced to the mass of the black hole (or earlier if the stellar wind is strong enough).

b) Supercritical accretion stage: Mass transfer continues at a lower rate of $\sim 3 \times 10^{-6} M_\odot \mathrm{yr}^{-1}$ for up to another few 10^6 years through a directed wind.

Nevertheless Podsiadlowski et al. [114] restrict accretion into the black hole to the Eddington limit. With this same assumption they were able to get GRS 1915+105 up to a spin parameter $a_\star = 0.9$. However, McClintock et al. [93] measured its spin parameter to be $a_\star \sim 0.98 - 0.99$. We were able to get the spin parameter up to the measured $a_\star > 0.98$ with hypercritical accretion [29] (see the discussion on p.355. of Bethe et al. [10]).

If the assumption that the rate of accretion is limited to the Eddington limit is suppressed, we observe that the binary Cyg X$-$1 in Podsiadlowski et al. [114] would be able to transfer up to $\sim 30M_\odot$ during the first thermal timescale (assuming there is that much mass in the system), and another few solar masses afterwards, during the period where the black hole and the secondary star are detached, before the secondary fills again its Roche lobe during the red giant stage.

Cyg X−1, V4641 Sgr and GRS 1915+105 are similar in that the donors in all cases were more massive than the black hole at the time the black hole was formed. The hypercritical accretion for GRS 1915+105 is necessary to bring $a_\star$ up to $a_\star > 0.98$ [29]. For the purposes of discussing Cyg X−1 including the hypercritical nature of the accretion the detailed evolution of Podsiadlowski et al. [114] is useful. Given hypercritical accretion, M33 X−7 and LMC X−1 can be straightforwardly discussed in a similar way as we show in the next sections.

In Sec. 6.3, we discuss what would be the consequences if the current spin of M33 X-7 were natal. We discuss a few problems in this scenario. In Sec. 6.4 we discuss the case for hypercritical accretion in M33 X-7 as an alternative way of making high spin of black hole in M33 X-7. Similarly, we discuss the consequences of a high natal spin in LMC X−1 in Sec. 6.5 and how its evolution would proceed if hypercritical accretion is allowed in Sec. 6.6. We summarize our conclusions in Sec. 6.7.

6.3 Consequences in M33 X−7

In this section, we ask for M33 X−7 what the consequences would be were the currently observed spin of the black hole all natal.

Most important for the binary evolution is that the helium-star (progenitor of the black hole) is spun up by the secondary star so that these "helium-stars will be fully synchronized with their orbital motion throughout their core-helium burning"; i.e., there is tidal locking of the helium-star with the secondary star [137]. Hence, the spin of helium-star and the orbital motion of binary being locked together, and the angular momentum of He-star is transferred to that of black hole as the helium star falls into the latter [81]). In that case, with the currently measured spin parameter $a_\star = 0.77$, the preexplosion orbital period of M33 X-7 would be essentially the same as for Nova Sco, which Lee et al. [81] predicted to be 0.4 days with spin parameter ~ 0.75 (see Figure 12 of Lee et al. [81]). This prediction was confirmed by Shafee et al. [124] with the measurement of $a_\star = (0.65 - 0.75)$ for Nova Sco.

Here we summarize a few problems with this scenario.

a) This tidal locking leads to a strange and complex situation for M33 X-7. Because of the short 0.4 day period, the helium star squashes inside of the $\sim 70 M_\odot$ star, with orbital separation $a \sim 10 R_\odot$ (McClintock, private communication). The tidal locking should be more effective in this case because these stars are much more massive and close together than those considered by Van den Heuvel & Yoon [137].

b) With high spin angular momentum of helium-star (black hole progenitor), the black hole is born in the Blaauw-Boersma explosion,[2] in which black hole binary can have system velocity due to the sudden mass loss during the explosion. With given preexplosion orbital period ~ 0.4 day and the present one of 3.45 days, M33 X-7 should have lost more than half of the system mass and couldn't survive the explosion if there were no hypercritical accretion as we discuss below.

In the case of Nova Sco, the explosion involved a mass loss of several solar masses [103]. The heliocentric radial system velocity of Nova Sco is -150 ± 19km s^{-1}. After correction for peculiar motion of the sum and differential Galactic rotation, the magnitude of the velocity stands out as being higher than any other dynamically identified Galactic black hole candidate [19]. Given the donor mass of $\sim 2 M_\odot$ and the black hole mass of $(5.1 - 5.7) M_\odot$, it lost nearly half of its system mass in the explosion [103]. The reason why Nova Sco is the most energetic explosion among the soft X-ray transient sources is that the explosion energy has to be big enough to expel nearly half of its system mass. We believe that this energy was provided by the black hole spin. The present remaining rotational energy is 430 bethes (1 bethe $= 10^{51}$ ergs). Lee et al. [81] found that in Nova Sco most of the rotational energy is natal.

Given the same spin parameter in the natal spin of M33 X$-$7, it would have ~ 3 times more rotational energy than Nova Sco, because of the ~ 3 times more massive black hole, about half of $M_\odot c^2$! In between the explosion and the present time, no forces act on the binary assuming the (negligible)

[2]For a detailed description of Blaauw-Boersma Kick, and a source of the relevant formula 6.1 see the appendix in Brown et al. [27]

sub-Eddington rate of accretion. In other words, the explosion must convert the originally 0.4 day period into the present one of 3.45 days. We take the black hole mass after the explosion to be the present one, since accretion at the Eddington limit changes its mass negligibly in 2-3 million years.

In the Blaauw-Boersma explosion, assuming rapid circularization,

$$P_2 = \left(1 + \frac{\Delta M}{M_{BH} + m}\right)^2 P_1 \tag{6.1}$$

where the black-hole mass $M_{BH} = M_{He} - \Delta M$ with the mass loss ΔM during the explosion, m is the mass of the secondary star at the time of the explosion, P_1 the pre-explosion period, and P_2 is the post-explosion period. It is well known that once the mass loss is half of the system mass ($\Delta M_{\rm breakup} = M_{BH} + m$) the binary becomes unstable and breaks up. This happens at

$$\left(\frac{P_{breakup}}{P_1}\right) = \left(1 + \frac{\Delta M_{\rm breakup}}{M_{BH} + m}\right)^2 = 4. \tag{6.2}$$

With $P_1 \sim 0.4$ days for $a_\star = 0.77$, the breakup period is 1.6 days, less than half the present 3.45 days observed in M33 X$-$7; i.e., the binary would break up during the explosion. Thus, there must be less mass loss in the explosion and mass must be transferred from the secondary star to the black hole following the explosion, as discussed by Bethe et al. [10], in order to achieve the observed spin parameter.

From the above consideration, we believe that the present value of spin parameter $a_\star = 0.77$ for M33 X-7 cannot be the natal one.

Nova Sco is completely different in the respect that the companion is $\sim 2M_\odot$, much lighter than that in M33 X-7, and the most of the black hole spin energy is natal. In Nova Sco, with the same natal period of $P_1 \simeq 0.4$ days, the breakup period should be the same $P_{breakup} = 1.6$ days. The $\sim 2M_\odot$ secondary star in Nova Sco is some billions of years old, so that even with accretion limited by Eddington, it could have transferred appreciable mass to the black hole. Lee et al. [81] found this to be $0.41 M_\odot$ which, if conservative, would have increased the period of Nova Sco following the explosion to be 1.5 days. It is the proximity of this period to the 1.6 days (breakup period) which makes the system velocity of Nova Sco to be higher than any other Galactic black hole candidate.

6.4 Evolution of M33 X−7

In the previous section, we have discussed that the present value for the spin parameter cannot be the natal one. In addition, from Figure 12 of Lee et al. [81], we see that the 3.45 day period corresponds to a low natal spin parameter $a_\star \sim 0.12$ which is much lower than the observed one $a_\star = 0.77$. So we believe that the spin up of the black hole has to be caused by the accretion from the companion. Knowing that the present day orbital period and spin parameter of the black hole in M33 X-7, we can estimate from Figure 6 of Brown et al. [25] that about $(40 - 50)\%$ of the mass of the black hole had to be accreted after its formation.

Now we can obtain the orbital period before the accretion in M33 X−7 assuming conservative mass transfer,

$$P_3 = \left(\frac{M_{BH,4} \times m_4}{M_{BH,3} \times m_3}\right)^3 P_4 \tag{6.3}$$

where subindex 3 indicates the recircularized values before the accretion starts and subindex 4 indicates the present values. Assuming that the black hole accreted $5M_\odot$ from its companion after its formation, one can obtain

$$P_3 = \left(\frac{15.65M_\odot \times 70M_\odot}{10.65M_\odot \times 75M_\odot}\right)^3 3.45 \text{ days} = 8.9 \text{ days} \tag{6.4}$$

or a spin parameter of $a_\star \sim 0.05$. This, of course, obligates us to reconstruct our calculation in eq. (6.1), nevertheless, the preexplosion period is no longer restricted by the present-day spin parameter and the mass loss can be much smaller as we discuss below.

The black-hole progenitor star had to be more massive than the secondary star in order for it to evolve into a black hole at least a few million years before the black hole formation and accrete hypercritically $\sim 5M_\odot$ after the black hole formation so we could observe the present-day configuration of the system. Given such massive stars, we know that the mass loss through winds has to be considerable. So we know that the ZAMS mass of the black-hole progenitor had to be larger than that of the secondary star, and that the ZAMS mass of the secondary star had to be larger than its mass anytime afterwards, i.e. $M_{\rm ZAMS} > m_{\rm ZAMS} > m_{\rm pre\ explosion} \gtrsim m_{\rm after\ explosion} \gtrsim m_{\rm before\ accretion} > m_{\rm after\ accretion} + 5M_\odot$ where M denotes the mass of black-hole progenitor and m denotes the mass of the secondary star. We have assumed that the mass of the secondary only changes drastically when it fills its Roche lobe for the first time (the second time will occur when it becomes a red giant, but the amount of mass transfer is much smaller) after the explosion of the primary star as explained by Podsiadlowski et al. [114]: *"After the brief turn-on phase, mass transfer occurs initially on the thermal timescale of the envelope reaching a peak mass-transfer rate of* $\sim 4\times10^{-3} M_\odot yr^{-1}$*."*, at which point it transfers hypercritically and in a conservative way, i.e., with little or no mass loss, $5M_\odot$ to the black hole. This means the ZAMS mass of the black-hole progenitor should be around $90M_\odot$. And probably between 10 and $35M_\odot$ right before the explosion, depending on the intensity of the winds (see **?**]). This means that ΔM in equation 6.1 must be between 0 and $25M_\odot$. So,

$$P_1 = \left(1 + \frac{\Delta M}{10.65 + 75}\right)^{-2} 8.9 \text{ days} \tag{6.5}$$

implies 5.3 days $\leq P_1 \leq$ 8.9 days, or a natal spin parameter in the $0.05 - 0.1$ range for the black hole.

This result shows that the energy available for the Blandford-Znajeck mechanism to produce an explosion is very small. Following Lee et al. [80], the black hole spin energy which can be extracted is given as,

$$E_{BZ} = 1.8\times10^{54} \epsilon_\Omega f(a_\star)\left(\frac{M}{M_\odot}\right) \text{ergs} \tag{6.6}$$

where

$$f(a_\star) = 1 - \sqrt{\frac{1}{2}\left[1 + \sqrt{1 - a_\star^2}\right]}. \tag{6.7}$$

Here $\epsilon_\Omega = \Omega_F/\Omega_H$ is the efficiency of extracting rotational energy which, for an optimal process is ~ 0.5, where Ω_F being the angular velocity of the magnetic field, and Ω_H the angular velocity of the black hole. We obtain "only" (as compared with the hundreds of bethes available in the Galactic transient sources [29]) between 3 to 11 bethes of available energy.

This means that most likely the explosion was a dark one and the amount of expelled material was small if not zero, analogous to that proposed by Mirabel & Rodrigues [97] for Cyg X−1.

6.5 Consequences in LMC X−1

The present day Kerr parameter of the black hole in LMC X−1 is $a_\star = 0.90^{+0.04}_{-0.09}$ and its orbital period is 3.9093 ± 0.0008 days (we shall use 3.91 days for our calculations), which means the distance between black hole and companion is roughly $36R_\odot$ or slightly more than twice the $17R_\odot$ radius of the companion, barely confined to its Roche lobe as Orosz et al. [110] point out. Given the present day mass of the secondary star ($30.62 \pm 3.22 M_\odot$) we know this system cannot possibly be much older than $\sim$ 5 million years old [60]. And not allowing for hypercritical accretion implies that the mass of the black hole has not been really altered since it was born.

Let us follow the assumption that this system has not had enough time to spin the black hole by mass transfer given the restrictions imposed by the Eddington limit on accretion onto a stellar mass black hole ($M_{acc} \lesssim 10^{-8} M_\odot$ yr^{-1}). We know from Van den Heuvel & Yoon [137] that such massive stars must be synchronized (i.e., the orbital and spin periods are the same) at the time of the explosion, when the orbit is smallest (since no mass has been lost yet to an explosion and the common envelope has steadily shrunk the orbit from several hundred solar radii to just a few). Lee et al. [81] predicted the Kerr parameters of GRO J1655−40 and 4U 1543−47 by assuming the synchronized core collapses without any angular momentum loss. And so we follow the same procedure here.

A natal Kerr parameter such as $a_\star = 0.90^{+0.04}_{-0.09}$ constraints the orbital period of the pre-explosion binary to a value smaller than 0.4 days (actually closer to 0.3 for the upper values of these measurements, see Fig.1 of Moreno Méndez et al. [100]), which for the current masses would imply an orbital separation smaller than $8R_\odot$, even assuming a tremendous mass loss ($86 M_\odot$, which we will justify in the following discussion) during the explosion, this distance is not larger than $11.5R_\odot$. This is, like in the scenario pictured for M33 X−7 in Moreno Méndez et al. [100], quite an unlikely situation.

Next, the formation of the black hole must be such that the $P_1 \lesssim 0.4$ days is transformed to the present $P_2 = 3.91$ days by the mass loss in the Blaauw-Boersma explosion [12, 17]. Nevertheless, this has, as in M33 X−7, unrealistic expectations since, as stated in Moreno Méndez et al. [100],

$$\left(\frac{P_{breakup}}{P_1}\right) = \left(1 + \frac{\Delta M}{M_{BH} + m}\right)^2 = \left(1 + \frac{M_{BH} + m}{M_{BH} + m}\right)^2 = 4. \tag{6.8}$$

and our scenario needs $P_2/P_1 \sim 10$ (or a mass loss of $\sim 86 M_\odot$!), that is, the binary system would not survive such a big mass loss, if there were enough mass to go through such a scenario in the first place.

Therefore, the evolution of the system must proceed with much less mass loss, the period expanding not further than to some 1.5 days, and mass must be transferred and accreted hypercritically into the black hole.

6.6 Evolutionary Path Of LMC X−1

The present day period of the binary is ~ 3.91 days which stands for a Kerr parameter of $a_\star \sim 0.1$ in Fig. 1 of Chapter 2 [99]. Mass transfer after the formation of the black hole can only occur from the now more massive companion towards the black hole, which implies the orbit could only have shrunk from the time of the collapse until the present day. This means that $a_\star \sim 0.1$ is an upper limit for the natal Kerr parameter of the black hole and so most of the present measured value has to be acquired through hypercritical mass accretion, that is, the original mass has to grow somewhere between 50 to 80% (or some 4 to $5 M_\odot$) from its natal mass to acquire its $a_\star = 0.90^{+0.04}_{-0.09}$ value.

Using the subindices $_3$ to represent the recircularized values, and $_4$ to indicate the present day values we can devolve LMC X−1 assuming conservative mass transfer:

$$P_3 = \left(\frac{M_{BH,4} \times m_4}{M_{BH,3} \times m_3}\right)^3 P_4 = \left(\frac{10.3 M_\odot \times 30.6 M_\odot}{6.3 M_\odot \times 34.6 M_\odot}\right)^3 3.91\text{days} = 11.8\text{days} \tag{6.9}$$

(or 18.22 days if we transfer $5 M_\odot$) which implies $a_\star < 0.05$. Nevertheless the pre-explosion period might have been somewhat smaller depending on the amount of mass lost during the explosion, respecting however the limit imposed by equation 6.5, i.e. the original period had to be larger than 3 days (in which case ΔM is similar to the present mass of the binary, i.e. $\sim 41 M_\odot$) or $a_\star < 0.15$ at the time of the collapse. More likely the mass loss is much smaller, since the available energy for the Blandford-Znajeck mechanism to produce an explosion is very small (formulae from Lee et al. [80]):

$$E_{BZ} = 1.8 \times 10^{54} \epsilon_\Omega f(a_\star) \left(\frac{M}{M_\odot}\right) erg \tag{6.10}$$

where

$$f(a_\star) = 1 - \sqrt{\frac{1}{2}[1 + \sqrt{1 - a_\star^2}]} \tag{6.11}$$

and

$$\epsilon_\Omega = \Omega_F / \Omega_H \tag{6.12}$$

is the efficiency of extracting rotational energy which, for an optimal process is ~ 0.5, where Ω_F being the angular velocity of the magnetic field, and Ω_H the angular velocity of the black hole. The available energy being below 2 bethes for the GRB and the hypernova, well below the hundreds of bethes available in the galactic sources [29].

Like in the case of M33 X$-$7 [100], the available information on LMC X$-$1 points towards a dark explosion where little or no mass was lost when the black hole was formed. Favoring the scenario of a pre- (and post-) explosion period roughly between 12 and 18 days which has been shortened, by hypercritical mass accretion onto the black hole, down to the presently observed 3.45 days.

6.7 Conclusions

In this letter, we discussed a few problems were the currently observed spins of the black holes in M33 X-7 and LMC X$-$1 all natal. Firstly the black-hole progenitors overlap with the companion stars, and secondly the binaries will break up during the explosions in which the black holes were born.

We suggest that the hypercritical accretion has happened in both, M33 X$-$7 and LMC X$-$1, after the black holes were formed spinning them up from low Kerr parameter values ($a_\star > 0.1$)to the high Kerr parameters observed today. LMC X$-$1 and even more so, M33 X$-$7 are the ideal systems to test hypercritical accretion on. Since $\dot{M}_{Edd} \equiv L_{Edd}/c^2 = 4 \times 10^{-8} M_\odot \mathrm{yr}^{-1}$ the necessary accretion to have the black holes torqued up by their companions require ~ 100 and ~ 120 million years respectively to achieve the 4 and $5 M_\odot$ necessary to spin the black holes up to Kerr parameter values of $a_\star = 0.90^{+0.04}_{-0.09}$ and $a_\star = 0.77$ [76] as we suggested. The age of the system is however only 5 and 2 to 3 million years [109].

We think that hypercritical accretion was already established in Houck & Chevalier [68], Brown & Weingartner [23] and Chevalier [38] by the disappearance of SN 1987A. However, it is good to have further proofs as given by M33 X-7 and LMC X−1.

Chapter 7

A Theory of Short-Hard GRBs.

7.1 Introduction

A theory of short-hard gamma-ray bursts (SHBs) in which the central engine uses the same Blandford-Znajek [13] theory as in the theory of the formation of black holes in the Galactic soft X-ray transient sources is presented here. In this theory the rotating black hole (BH) is the central engine for the γ-rays.

The comprehensive review by Nakar [102] is followed. A major difference between the long and short gamma-ray bursts (GRBs) is that the engine of the long GRBs is believed to operate at the center of a collapsing star, while in all current SHB-progenitor models the engine is "exposed". Differences between long and short GRB afterglows are expected mostly due to different total energies, different properties of the ambient medium and possibly, different afterglow geometries. In general, the central engines seem to have great similarity except in available energy and energy fluence.

The main sequence or evolved star companions in the soft X-ray transient sources are replaced by the neutron star (NS) in SHBs. Because of the binary evolution paths followed by the binaries that produce these two GRB populations, we expect the black holes in the transient sources (where the BH is formed from the collapse of a massive star) to be a few times more massive than those in the NS-BH binaries (where the BH might be formed from the mass transfer of a few tenths of a $M_\odot$ into a first born NS, these are called low-mass black holes, or LMBHs). Because the Blandford-Znajek energy depends linearly on the black hole mass, the SHBs are, therefore, an few times less energetic, for equal spin parameters, than the longer GRBs, but otherwise similar in many respects.

The main point of this note is that the restrictions on the masses of the NS-NS binaries are so severe that there are many more, a factor of ~ 5, BH-NS binaries. These form a viable alternative for the SHB progenitors.

The scenario in which many more binaries evolved into 'low-mass BH'-NS binaries was generally discounted, because the latter had not been observed (for reasons we discuss later). Armitage & Livio [2], in a two-dimensional hydro calculation, showed that an accretion disk reformed inside of the accretion shock, allowing matter to accrete onto the NS. They suggested that jets might drive the hypercritically-accreting matter off, saving the NS from collapsing into a BH. This was only one of many arguments made against the Bethe & Brown [8] scenario for forming low-mass BH, NS binaries, so it is important to show that observations require the two NSs in a binary come from two ZAMS mass progenitors that are within 4% of each other in mass.

This point could have been made already in 1995 since Brown [22] showed that the progenitors of a NS-NS binary must have ZAMS masses within 4% of each other so that they burn He at the same time and avoid common envelope which otherwise evolves the first-born NS into a 'low-mass BH'. The 4% restriction is what limits the number of NS-NS binaries. The Brown [22] work was essentially a prediction, because only two or three binary NSs with well measured masses were known at that time, the near equality in masses in each binary being viewed as a curiosity.

However, now 7 binaries with well-measured NS masses have been observed (see references in Table 7.1) and the close difference in measured mass of the NSs within each of the binaries appears to be a general result.

In fact, observers have pointed out a "crawl space". The mass differences in the lower-mass NS-NS binaries like the double pulsar, J0737−3039, is $\sim (0.1-0.2)M_\odot$, far greater than would follow from the 4%-mass difference in the ZAMS masses. One purpose of this note is to show that this $(0.1-0.2)M_\odot$ difference does not come from the ZAMS masses of the progenitors, but results from the He-red giant shell burning that follows the He core burning.

In calculating the common envelope evolution when the ZAMS masses of the progenitors are more than 4% different, we find that the probability of

BH-NS binaries is 84% compared with 16% for the NS-NS binaries. One may question why the former are not seen since, as we shall show, they do participate in the SHBs.

The random probability that all 7 well-measured NS-NS binaries would end up within 4% of each other in mass is $(0.16)^7 = 2.7 \times 10^{-6}$. The more correct statement is that the ZAMS masses of their progenitors are within 4% of each other, but these masses are not directly observable.

A much more detailed and complete double core evolution of nearly-equal-mass NS-NS binaries has been carried out by Dewi et al. [44]. They point out all of the difficulties in such a calculation, but the resulting $\lesssim 4\%$ difference in the two progenitors of double binaries unquestionably occurs in the progenitors.

One of the problems with the merger of NS-NS model to explain SHBs is that no spectral lines of excited elements are seen during a SHB. The same model with a BH-NS binary, can sweep the elements behind the event horizon, solving this problem.

In Section 7.2 we will bring up to date the work of Brown [22] which shows that in the evolution of binary NSs, the ZAMS masses of the two giant progenitors must be within 4% of each other so that they burn He at the same time. Otherwise, the first-born NS is in common envelope with the companion, and accretes sufficient matter to evolve into a BH. Since this requirement for evolving NS-NS binaries is very stringent, the usual result is to form a BH, i.e. a BH-NS binary. As noted above, the ratio of BH-NS binaries to NS-NS ones is $\gtrsim 5$ to 1.

The argument of Brown [22] has become beclouded because small changes in the first-born-NS masses in binaries may occur because its progenitor may accrete matter (during its CO core stage) when the companion fills its Roche Lobe (during the He-shell burning stage of the second-born-NS). This results in the He star which accepts the transferred mass adding $\sim (0.1 - 0.2)M_\odot$ to its mass. We identify these in a schematic model. This does not affect the constraint that the giant progenitors of the two NSs in a binary must be within $\sim 4\%$ of each other in mass. We show that the number of well-measured masses of NS-NS binaries has now grown to 7 and that the near equality in their masses by chance alone would occur with probability $P \sim 2.7 \times 10^{-6}$.

Name	M_1 $[M_\odot]$	M_2 $[M_\odot]$	Difference [%]	Pulsar period [ms]	Ref.
B1534+12	$1.3332^{+0.0010}_{-0.0010}$	$1.3452^{+0.0010}_{-0.0010}$	0.90	37.9	(1)
J1906+0746	$1.248^{+0.018}_{-0.018}$	$1.365^{+0.018}_{-0.018}$	9.0	144.1	(2)
J0737−3039	$1.337^{+0.005}_{-0.005}$	$1.250^{+0.005}_{-0.005}$	6.7	22.72	(3)
J1756+2251	$1.40^{+0.02}_{-0.03}$	$1.18^{+0.03}_{-0.02}$	17	28.5	(4)
J1829+2456	$1.30^{+0.05}_{-0.05}$	$1.27^{+0.11}_{-0.07}$	2.3	4.1	(5)
B1913+16	$1.4408^{+0.0003}_{-0.0003}$	$1.3873^{+0.0003}_{-0.0003}$	3.8	59	(1)
2127+11C	$1.349^{+0.040}_{-0.040}$	$1.363^{+0.040}_{-0.040}$	1.0	30.5	(6)
J1518+4904	$1.56^{+0.13}_{-0.44}$	$1.05^{+0.45}_{-0.11}$	39	40.9	(7)
J1811+1736	$1.62^{+0.22}_{-0.55}$	$1.11^{+0.53}_{-0.15}$	39	104	(1)

Table 7.1: Compilation of double-neutron-star binaries. The first seven NS-NS binaries are included in Fig. 7.1, the last two have error bars which are to large to tell whether the ZAMS masses of the progenitors were within 4% or not. The % of the mass difference is calculated from the average of the two masses. The references in the last column are as follows: (1) Stairs [131]; (2) Kasian et al. [74]; (3) Lyne et al. [87]; (4) Faulkner et al. [48]; (5) Champion et al. [36]; (6) Bethe et al. [11], period from [117]; (7) Nice et al. [104].

The BHs are established by the NS-NS absence at higher mass difference.

7.2 Evolution of Neutron-Star Binaries

Brown [22] outlined the evolution that leads to binary NSs. If they burn He at the same time, the first star to begin burning it sends the (hot) He to the companion, but then as the companion starts burning it, it sends the He back. A common envelope is established by the He between the two carbon-oxygen cores, before the He is expelled from both stars. A crucial development by Braun & Langer [20] was that the time for He burning is too short for the He to be accepted by either core, so that there is near equality of the CO core masses, given that of the progenitors. The detailed evolution including the role of hydrogen is given in the original paper [20].

In Table 7.1 we show the masses of the well-measured NS-NS binaries. Obviously these NS masses do not follow precisely the 4% proximity from the ZAMS masses although they are strikingly close. This is more visually shown in Fig. 7.1 constructed from the calculations of Lee et al. [82] using hypercritical accretion. As noted earlier, the probability that by chance all 7 masses lie close to the lower curve, the companion mass, is 2.7×10^{-6}, vanishingly small.

There are, however, details in the Table which indicate one must go beyond the work of Brown [22]. We have to invoke He shell burning to transfer mass from the first-pulsar progenitor to the companion and this additional mass must be accepted by the companion so as to take it above the range of He-star masses that undergo He shell burning.

He shell burning takes place in only the lower-mass He stars (we choose the dividing mass to be that which results in a NS of $1.335 M_\odot$). The surface energy loss by wind in the lower-mass stars is more important in the He burning, compared with the volume, than in heavy He stars. The lower-mass stars have to burn hotter, and, therefore, have shell-burning which follows the core burning. A good review of helium shell burning is given by Van den Heuvel [136], Sec. 4.8.

From Van den Heuvel [136] one sees that the He shell burning leads to an extension of the He envelope to $\sim 250 R_\odot$ for a $2 M_\odot$ He star, but only to $\sim 1 R_\odot$, the order of magnitude of the core He star, for a $4 M_\odot$ He star. This corresponds to ZAMS masses of $\sim 10 M_\odot$ and $15 M_\odot$, respectively.

In the case of double pulsar 0737−3039A,B the helium red giant mass transfer was included in the evolution of Dewi & van den Heuvel [43] and by Willems & Kalogera [142] and Willems et al. [143]. Helium is accreted onto the higher mass pulsar by Roche transfer in increasing the mass difference between the two NSs. The same is true in J1756−2251. Lee et al. [82] carried out calculations in hypercritical accretion to check that the necessary accretion to obtain the pulsar-companion difference in masses, assuming the ZAMS mass of the progenitors to be within 4% of each other. It should be noted that the hypercritical accretion can be carried out essentially exactly, following Belczynski et al. [5].

From Fig. 7.1 we see that stars in the 84% probability region, which follows when the two progenitors do not burn He at the same time, would have an added mass of 0.6 to $0.8 M_\odot$ due to the accretion from the companion in common envelope evolution. The binary evolution of Lee et al. [82] shows that the ratio of BH-NS binaries to NS-NS to be ~ 5. The fact that none of the well measured double NSs populate this more probable region can only be in-

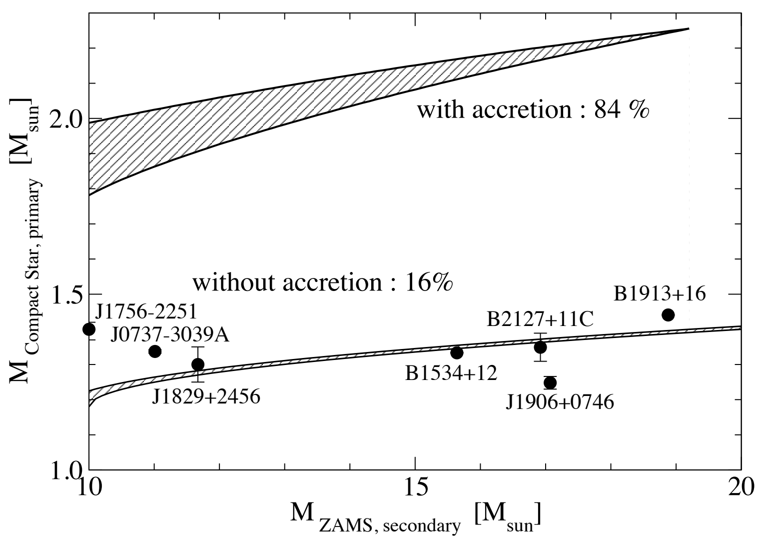

Figure 7.1: Masses of primary compact stars in 7 well-measured NS-NS binaries with and without hypercritical accretion during H red giant stage of secondary star. Note that the 84% corresponds to $M_{\rm compactstar,primary} > 1.8 M_\odot$ [82]. There is uncertainty in the final primary compact star masses due to the extra mass accretion, $\sim (0.1 - 0.2) M_\odot$, during the He giant stage. This may increase the primary compact star masses for both "with" and "without" hypercritical accretion. Note that with our maximum NS mass of $1.8 M_\odot$, all primary compact stars with hypercritical accretion would become 'low-mass BHs'.

terpreted as showing that this region is filled by BH-NS binaries. In principle these should be observable as single pulsars with large Doppler curves, but, as well show in Section 7.4, these will be difficult to find given that the pulsar is never recycled and, therefore, has no observability premium.

As we noted earlier, the large effect of He shell burning has disappeared by the $4M_\odot$ He core, because the He shell burning does not exceed the He core burning in radius. This corresponds to ZAMS $\sim 15M_\odot$. In B1534+12 the masses of pulsar and companion are very close, the mass of the companion being possibly slightly greater. This may come from small fluctuations in the double He star common envelope.

Looking at Table 7.1, J1906+0746 is somewhat different, in that the companion NS is more massive than the pulsar. Of course, earlier in the evolution the (recycled) pulsar progenitor must have been more massive than the companion NS progenitor, in that it would have been the first-born compact object that was speeded up in common envelope evolution. In general, He red giant winds are very fragile and wind could have blown off the matter in the He red giant phase of the pulsar, whereas the red giant phase of the companion would not be blown off by wind, but would add to the mass of the companion. This explanation is obviously contrived and we will watch the observations on this binary with interest.

Above we have made the case that the massive progenitors of binary NSs must be within 4% of each other in ZAMS mass. This prediction by Brown [22] has been observed by all 7 presently well measured NS-NS binaries. This clearly is not just remarkable coincidence, but shows that the maximum NS mass is $< 1.7M_\odot$. This does not quite reach the claim of Bethe & Brown [8] that SN1987A went into a BH with the compact object mass limited to $1.56M_\odot$. However, it is not so far from it, and indicates that ZAMS $\sim 18 - 20M_\odot$ are at the lower limit of BH progenitors. We expand on the consequences of this below.

7.3 Relationship of the Short-Hard Bursts to the Galactic Soft-X-Ray Transient sources and GRB060218/SN 2006aj

We believe we can use similar methods in evolving SHBs to those applicable to the soft-X-ray-transient sources in our Galaxy and to GRB 060218. Lee, Brown & Wijers [81] showed how to calculate the spin parameters of the transient sources, and Brown et al. [29] showed that the subluminous GRBs in neighboring galaxies could come from the soft-X-ray-transient sources in these galaxies.

We use GRB 060218 as an example of GRBs in which we have proposed [29] that they come from soft-X-ray-transient sources. In Lee et al. [81] we showed how to work out the spin parameters of Galactic black-hole binaries and in Brown et al. [29] we gave the results for Galactic binaries. Given the spin parameters we could work out quantitatively the rotational energy of black holes, which in the Blandford-Znajek mechanism [13] could be used to power the GRBs.

In the previous section, we showed that there are $\gtrsim 5$ times more BH-NS binaries than NS-NS binaries. We will show that the same central engine can power these as well as the soft-X-ray-transient sources and long-soft GRBs. The difference is that the collapsing NS, rather than main sequence or evolved companion (we expect about equal numbers in proportion in our Galaxy), occurs in the binary, and spins up the BH on the way to the merger. Of course, we observe the Galactic soft-X-ray binaries before they merge (if they do so within a Hubble time).

A formal scenario for the Bethe & Brown [9] work of NSs merging with BHs has been given by Popham et al. [115]. We choose a model in the middle of their parameter set, because we wish to expand their discussion of the Bethe & Brown [9] scenario.

With the accretion from the NS to the BH during binary merger, the BH can be spun up to a high spin parameter $a^\star \sim 0.8$ [39]. This can be converted into an available rotational energy in Blandford-Znajek mechanism through

$$E_{rot} = \epsilon_\Omega f(a^\star) M_{BH} c^2 \tag{7.1}$$

where the efficiency $\epsilon_\Omega = \Omega_F/\Omega_H$ where Ω_F and Ω_H are the angular velocities of field lines and the BH, respectively, and

$$f(a^\star) = 1 - \sqrt{\frac{1}{2}(1 + \sqrt{1 - a_\star^2})}. \tag{7.2}$$

Usually the efficiency ϵ_Ω is taken from impedance matching to be 1/2. Note that it is not clear at all as to how much of this energy can be accepted. In the soft-X-ray-transient sources, only a small part of a much greater rotational energy can be accepted because the large rotational energy dismantles the central engine after a short time [29]. In LMC X−3 which Brown et al. [30] took to model a cosmological GRB, our estimate was that $\sim 1/2$ of the rotational energy was accepted, but something like half of this was used up in ram pressure work to clear out matter for the γ-rays to leave; i.e., in the gamma-ray efficiency. The ram pressure may not be needed in the SHBs, but there still will be an efficiency to pay. Therefore, given the rotational energy (eq. 7.1), some fraction of it should be measured in the SHBs. This clearly does not mean that we calculate the energy quantitatively, but we do end up in the right ballpark with reasonable assumptions as to parameters. Our main argument is that the SHBs seem to be similar to the longer ones in many ways, only of lower energy and in an environment where they do not have to burst out of a star because they are the merger of a BH-NS binary.

7.4 Population Synthesis

It should be remarked that the concept of a large number of 'low-mass BH'-NS (LMBH-NS) binaries originated with Bethe & Brown [9]. Their Galactic birth rate of compact object binaries of 10^{-4} yr^{-1} was at the high end of those in the literature. Their use of hypercritical accretion to evolve the pulsar in NS-NS binaries into a BH during common envelope evolution with the companion red giant violated, however, the standard scenario for NS-NS evolution accepted at the time.

The main argument against the large number of LMBH-NS binaries was that none had been observed, and this argument is still valid. However, the NS in such a binary is necessarily "fresh", or unrecycled (being the product of the collapse of the second star it has no mass donor to recycle it), whereas the (millisecond) pulsar in the observed NS-NS binaries is "recycled", with its magnetic field and its spin period brought down (two or more orders of magnitude for the former and up to three or four orders of magnitude for the later) by accretion. The "observability premium" of Wettig & Brown [141]

$$\Pi = 10^{12} G/B \tag{7.3}$$

where B is the magnetic field of the pulsar, incorporates the result of recycling [133], namely, that the length of time a pulsar is observable depends linearly on Π. Thus, NS-NS binaries with a recycled pulsar are observable at least two to three orders of magnitude longer than LMBH-NS binaries.

In some papers, the merging of BH-NS binaries is hindered by BH spins and spin-orbit orientation, but in the Blandford-Znajek central engine the spin of the BH is perpendicular to the orbit, so there is no such hinderance, except possibly in the case of binaries with a low-mass black-hole progenitor (of $\sim (18-25)M_\odot$) which go through delayed explosions. (In this case there is an intermediate neutron star which will have a kick velocity).

We cannot produce LMBH-NS binaries with probabilities favoring their observation at the present, but we have shown that for the binaries with two progenitors in the $\sim (10-20)M_\odot$ ZAMS-mass region $\sim 84\%$ of them go into $\sim 2M_\odot$ BHs by our calculation.

Initially the long lifetime of the binaries that merged in the SHBs of $\tau \approx$ 6Gyr [101] awoke some opposition, but we shall show quite simply that this should be considered a lower limit.

Bethe & Brown [9] used the relation of the merging time to the binary separation in the approximation of circular orbits; Shapiro & Teukolsky [125] find (formula 16.4.10)

$$T = \frac{5}{256} \frac{c^5}{G^3 M^2 \mu} R^4 \tag{7.4}$$

where R is the orbital radius, assumed circular for the moment, $M = M_A + M_B$ and

$$\mu = \frac{M_A M_B}{M_A + M_B} \tag{7.5}$$

where M_A and M_B are the masses of the compact objects which merge. For the moment we assume that for a given compact star the probability of forming compact binary system with separation R depends only on volume (random distribution in the local 3-dimensional space), so that

$$4\pi R^2 dR = C\sqrt{T}\frac{dR}{dT}dT = C' T^{-1/4} dT. \tag{7.6}$$

Nakar et al. [101] found that short progenitor lifetimes $\tau \sim 3$ Gyrs are ruled out for the SHBs, and lifetimes of $\tau \sim 6$ Gyrs are favored. This constraints the birthrate to be $> T^{-1/2}$, which is satisfied by our $T^{-1/4}$.

The Bethe & Brown [9] population of merging compact objects of 10^{-4} per galaxy per year is based on our Galaxy. This translates into $1250\text{Gpc}^{-3}\text{yr}^{-1}$. Multiplying by a factor of 8 to take into account the higher star formation rate at redshift $\sim Z = 1$, gives a merger rate $10^4\text{Gpc}^{-3}\text{yr}^{-1}$. We believe from the similarity of these SHB binaries to the soft-X-ray transient sources that the beaming should be $\sim 10\%$ [29], so that the observed rate should be $\sim 10^3\text{Gpc}^{-3}\text{yr}^{-1}$. This estimate falls about midway between the estimates of "at least" $R_{SHB} \approx 10\text{Gpc}^{-3}\text{yr}^{-1}$ and $R_{\text{SHB}} \approx 10^5\text{Gpc}^{-3}\text{yr}^{-1}$ of Nakar et al. [101]. Nakar's later estimate (eq.47 of Nakar [102]) is $10 < R_{SHB} = \text{merger} < 10^4\text{Gpc}^{-3}\text{yr}^{-1}$, where the higher value is a firm upper limit showing that our SHB merging rate is essentially Nakar's upper limit from population syntheses. This conclusion is supported strongly by the observational results in Fig. 1.

Obviously these results will have important consequences for LIGO.

7.5 Conclusions

A theory of SHBs has been presented in which the merger of LMBH-NS binaries are proposed as the progenitors. The central engine is powered by the Blandford-Znajek mechanism in a similar fashion as in the case of the long-soft GRBs. The main difference being that the low-mass companion star in the binary is replaced by a NS, and the collapsing core inside the giant is replaced by a LMBH spun up by the merging NS. The same NS produces the accretion disk which completes the central engine.

The work of Brown [22] and Bethe & Brown [8] has been updated to clarify that the original $\lesssim 4\%$ difference in ZAMS masses for NS-NS binaries can be somewhat altered such that the resulting NS masses in the binary can differ by a considerably larger percent. This effect can be such as to produce binaries where the second-born NS is more massive than the first-born one. However the ZAMS-mass original difference remains unaltered as well as the estimated 5 to 1 ratio of LMBH-NS to NS-NS binaries. This, through population synthesis, gives us a LMBH-NS merger rate within the estimated SHBs rate.

Chapter 8

Conclusions.

8.1 Synopsis

In this work I have reviewed the evolutionary model of binaries with a massive star whose core, through Case C mass transfer and tidal locking, has spun up, allowing for angular momentum from the binary to be transferred back to the black hole progenitor at a time when it will no longer be able to lose it to mass loss through winds (or common envelope evolution). This evolutionary path leads to a core collapse capable of producing, through a Goldilocks scenario for the mass of the companion, dark explosions (like in the cases of Cyg X−1, LMC X−1 or M33 X−7, where the companions are too massive), likely subluminous GRBs (e.g., GRO J1655−40 or 4U 1543−47, where the companions are too light) and, given the right mass companion, it can account even for the brightest and most energetic long-soft (cosmological) GRBs (like, as suggested in Section 4, LMC X−3). This setting also seems to suggest, because of the stripping of the outer layers of the supergiant, that the hypernova explosion will be deficient in hydrogen, explaining in a simple way the Type I_{bc} nature of the explosions.

Performing a population synthesis study, it has been shown as well that our model predicts values well within the observed abundances of both subluminous and cosmological GRBs.

Furthermore, our model is at the stage where estimates for the Kerr parameters of black holes and system velocities for the binaries they reside in can be produced for black holes with low mass companions while more complicated inferences of progenitors and binary evolution can be produced for black holes with high mass companions. The model correctly predicts the independently observed values in two Galactic systems, and it further allows us to calculate estimates for the Kerr parameters (and system velocities) and evolution of the other known Galactic and extragalactic black hole binaries.

We have extrapolated our model in an attempt to shed some light on the mechanisms involved in producing a short-hard GRB on the hypothesis that the merger of a BH-NS binary will produce a central engine similar to those of the long-soft GRBs. We have given, as well, a population synthesis estimate of the rate of short-hard GRB occurrence that is in accordance with those in the literature.

Table 8.1 summarizes the observations and the predictions of our model for the Kerr parameter values in most known black hole binaries where important parameters (such as masses and orbital periods) have been observationally established. It must be noted that these Kerr parameter values are natal, i.e., they could decrease (right after the collapse) if the system powers a strong GRB/HN explosion (although this may have occurred in a somewhat smaller explosion in GRO J1655−40 and 4U 1543−47, the error bars in the measurements are larger than what our model predicts would be lost through the powering up of the explosions). On the other hand, they could substantially increase where mass transfer from the companion star significantly increases the black hole mass (e.g., GRS 1915+105, LMC X−1 and M33 X−7). This latter situation has even brought strong support to the case of hypercritical accretion onto these compact objects.

8.2 Future Avenues

Several important directions need to be further studied regarding the study of black hole binary formation and the study of GRB/HN explosions:

- Our model has brought into focus the need to further investigate hypercritical accretion onto compact objects in situations other than core collapse, with M33 X−7 and LMC X−1 as two examples illuminating the need to include this important phenomenon in future binary evolution with compact object studies. We believe the use of Bondi accretion onto

Name	$M_{BH,2}$ $[M_\odot]$	$M_{d,2}$ $[M_\odot]$	$M_{BH,now}$ $[M_\odot]$	$M_{d,now}$ $[M_\odot]$	Model $a_{\star,2}$	Measured $a_\star$	$P_{Orbit,now}$ [days]	$E_{\rm BZ}$ $[10^{51}$ ergs]	Refs
AML: with main sequence companion									
J1118+480	~ 5	< 1	$6.0 - 7.7$	$0.09 - 0.5$	0.8	-	0.169930(4)	~ 430	(1)
Vel 93	~ 5	< 1	$3.64 - 4.74$	$0.50 - 0.65$	0.8	-	0.2852	~ 430	(2)
J0422+32	$6 - 7$	< 1	$3.4 - 14.0$	$0.10 - 0.97$	0.8	-	0.2127(7)	$500 \sim 600$	(3)
1859+226	$6 - 7$	< 1	$7.6 - 12$		0.8	-	0.380(3)	$500 \sim 600$	(4)
GS1124−683	$6 - 7$	< 1	6.95(6)	$0.56 - 0.90$	0.8	-	0.4326	$500 \sim 600$	(5)
H1705−250	$6 - 7$	< 1	$5.2 - 8.6$	$0.3 - 0.6$	0.8	-	0.5213	$500 \sim 600$	(6)
A0620−003	~ 10	< 1	11.0(19)	0.68(18)	0.6	-	0.3230	~ 440	(7)
GS2000+251	~ 10	< 1	$6.04 - 13.9$	$0.26 - 0.59$	0.6	-	0.3441	~ 440	(8)
Nu: with evolved companion									
GRO J1655−40	~ 5	$1 - 2$	$5.1 - 5.7$	$1.1 - 1.8$	0.8	$0.65 - 0.75$	2.6127(8)	~ 430	(9)
4U 1543−47	~ 5	$1 - 2$	$2.0 - 9.7$	$1.3 - 2.6$	0.8	$0.75 - 0.85$	1.1164	~ 430	(10)
XTE J1550−564	~ 10	$1 - 2$	$9.68 - 11.58$	$0.96 - 1.64$	0.5	-	1.552(10)	~ 300	(11)
GS 2023+338	~ 10	$1 - 2$	$10.3 - 14.2$	$0.57 - 0.92$	0.5	-	6.4714	~ 300	(12)
XTE J1819−254	$6 - 7$	~ 10	$8.73 - 11.69$	$5.50 - 8.13$	0.2		2.817	$10 \sim 12$	(13)
GRS 1915+105	$6 - 7$	~ 10	14(4)	1.2(2)	0.2	> 0.98	33.5(15)	$10 \sim 12$	(14)
Cyg X−1	$6 - 7$	$\gtrsim 30$	~ 10.1	17.8	0.15	-	5.5996	$5 \sim 6$	(15)
Extragalactic									
LMC X−1	~ 10	~ 35	$8.96 - 11.64$	30.62 ± 3.22	< 0.05	$0.81 - 0.94$	3.91	< 2	(16)
LMC X−3	7	4	$5 - 11$	6 ± 2	0.43	< 0.26	1.70	~ 155	(17)
M33 X−7	~ 15	~ 80	$14.20 - 17.10$	70.0 ± 6.9	~ 0.05	$0.72 - 0.82$	3.45	$3 - 11$	(18)

Table 8.1: Parameters at the time of formation of the black hole and at present time. Subindex 2 stands for values at the time the black hole is formed, whereas subindex *now* stands for recently measured values. For LMC X−1 we estimate a $M_{ZAMS} \sim 40 M_\odot$, and for M33 X−7 we estimate a $M_{ZAMS} \sim 90 M_\odot$ for the black hole progenitors. $E_{\rm BZ}$ is the rotational energy which can be extracted via Blandford-Znajek mechanism with optimal efficiency $\epsilon_\Omega = 1/2$ (see Appendix 8.3). REFERENCES: (1)[90], [139], [99]; (2)[49], [99]; (3)[3], [99]; (4)[50], [99]; (5)[3], [56], [99]; (6)[3], [99]; (7)[3], [57], [99]; (8)[3], [62], [99]; (9)[3], [4], [124], [99]; (10)[106], [108], [124], [99]; (11)[107], [99]; (12)[3], [121], [122], [99]; (13)[107], [99]; (14)[61], [93], [99]; (15)[66], [99]; (16)[60], 6; (17)[42], [30]; (18)[86], [100].

compact objects is justified under certain conditions where the energy and angular momentum of the infalling material can be dissipated away fast enough.

Furthermore, detailed studies (in which the compact object is embedded in a stream of material, at different densities and velocities) are necessary in order to justify the removal of angular momentum during the accretion process.

- New GRB/HN data is being constantly obtained, and even more so with the advent of new faster and more sensitive telescopes, so the need to model the progenitors of GRBs (both, short and long) is paramount. Our model is likely the only one which can explain the systems and conditions from which subluminous and cosmological long-soft GRBs come from by "dialing" the donor (secondary star) mass in order to obtain the observed energies and T_{90}'s (time required to deliver 90% of the bursts energy), both of which are calculated in the Blandford-Znajek mechanism according to the available rotational energy in the black hole obtained from our binary evolution model.

  In a similar fashion we can also apply the same concepts to estimate energies for the short-hard GRBs, produced during the merger of a double compact-object binary. Follow up calculations will further help refine the details of our model and will eventually lead to a better estimate of the mass of the companion that establishes the cutoff between subluminous and cosmological GRBs.

- The recent SN 2008D/XRF 080109 had certain interesting characteristics which make it oddly similar to SN 1987A, which is argued to have produced a $\sim 1.5 M_{\odot}$ black hole. If this is the case, there is a very likely possibility that the equation of state (EOS) of nuclear matter at high densities will be softened by the appearance of a kaon condensate. Given that the symmetry energy can now be pinned down by utilizing recent results from the application of Brown-Rho scaling (BRS) to the ^{14}C problem, the study of a repulsive EOS, due to BRS, preceding kaon condensation, might be another promising avenue for the study of both, the nuclear EOS and the maximum mass of a compact star.

- The next step in evolution of massive binaries is to consider a massive secondary star. The abundance ratio of NS-NS to BH-NS binaries directly depends on the EOS as well as on whether hypercritical accretion can occur onto a compact object during common envelope evolution, therefore this topic has very important consequences to nuclear physics, astrophysics (including gravitational wave detection) and cosmology.

- The study of Gamma-Ray Bursts and stellar-mass compact objects directly addresses fundamental questions about the physics of the Cosmos in more than one important way. Not only is their formation and occurrence what gives our Universe the observed abundances of heavy elements, but GRBs might as well become one of the best cosmological standard candles once they are better understood. Knowing the boundary mass between the formation of NSs and BHs will give us, through population syntheses, abundances of NSs, BHs, and GRBs as well as those of chemical elements and might help us put better constraints in the cosmological natural selection.

- Further measurements of Kerr parameters are being carried out on soft X-ray transient sources with the X-ray continuum fitting. Different techniques (e.g., Fe K line profile fitting) are already underway, and new ones (X-ray polarimetry and QPOs) will also be applied to the observed black hole binary systems as well as newly discovered ones. Spatial peculiar velocities will also be measured for some of these systems. All these observations will further test our model and will allow us to correct and perfect it to the point where it may become a useful tool explaining the evolution of black hole binaries as the relics of GRB/HN explosions.

8.3 Concluding Remarks

The model I have displayed in this work has produced clear estimates of the values of the Kerr parameters of over a dozen black holes in soft X-ray transient sources (Table 8.1). Of these, 6 have measured Kerr parameters. Two of these sources have matched the predictions and 4 more can be consistently explained by post-explosion accretion from the companion star onto the black hole. It will be very exciting to compare future measurements of these sources with the remaining values presented on Table 8.1. However, as mentioned earlier (Section 8.1), the predicted Kerr parameters might have evolved from the natal ones, which we predict, therefore computation of the post-explosion evolution of these binaries will be required.

Of great relevance and great help will be to have accurate measurements of peculiar system velocities for these binaries, since, as it has been remarked in several places throughout this work, although the mass loss during a GRB/HN event might be cylindrically symmetric with respect to the exploding star, it is not symmetric with respect to the center of mass of the binary, and thus a Blauuw-Boersma kick is imparted to the binary along the orbital plane. The outlined model allows an accurate estimate of these velocities if the mass lost during the explosion can be obtained. Thus, such measurements would provide further constraints on the evolution of each binary and would allow further scrutiny of our evolutionary model for black hole binaries as sources of GRB/HN explosions.

Binary evolution of compact objects allows, in a natural way, the formation of rapidly spinning black holes or neutron stars, which are the energetic central engines for the Blandford-Znajek mechanism. Therefore the evolution of compact object binaries is a very likely scenario for both long-soft and short-hard GRBs. This makes the problem of hypercritical accretion onto black holes and compact stars highly relevant.

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Appendix A: Blandford-Znajek Mechanism.

In Fig. 8.1 we show that a rotating black hole operates like a generator of electricity; this is the Blandford-Znajek [13] mechanism, which is summarized in the caption. We rely on the relatively complete review by Lee et al. [80].

The rotational energy of a black hole with angular momentum J is a fraction of the black hole mass energy

$$E_{\rm rot} = f(a_\star) M_{\rm BH} c^2, \tag{8.1}$$

where

$$f(a_\star) = 1 - \sqrt{\frac{1}{2}\left(1 + \sqrt{1 - a_\star^2}\right)}. \tag{8.2}$$

For a maximally rotating black hole ($a_\star = 1$) $f(1) = 0.29$. In Blandford-Znajek mechanism the efficiency of extracting the rotational energy is determined by the ratio between the angular velocities of the black hole Ω_H and the magnetic field velocity Ω_F,

$$\epsilon_\Omega = \Omega_F / \Omega_H. \tag{8.3}$$

For optimal energy extraction $\epsilon_\Omega \simeq 0.5$. This just corresponds to impedance matching between that in the generator (Fig. 8.1) and that in the perturbative region where the energy is delivered. One can have the analytical expression for the energy extracted

$$E_{BZ} = 1.8 \times 10^{54} \epsilon_\Omega f(a_\star) \frac{M_{BH}}{M_\odot} \text{erg}. \tag{8.4}$$

The rest of the rotational energy is dissipated into the black hole, increasing the entropy or equivalently irreducible mass.

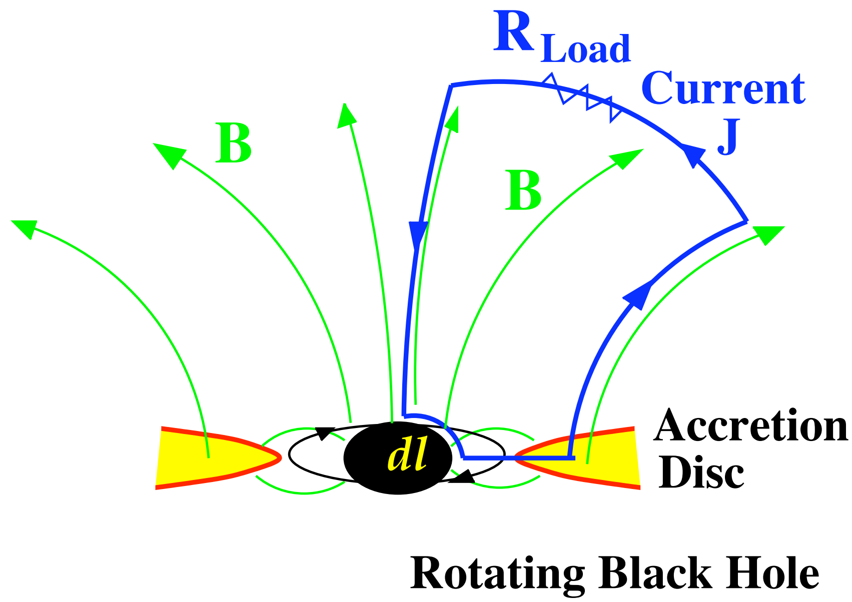

Figure 8.1: The black hole in rotation inside the accretion disk, formed by what is left of the He star. A wire loop can be drawn, coming down a field line from the top (this particular field line is anchored in the black hole.) to the north pole of the black hole. The black hole has a surface conductivity, so the wire can be extended from the north pole of the black hole to the equator and further extended into the (highly ionized) accretion disk, in which the magnetic field lines are frozen. The wire can be continued on up out of the accretion disk along a field line and then connects back up to form a loop. As this wire loop rotates, it generates an electromagnetic force, by Faraday's law, sending electromagnetic radiation in the Poynting vector up the vertical axis. The region shown has condensates of charge designed to allow free rotation of the black hole, although the black hole when formed rotates much more rapidly than the accretion disk, since the compact object angular momentum must be conserved as the inner part of the He star drops into the black hole. In trying to spin the accretion disk up, the rotation engendered by the field lines is converted to heat by viscosity, the resulting hypernova explosion taking place in a viscous time scale. The gamma ray burst is fueled by the deposition of Poynting vector energy which is sent up the rotational axis into a fireball.

Although the use of $\epsilon_\Omega = 0.5$ is close to the actual efficiency for high $a_\star$, it decreases with $a_\star$ from 0.69 at $a_\star \sim 0.8$ to 0.46 at $a_\star \sim 0.4$ ([25], see Ω_{disk}/Ω_H, $r = r_{ms}(\tilde{a})$ where r_{ms} is the marginally stable radius). There would be a further decrease of $\gtrsim 50\%$ to the $a_\star$ of 0.15 of Cyg X$-$1 and M33 X$-$7. This low efficiency virtually ensures that these two binaries will have gone through dark explosions.

The hypernova results from the magnetic field lines anchored in the black hole and extending through the accretion disk, which is highly ionized so the lines are frozen in it. When the He star falls into a black hole, the latter is so much smaller in radius that it has to rotate much faster than the progenitor He star so as to conserve angular momentum.

Initially, large amounts of energy, up to 10^{52} ergs, were attributed to gamma ray bursts. However, when the correction is made for beaming, the actual gamma ray burst energy is distributed about a "mere" $\sim 10^{51}$ ergs [113]. However, $\sim 10^{52}$ergs are required to clear the way for the jet. Hypernova explosions are usually modeled after the nearby supernova 1998bw. The hypernova by Nomoto et al. [105] that Israelian et al. [70] compared with GRO J1655$-$40 had an energy of 3×10^{52} ergs.

Appendix B: Dismantling the Accretion Disk by High Energy Input.

The amount of energy poured into the accretion disk of the black hole is almost unfathomable, the 4.3×10^{53} ergs$\simeq \frac{1}{4} M_\odot c^2$ being 430 times the energy of a strong supernova explosion. Near the horizon of the black hole, the physical situation might become quite complicated [134]. Field-line reconstruction might be common and lead to serious breakdowns in the freezing of the field to the plasma; and the field on the black hole sometimes might become so strong as to push it back off the black hole and into the disk (Rayleigh-Taylor Instability) concentrating the energy even more. During the instability the magnetic field lines will be distributed randomly in "globs", the large ones having eaten the small ones. It seems reasonable that the Blandford-Znajek mechanism is dismantled. Later, however, conservation laws demand that the angular momentum not used up in the GRB and hypernova explosion be reconstituted in the Kerr parameter of the black hole. The radius of the (Kerr) black hole is

$$R = \frac{GM}{c^2} = 1.48 \times 10^6 \frac{M}{10 M_\odot} \text{cm}.$$

Given the above scenario of the very high magnetic couplings dismantling the disk by Rayleigh-Taylor instability (e.g. in GRO J1655−40), we wish to make a "guestimate" of the same effect for XTE J1550−564, since its Kerr parameter is being measured by the Smithsonian-Harvard collaboration.

The black hole radius is proportional to M_{BH}. The ratio of $M_{BH}(1550 -$

564)/M_{BH}(1655 − 40) $\simeq$ 2; but we use a ratio more like 2.5 because GRS 1655−40 has a Kerr black hole and, in XTE J1550−564, the black hole is about halfway between Kerr and Schwarzschild. Thus, the area of the circle inside the last stable circular orbit is $\sim$ 6 times larger for XTE J1550−564. Taking the field strength for Rayleigh-Taylor instability, to go as the inverse of the area, this means that its effect would be cut down by a factor of 6 in going from GRS 1655−40 to XTE J1550−564. If our scenario that the magnetic field coupling is correct for GRO J1655−40, not much of the effect would be left in XTE J1550−564 which should accept most of the energy for a cosmological GRB.

The question then is, what is the latter? We know that SN1998bw had a hypernova energy of $\sim$ 30bethe. From the equipartition of energy, we would estimate the kinetic energy to clear out a path for the jet in the GRB to be about equal to the thermal energy, which is roughly true in MacFadyen [89]. So the total energy would be $\sim$ 60bethes, which we know can be accepted by the binary. MacFadyen [89] suggested that the accompanying GRB 980425 was "smothered", so that may be a lower limit, although it is larger than estimates we have seen to date, so we choose it as the energy of a cosmological GRB.

Now, 60 bethes is $\sim$ 20% of our estimated angular momentum energy for XTE J1550−564. However, the decrease in Kerr parameter necessary to go from 300 to 240bethes is only 0.06 from our natal $a_\star \sim 0.5$, or $\sim$ 12%, which requires an accurate measurement in Kerr parameter, but is none the less much larger than the $\sim$ 6% decrease estimated for GRO J1655−40. (Note that for Schwarzschild black holes, the rotational energy goes approximately as $a_\star^2$.) On the other hand, this may be the minimal difference between natal and present angular momentum energies because no other model leaves the system spinning so rapidly. Thus, if no difference is conclusively seen because observational errors are $\sim$ 12% then this is also very interesting.

Suppose a decrease of more than 12% is seen. Then this means that we have underestimated the energy of the explosion, but our above estimates are as large as any proposed ones. In any case, the possibility that the system, following the explosion, is left rotating is a new and interesting one.

We thus see that we can fit the Fruchter et al. [51] condition for no high luminosity GRBs in our high metallicity stars in it, because all of the donor masses of the GRBs that received the highest rotational energies had companion masses of $\sim (1-2)M_\odot$ and the rotational energy furnished to them was so great that the accretion disks were dismantled. The result is that the GRBs were subluminous, like the vast majority of GRBs. The ~ 6 times larger surface area should be helpful in allowing XTE 1550−564 to accept the rotational energy, but the amount is still tremendous. We do not have any donors of $\sim 5M_\odot$ in our Galaxy, but such a donor would bring the energy down by $\sim 1/5$, since it goes inversely with donor mass, to ~ 60 bethes, that we estimate for GRB 980425, which as noted has nearly Galactic metallicity. Although GRB 980425 is essentially "smothered" [89], 60bethes is the highest energy anyone has attributed to the GRB and Hypernova explosion. Thus, our estimates would say that donors of $\sim 5M_\odot$ would give the most luminous GRBs, and that XTE 1550−564 may or may not have been able to accept $\sim$ 60bethes, but since we view this as an upper limit, this black hole should still be spinning furiously.

We should enter a proviso here. Our main thesis is that the black hole binary must have a donor sufficiently massive to slow it down enough so that the black hole can accept the strong magnetic coupling through its accretion disk without the Rayleigh-Taylor instability entering. By increasing the area of the accretion disk by a factor ~ 6 in going from the $\sim 5M_\odot$ black hole in GRO J1655−40, to the $\sim 10M_\odot$ black hole in XTE 1550−564 the density of magnetic coupling is decreased by a factor ~ 6. However, the donor in XTE 1550−564 is only $\sim 1.3M_\odot$, about the same size as in GRO J1655−40. Therefore, the GRB in XTE 1550−564 may still have been subluminous, but less so than GRO J1655−40 because of the larger disk area.

The donor masses at the time of common envelope evolution of XTE 1819−254 and GRS 1915+105 were $\sim 10M_\odot$. However they have only $\sim 1/3$ of the energy of the binary with $5M_\odot$ donor, $\sim (10-12)$bethes. These black holes accrete a lot of matter, more than doubling the black hole mass in the case of GRS 1915+105 and, ultimately will double that mass in XTE 1819−254. Such binaries with substantial mass exchange do not obey our simple scaling which is designed for natal angular momentum. They must be evolved in detail.

We believe, that similar evolutionary arguments will apply to binaries with higher mass donors, and, anyway, the angular momentum energy will be cut down by the higher donor masses. Thus we expect that only binaries with donor masses $\sim 5M_\odot$ will give highly luminous GRBs with rotational energy ~ 60bethes.

In a very rough estimate, using a flat distribution of donors with mass up to $80M_\odot$ we can estimate that the number of cosmological GRBs, those with high luminosity will be $\sim 5/80$ of the total, not far from the ratio of GRBs with high luminosity to subluminous ones found by Liang et al. [85].

In summary, the commonly accepted estimates of the explosion energies in GRBs are orders of magnitude less than the natal angular momentum energy in GRO J1655−40. The measured Kerr parameter of $a_\star \sim 0.8$ [124] has a present rotational energy indistinguishable from the natal one, within error bars. However, the dismantling of the accretion disk from the very high magnetic couplings should be less in XTE J1550−564 and the rotational energy is somewhat less, the donor being about the mass of the He star in the Woosley Collapsar model, so we propose that XTE J1550−564 can accept substantial rotation energy but probably less than the energy of a cosmological GRB. Measurement of this energy is being carried out by the Smithsonian-Harvard collaboration.

If our suggested scenario is confirmed, then this should be strong support for the Brown et al. [25] binary scenario for GRBs.

Appendix C: Hypercritical Accretion.

Brown & Weingartner [23] calculated analytically hypercritical spherical accretion onto compact objects and in particular for the fall back in SN 1987a and obtained the same result Houck & Chevalier [68] obtained numerically. Along this paper we have used the evolution of Cyg X−1 by Podsiadlowski et al. [114] where they find mass transfers from the secondary to the black hole as large as $10^{-4}M_\odot$ per year for a period on the order of 10^4 years only to limit the accretion to Eddington's limit in the end, ejecting most of the transferred mass out of the binary. We have used the Podsiadlowski et al. [114] path in order to evolve not only Cyg X−1, but also XTE J1819−254 and GRS 1915+105, except we believe a large fraction of this mass gets accreted hypercritically into the black hole. Similarly, Moreno Méndez et al. [100] show that M33 X−7 can only be evolved into its current state [86], if hypercritical accretion takes place in this system. Next we outline the calculation by Brown & Weingartner [23]:

Brown & Weingartner [23] obtain, for SN 1987a, an accretion rate of

$$\dot{M} = 1.15 \times 10^{22}\text{g s}^{-1} = 1.81 \times 10^{-4} M_\odot \text{yr}^{-1}, \tag{8.5}$$

which is the same order of magnitude as the one in Podsiadlowski et al. [114], so we simply follow their results. The Eddington accretion rate is

$$\dot{M}_{\text{Edd}} = \frac{4\pi c R}{\kappa_{\text{es}}} = 5.92 \times 10^{-8} M_\odot \text{yr}^{-1}, \tag{8.6}$$

where $R \simeq 10^6$ cm is the radius of the compact object and κ_{es} is the opacity, which we take to be $\kappa_{\text{es}} \simeq 0.1$ cm^2 g^{-1}. This is an estimate that applies over a range of temperature and density similar to that present here [37]. The Eddington luminosity L_{Edd}, the luminosity for which the pressure of outward traveling photons balances the inward gravity force on the material, is obtained from $L_{\rm Edd} = \dot{M}_{\rm Edd} c^2$. If $\dot{M}$ exceeds $\dot{M}_{\rm Edd}$ then some of the accretion energy must be removed by means other than photons. In the present case,

$$\dot{m} \equiv \frac{\dot{M}}{\dot{M}_{\rm Edd}} = 0.31 \times 10^4. \tag{8.7}$$

When $\dot{M}$ exceeds $\dot{M}_{\rm Edd}$ by so large a factor, the accretion rate is called hyper-critical, and was considered by Blondin et al. [15].

Blondin et al. [15] finds a trapping radius r_{tr} such that photons within $r_{\rm tr}$ are advected inward with accreting matter faster than they can diffuse outward. We follow his derivation in slightly modified form. We start with the same type of equation as is used in deriving the Bondi $\dot{M}$:

$$\dot{M} = 4\pi r^2 \rho v. \tag{8.8}$$

Here, however, we consider any radius r and take v to be the free fall velocity, that is $v = (GM/r)^{1/2}$. We divide equation 8.8 by equation 8.6 to find

$$\rho\kappa_{\rm es} = R r_s^{-1/2} r^{-3/2} \dot{m}, \tag{8.9}$$

where $r_{\rm s} = 2GM/c^2$ is the Schwarzchild radius; $r_{\rm s}$ is about 4.4 km for a $1.5 M_\odot$ compact object. We now find the optical depth for electron scattering to be

$$\tau_{\rm es} \equiv \int_r^\infty \rho\kappa_{\rm es} dr = 2R\dot{m}(r r_{\rm s})^{-1/2}, \tag{8.10}$$

where r is the distance from the compact object where the photon begins its journey.

From random walk, the time for the photon to diffuse over a distance d is

$$\tau_{\rm diff} = \frac{d}{c}\frac{d}{\lambda_{\rm es}} \simeq \frac{d}{c}\tau_{\rm es}, \tag{8.11}$$

where $\lambda_{\rm es}$ is the photon mean free path. The approximation in the last step follows if we assume d is large enough that most of the scatterings that the photon undergoes on its trip to infinity have already occurred by the time it has traveled distance d. In order to obtain the photon trapping radius, we set this equal to the dynamical time, which is the time it would take for a piece of matter at position r to travel the distance d at its instantaneous speed at r, and approximates the time it takes for a photon to be advected inwards distance d:

$$t_{\rm dyn} = \frac{d}{v(r)} = d\left(\frac{r}{2GM}\right)^{1/2}, \tag{8.12}$$

where $v(r)$ is the free-fall velocity at r. Using equation 8.10 for $\tau_{\rm es}$ and substituting $r_{\rm tr}$, the trapping radius, for r we find

$$r_{\rm tr} = 2R\dot{m}. \tag{8.13}$$

Since the solution of the diffusion equation for radial diffusion in three dimensions decreases the diffusion time by a factor of $\pi^2/3$, we must likewise decrease $r_{\rm tr}$ by this factor. We finally obtain

$$r_{\rm tr} = 0.6R\dot{m} = 1.86 \times 10^9 {\rm cm}. \tag{8.14}$$

Any photon flux emitted much below $r_{\rm tr}$ is unable to diffuse upstream and thus can have no effect on the luminosity reaching the observer at infinity.

Chevalier and collaborators took into account that neutrinos can carry away accretion energy and developed self consistent solutions for hypercritical accretion (Chevalier (1989), (1990); [68]). In particular they find an expression for the radius of the accretion shock in terms of $\dot{M}$, for a neutron star. We follow the derivation in Chevalier (1989, p.854) with small modification. We first derive an expression for $p_{\rm ns}$, the pressure at the surface of the neutron star, in terms of $\dot{M}$ and $r_{\rm sh}$, the shock radius. We then examine neutrino cooling near the surface of the neutron star, producing an equation in terms of $p_{\rm ns}$. Insertion of our expression for $p_{\rm ns}$ gives $r_{\rm sh}$ in terms of $\dot{M}$.

Since the pressure is radiation dominated, the accretion envelope forms an $n = 3$ ($\Gamma = 4/3$) polytrope. Thus, inside the shock

$$\rho = \rho_{\rm sh}\left(\frac{r}{r_{\rm sh}}\right)^{-3}; p = p_{\rm sh}\left(\frac{r}{r_{\rm sh}}\right)^{-4}; v = v_{\rm sh}\left(\frac{r}{r_{\rm sh}}\right), \tag{8.15}$$

where the subscript $_{sh}$ refers to the value at the shock front. Because of the adiabatic compression by factor $(\Gamma+1)/(\Gamma-1)$,

$$\rho_{\rm sh} = 7\rho_0, \tag{8.16}$$

where ρ_0 is the density just outside the shock front. We neglected the (small) decrease in Γ because of increased ionization of the material going through the shock. As in Blondin et al. [15], ρ_0 is calculated as follows:

$$\rho_0 = \frac{\dot{M}}{4\pi r_{\rm sh}^2 v_{in}}, \tag{8.17}$$

where $v_{\rm in}$ is the free-fall velocity at the shock radius. From conservation of mass flow across the shock front,

$$v_{\rm sh} = -\frac{1}{7}v_{\rm in}. \tag{8.18}$$

Thus, the kinetic energy of the accreting matter is diminished by a factor of 49; that is, it is almost entirely converted into thermal energy, so we can estimate the thermal energy density as

$$\epsilon_{\rm sh} \simeq \frac{7}{2}\rho_0 v_{\rm in}^2, \tag{8.19}$$

the factor 7 entering because of compression (equation 8.16. In this derivation we neglected $v_{\rm sh}$ as compared with $v_{\rm in}$. The small correction that would be generated had we taken account of the decrease in Γ due to ionization across the shock in obtaining eq. 8.16. Since the pressure is radiation dominated,

$$p_{\rm sh} = \frac{1}{3}\epsilon_{\rm sh} \simeq \frac{7}{6}\rho_0 v_{\rm in}^2. \tag{8.20}$$

Inserting equation 8.17 and $(2GM/r_{\rm sh})^{1/2}$ for $v_{\rm in}$, and using the second equation of 8.15, we find the pressure at the surface of the neutron star

$$p_{\rm ns} = 1.86\times 10^{-12}\dot{M} r_{\rm sh}^{3/2} {\rm dyn\ cm^{-2}}, \tag{8.21}$$

where $\dot{M}$ is expressed in g s^{-1} and $r_{\rm sh}$ in cm. We took the radius of the neutron star to be $\simeq 10$ km.

The energy loss by neutrino pair production per unit volume is [46]

$$\dot{\epsilon}_n = 1.06 \times 10^{25} T^9 C\left(\frac{\mu_e}{T}\right) \text{ergs cm}^{-3}\text{s}^{-1}, \tag{8.22}$$

where μ_e is the electron chemical potential, T is in MeV, and $C(x)$ is a slowly varying function of x which can be computed from the paper of Dicus [46]. For $x = 0$, $C = 0.92$; we shall use this value, because the electrons are not very degenerate. In the region where $\dot{\epsilon}_n$ is operative, $T \sim 1$ MeV so that e^+, e^- pairs as well as photons, contribute to the radiation pressure. With T in MeV, the photon blackbody energy density is

$$w = 1.37 \times 10^{26} T^4 \text{ergs cm}^{-3}. \tag{8.23}$$

Inclusion of the e^+, e^- pairs multiplies this by a factor of 11/4, and we divide by three to obtain the pressure

$$p = 1.26 \times 10^{26} T^4 \text{ergs cm}^{-3}. \tag{8.24}$$

Combining equations 8.22 and 8.24 yields

$$\dot{\epsilon}_\text{n} = 1.83 \times 10^{-34} p^{2.25}, \tag{8.25}$$

where $\dot{\epsilon}_\text{n}$ is in ergs cm^{-3} s^{-1} when p is in ergs cm^{-3}. The neutrino cooling is taken to occur within a pressure scale height $r_\text{ns}/4$ of the neutron star, or in a volume of $\simeq \pi r_\text{ns}^3$. Energy conservation gives

$$\pi r_\text{ns}^3 \times 1.83 \times 10^{-34} p_\text{ns}^{2.25} = \frac{GM\dot{M}}{r_\text{ns}}, \tag{8.26}$$

with, again, everything in cgs units. Inserting equation 8.21 into equation 8.26, we solve for r_sh:

$$r_\text{sh} \simeq 6.4 \times 10^8 \left(\frac{\dot{M}}{M_\odot \text{yr}^{-1}}\right)^{-10/27} \text{cm}. \tag{8.27}$$

The power $-10/27 = -0.370$ is the same as that obtained by Houck & Chevalier [68] using the accurate neutrino cooling function, and not that of Chevalier

(1989).

The detailed calculation of Houck & Chevalier [68] finds the only substantial correction to our schematic calculation to arise from general relativity, which can be taken into account by multiplying the expression for $r_{\rm sh}$ by 0.4. Thus,

$$r_{\rm sh} \simeq 2.6 \times 10^8 \left(\frac{\dot{M}}{M_\odot {\rm yr}^{-1}} \right)^{-0.370} {\rm cm} = 6.3 \times 10^9 {\rm cm}, \tag{8.28}$$

where we have used equation 8.5 for $\dot{M}$.

Brown & Weingartner [23] go on to calculate the critical time after which the neutron star left behind by SN 1987a should be visible and find it is less than a year after the explosion. Up to this moment, no neutron star has been observed. Nevertheless the hypercritical accretion onto a black hole will be just as efficient, if not more than onto a neutron star, since the black hole has no surface against which matter would hit and produce radiation pressure, but only an event horizon which it would ram through without resistance.

Appendix D: Energies.

The energies of cosmological GRBs and hypernovae are generally taken to be ~ 1 bethe and a few times 10 bethes, respectively. However, the GRB involves, in order that it can begin, the clearing of the matter in the way of the jet in order that it can leave the star. This energy is roughly the mass still contained within the beaming angle of the jet times the square of the velocity with which it is displaced. This is about $1M_\odot$ (including both poles) times $(1\% - 10\%)c^2$ [89]. MacFadyen estimates this to be a few $\times 10^{51}$ ergs. However, the$(1\% - 10\%)M_\odot c^2$ is $\sim (2 \times 10^{52} - 2 \times 10^{53})$ ergs. We compare this with the thermal energy of $\sim 3 \times 10^{53}$ ergs in the hypernova 1998bw. The mechanism for heat production is viscous heating, from the field lines frozen in the rotating disk, different from ordinary supernovae. In the latter, the kinetic and heat energy are roughly equal, from energy equipartition. We assume this to be true also in the GRB-Hypernova case. Otherwise we cannot make sense out of our energies. In other words, we believe that GRB 980425/SN 1998bw had essentially the same rotational energy as cosmological GRBs, the GRB, however, being nearly smothered by the work from the ram pressure.

This means that the kinetic energy is roughly equal to that of the hypernova, ~ 30 bethes in the case of GRB 980425. The GRB is, however, a "smothered" one [89], which means that the ram pressure to free the jet uses up nearly all of the kinetic energy.

In cosmological GRBs the typical GRB energy is ~ 1 bethe, resulting from a near cancellation of work from ram pressure work and other kinetic energy.

If such a near cancellation is to be engineered, the helium stars in which the jet is produced cannot be very different in properties. However, the cosmological GRBs are only a tiny population compared with subluminal ones, which are probably formed with helium stars that do not show this near cancellation of kinetic energies with ram pressure work.

To summarize, we are saying that there is a large "invisible energy" which goes into work done by the ram pressure, so that our energies calculated in the Blandford-Znajek mechanism look much larger than those found empirically in the observations, although it is generally realized that extra energy must be furnished to get photons out in the GRB, by saying that there is an efficiency, often taken to be $\sim 10\%$.

Appendix E: A New Constraint for GRB Progenitor Mass.

We wish to point out that there is observational support for our choice of the LMC as a site for the long gamma-ray bursts. Larsson et al. [79] provide a new constraint for gamma-ray burst progenitor mass. They show that long-duration gamma-ray bursts (L-GRBs) are much more concentrated on their host galaxy light than core collapse supernova explosions. From this they say *"Assuming core collapse supernova arise from stars with main-sequence masses* $> 8M_{\odot}$*, GRBs are likely to arise from stars with initial masses* $> 20M_{\odot}$*. This difference can naturally be explained by the requirement that stars which create a L-GRB must also create a black hole."*

As a template and close by natural analogue of starburst galaxies they use NGC 4038/39. They say, however, *the luminosity function and surface density of clusters on NGC 4038/39 is comparable to that seen in other local star-forming galaxies of varying morphology in which they include the LMC.*