THE THERMAL EVOLUTION FOLLOWING A SUPERBURST ON AN ACCRETING NEUTRON STAR

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ABSTRACT

Superbursts are very energetic type I X-ray bursts discovered in recent years by long-term monitoring of X-ray bursters and are believed to be due to unstable ignition of carbon in the deep ocean of the neutron star. In this Letter, we follow the thermal evolution of the surface layers as they cool following the burst. The resulting light curves agree very well with observations for layer masses in the range \(10^{17} - 10^{18}\) g expected from ignition calculations and for an energy release \(\gtrsim 10^{47}\) ergs cm\(^{-2}\) during the flash. We show that at late times the cooling flux from the layer decays as a power law \(F \propto r^{-\alpha}\), giving timescales for quenching of normal type I bursts of weeks, in good agreement with observational limits. We show that simultaneous modeling of superburst light curves and quenching times promises to constrain both the thickness of the fuel layer and the energy deposited.

Subject headings: accretion, accretion disks — stars: neutron — X-rays: bursts

1. INTRODUCTION

Type I X-ray bursts from accreting neutron stars in low-mass X-ray binaries involve unstable thermonuclear burning of accreted hydrogen (H) and helium (He; Lewin, van Paradijs, & Taam 1995). In the last few years, long-term monitoring of X-ray bursters by BeppoSAX and the Rossi X-Ray Timing Explorer has revealed a new class of very energetic type I X-ray bursts, now known as “superbursts” (see Strohmayer & Bildsten 2004 for reviews). The \(10^{47}\) ergs energies and several hour durations of superbursts are \(100-1000\) times greater than usual type I bursts. In addition, they are rare: so far, eight have been seen from seven sources, with recurrence times not well constrained but estimated as \(\sim 1\) yr (Kuulkers 2002; in ’t Zand et al. 2003; Wijnands 2001), instead of hours to days for usual type I bursts.

The current picture is that superbursts are due to unstable ignition of carbon at densities \(\rho \sim 10^{8} - 10^{9}\) g cm\(^{-3}\). Hydrogen and helium burn at \(\rho \sim 10^{5} - 10^{6}\) g cm\(^{-3}\) via the rp-process (Wallace & Woosley 1981), producing chiefly heavy elements beyond the iron group (including nuclei as massive as \(A = 104\); Schatz et al. 2001) but with some residual carbon (mass fraction \(X_c \sim 0.01-0.1\); Schatz et al. 2003b). Cumming & Bildsten (2001, hereafter CB01) showed that this small amount of carbon can ignite unstably once the mass of the ash layer reaches \(10^{3} - 10^{4}\) g (see also Strohmayer & Brown 2002). This fits well with observed superburst energies for \(X_c \approx 0.1\) and an energy release from the nuclear burning of 1 MeV per nucleon. The low thermal conductivity of the rp-process ashes gives a large temperature gradient and ignition at the required mass (CB01). The heavy nuclei may also photodisintegrate to iron group during the flash, enhancing the nuclear energy release (Schatz, Bildsten, & Cumming 2003a). Therefore superbursts offer an opportunity to study the rp-process ashes.

Previous authors used one-zone models to estimate the time dependence of the flash (CB01; Strohmayer & Brown 2002). In this Letter, we present the first multi-zone models of the cooling phase of superbursts. Unlike normal type I bursts, the time to burn the fuel is much less than the convective turnover time. We therefore assume that the fuel burns locally and instantaneously in place, without significant vertical mixing. We do not calculate ignition conditions but rather treat the amount of energy deposited and the thickness of the fuel layer as free parameters.2 In § 2, we describe our calculations of the subsequent thermal evolution of the layer and present a simple analytic model that helps to understand the numerical results. At late times, the flux evolves as a power law in time rather than the exponential decay found by CB01 for a one-zone model. In § 3, we use the long-term flux evolution of the layer to predict the timescale of quenching of type I bursts after the superburst and compare to observations.

2. TIME EVOLUTION OF THE SUPERBURST

After the fuel burns, the cooling of the layer is described by the entropy equation

\[
\frac{\partial s}{\partial t} = -\frac{1}{\rho} \frac{\partial F}{\partial r},
\]

where the heat flux is \(F = -K(\partial T/\partial r)\) and \(s\) is the neutrino energy loss rate. The layer remains in hydrostatic balance, in which case a useful independent coordinate is the column depth \(y\) into the star (in units of g cm\(^{-2}\)), where \(dy = -\rho dr\), giving a pressure \(P = gy\). The surface gravity is \(g = (GM/R^2)(1 + z)\), where \(1 + z = (1 - 2GM/Rc^2)^{-1/2}\) is the gravitational redshift factor. In this Letter, we assume \(M = 1.4 M_\odot\) and \(R = 10\) km, giving \(z = 0.31\) and \(g_{14} = g/10^{14}\) cm s\(^{-2}\) = 2.45.

To find the temperature profile just after the fuel burns, we deposit an energy \(E_{\text{fuel}} = Ec_r 10^{47}\) ergs g\(^{-1}\) throughout the layer. Since carbon burning to iron gives \(\sim 10^{46}\) ergs g\(^{-1}\), we expect \(E_{\text{fuel}} = 1\) to correspond to \(X_c \approx 0.05-0.1\), depending on how much energy is contributed by photodisintegration (Schatz et al. 2003a). At each depth, we calculate the temperature of the layer, \(T_y\), from \(\int_0^y c_p dT = E_{\text{fuel}}\), where \(T_y\) is the initial temperature. Following CB01, a simple analytic estimate of \(T_y\) is as follows. The electrons are degenerate and relativistic for \(\rho \gtrsim 10^7\) g cm\(^{-3}\), giving \(n_e = (1.5 \times 10^8\) g cm\(^{-3}\)) \(P_{56}^{1/4}\). Fermi energy \(E_F = (2.7\) MeV\) \(P_{56}^{1/4}\), and pressure scale height \(H = y/\rho = (67\) m\) \(P_{56}^{1/4}\gamma_{8/4}\), where \(P_{56} = P_{56}^{1/4}\) ergs cm\(^{-3}\). The heat capacity \(c_p\) is determined mainly by the electrons, \(c_p \approx\)

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2 This is a similar approach to Eichler & Cheng (1989), who studied the thermal response of a neutron star to energy deposition at different depths. However, the transient events that they consider are less energetic than superbursts.
\[ \pi^2 \left( Y_n k_B/m_p \right) (k_n T E_n) = (2.6 \times 10^7 \text{ ergs g}^{-1} \text{ K}^{-1}) T_B P_{10}^{-1/4}. \]

Integrating and assuming \( T_s \ll T_f \), we find the convectively stable temperature profile \( T_f = (3.6 \times 10^9 \text{ K}) E_17^{-1/5} y_{10}^{-1/8} \), insensitive to depth and depending mainly on \( E_{17} \).

Our thermal evolution code uses the method of lines, in which the right-hand side of equation (1) is differenced over a spatial grid and the resulting set of ordinary differential equations integrated using a stiff integrator. We choose a uniform grid in depth and the resulting set of ordinary differential equations is integrated using a stiff integrator. We choose a uniform grid in depth and the resulting set of ordinary differential equations.

\[ \frac{\partial T}{\partial t} = \frac{1}{\sigma T} \frac{\partial}{\partial x} \left( Y_k \frac{1}{\sigma} (y k_B/m_p) \frac{D T}{D x} \right) = \frac{1}{\sigma T} \frac{\partial}{\partial x} \left( \frac{D T}{D x} \right). \]

\[ T(x, t) = \frac{1}{\sqrt{\pi Dt}} \sinh \left( \frac{ax}{2Dt} \right) \exp \left( -\frac{x^2 + a^2}{4Dt} \right). \]

A similar problem is ohmic decay of crustal magnetic fields, where power-law decay is also expected (Sang & Channugam 1987; Upin, Channugam, & Sang 1994). A simple way to obtain this result is to apply the method of images to the Green’s function for an unbounded domain \( T(x, t) \sim t^{-1/2} \exp (\sqrt{2} \sqrt{4Dt}) \).

Eichler & Cheng (1989) derive a similar result for a power-law dependence of conductivity on depth, which also shows self-similar behavior at late times (see Lyubarsky, Eichler, & Thompson 2002 for a recent application to cooling of SGR 1900+14 after an outburst).
and the surface flux is $F \propto (\partial T/\partial x)_{x=a} \propto r^{-3/2} \exp (-\tau/\theta)$, where $\tau$ is the thermal time at the initial heating depth $\tau = 4a^3/D$. For an initial “top hat” temperature profile $T(x < a) = 1$, $T(x > a) = 0$, the surface flux is $F \propto (\tau/\theta)^{1/2} \exp (-\tau/\theta)$. For $\tau > \theta$, before the cooling wave reaches the base of the layer, $F \propto r^{-1/2}$; for $\tau > \theta$, the solution is independent of the initial temperature profile, and $F \propto r^{-3/2}$.

The numerical results show a similar behavior, although with different power-law indices. The relevant timescale in this case is $t_{\text{cool}} = H^2/D$, where $D = K l \rho c_p$. The electron conductivity is $K = \pi n_e kT/\tau_{\text{rel}}$, where $\tau_{\text{rel}} = E_p/c^2$ and $\tau_{\text{rel}}$ is the electron collision frequency. When electrons dominate the heat capacity, the thermal diffusivity takes the particularly simple temperature-independent form $D = c^2/\tau_{\text{rel}}$. For electron-ion collisions, $\tau_{\text{rel}} = (9.3 \times 10^{16} \text{ s}^{-1}) P_{10}^{1/4} (Z^2/|A|) A_e/|Y_e|$, (e.g., see Appendix of Schatz et al. 1999), giving

$$t_{\text{cool}} = (3.8 \text{ hr}) Y_{12}^{1/4} \frac{Y_e (Z^2/|A|) A_e}{6} \left( \frac{B_{14}}{2.45} \right)^{-5/4}$$

(3)

(see also eq. [10] of CB01), where we insert the appropriate numbers for $^{56}$Fe composition. The simple top hat solution for constant conductivity motivates a fit to the numerical solutions,

$$F_{25} = 0.2 t_{\text{hr}}^{1/4} E_{17}^{3/4} [1 - \exp (-0.63 t_{\text{cool}}^{3/4} E_{17}^{3/4} t_{\text{hr}}^{-1/3})]$$

(4)

where $t_{\text{hr}} = t/1 \text{ hr}$. For $t > t_{\text{cool}}$, $F_{25} = 0.13 E_{17}^{1/2} (\text{hr} t_{\text{cool}})^{-4/3}$. The transition from $F \propto r^{-3/2}$ to $F \propto r^{-4/3}$ occurs when $t_{\text{cool}} \approx E_{17}^{-1} r_{12}^{1/2}$.

Equation (4) fits the numerical results to better than a factor of 2 for models with substantial neutrino emission. As emphasized by Strohmayer & Brown (2002), neutrino cooling is important for large carbon fractions: it depresses the flux at $t \approx 5-10$ hr for the models with $E_{17} = 2$ and $3$, $y = 10^{15} \text{ g cm}^{-2}$ in Figure 2. Whenever neutrinos dominate the cooling, the peak temperature is large enough that emission is by pair annihilation. A good fit to the neutrino energy loss rate is $\epsilon_p \approx (10^4 \text{ ergs g}^{-1} \text{ s}^{-1}) T_{12}^{12} Y_1^{1/2}$, giving a cooling time $t_c = \epsilon_p T_1 e_p = (2.5 \times 10^{12} \text{ s}) Y_1^{-1/2} T_{12}^{-10}$. Inserting the peak temperature from equation (2) gives $t_c \approx (300 \text{ hr}) E_{17}^{3/5}$. Neutrinos dominate when $t_c < t_{\text{cool}}$ or when $E_{17} > 2.3 y_{12}^{-3/20}$.

3. COMPARISON TO OBSERVATIONS

The cooling curves in Figure 2 compare well with observed light curves, including a rapid initial decay on hour timescales, followed by an extended tail of emission (as observed following some superbursts, e.g., KS 1731–260: Kuulkers et al. 2002; Ser X-1: Cornelisse et al. 2002). We will present a detailed comparison with the observed superburst light curves in a future paper. The initial decay from the peak depends mostly on $E_{17}$, and so it should be possible to constrain the amount of fuel consumed in the superburst. Our models do not resolve the peak itself, since this depends on the details of how the burning propagates out to the surface; however, for $E_{17} \approx 2–3$, the flux exceeds the Eddington flux, $F_{\text{Edd}} = 3 \times 10^{25} \text{ ergs cm}^{-2} \text{ s}^{-1}$ $(1 + X)$, where $X$ is the H fraction, for timescales of minutes. Superburst peak luminosities are generally less than the Eddington luminosity (Kuulkers 2004), implying $E_{17} \approx 2$. The one exception is the superburst from 4U 1820–30, which showed dramatic photospheric radius expansion lasting for several minutes (Strohmayer & Brown 2002). This is consistent with the proposal that this source, which accretes and burns He-rich material, produces large quantities of carbon (Strohmayer & Brown 2002; Cumming 2003). The transition to the late-time power law occurs after $t \approx (4 \text{ hr}) E_{17}^{-1} y_{12}^{1/4}$ (eq. [5]), which corresponds to $F_{25} \approx 0.13 E_{25}$. It may therefore be possible to measure the power-law decay using superburst tails, although this depends on being able to subtract out the underlying accretion luminosity, $F_{\text{accretion}} \approx 0.1 (M/0.1 M_{\odot})$, in a reliable way.

Another way to probe the late-time cooling is to use the remarkable observation that type I bursts disappear (are “quenched”) for $t_{\text{quench}} \approx 7 \text{ weeks}$ following the superburst (e.g., Kuulkers 2004). CB01 proposed that the cooling flux from the superburst temporarily stabilizes the H/He burning. An estimate of the critical stabilizing flux, $F_{\text{crit}}$, is as follows. When the condition for temperature fluctuations to grow and unstable He ignition to occur is $\rho c_s \approx \eta c_\epsilon$ (Fushiki & Lamb 1987), where $\epsilon_\alpha$ is the triple alpha $(3\alpha)$ energy production rate, $\epsilon_\text{cool}$ is a local approximation to the cooling rate, and $\rho$ and $c_\epsilon$ are the respective temperature sensitivities. For a large flux from below, the He burns stably before reaching this ignition condition, at a depth where the time to accumulate the layer equals the He-burning time, $y/\rho = Y_{\alpha} \rho c_\epsilon$, where $Y_\alpha$ is the mass fraction, $\rho$ is the local accretion rate per unit area, and $Q_{\alpha} = 5.8 \times 10^{17} \text{ ergs g}^{-1} = 0.606 \text{ MeV}$ nucleon$^{-1}$ is the $3\alpha$ energy release. At the transition from unstable to stable burning, both criteria are satisfied at the base of the H/He layer. Using the first condition to eliminate $\epsilon_\alpha$ from the second, and writing $\epsilon_\text{cool} \approx F/\rho c_\epsilon$, gives $F = \rho \epsilon_\text{cool} \rho c_\epsilon = 6.2 \times 10^{22} \text{ ergs cm}^{-2} \text{ s}^{-1} (\rho \epsilon_\text{cool}) (Y_{\alpha} / 0.5)$. Some of this flux is provided by hot CNO burning of accreted H. $F_{\text{crit}} \approx \rho c_\alpha = (5.8 \times 10^{17} \text{ ergs cm}^{-2} \text{ s}^{-1}) y_{12} (Z/0.01)$ (Cumming & Bildsten 2000; Z is the metallicity); the remainder is $F_{\text{crit}} = F - F_{\text{crit}}$. This estimate agrees well with a more detailed calculation using the ignition models of Cumming & Bildsten (2000), in which we find $F_{\text{crit}} \approx Q_{\alpha} = 0.7 \text{ MeV per accreted nucleon, almost independent of } M$. Therefore,

$$F_{\text{crit}} \approx 6 (\rho \epsilon_\text{cool})$$

(5)

(see also Paczyński 1983; Bildsten 1995). Equation (5) in the limit $t > t_{\text{cool}}$ gives $t_{\text{quench}} = 38 t_{\text{cool}}^{3/4} E_{17}^{3/8} = (6 \text{ days}) Y_{12}^{3/4} F_{\text{crit}}^{3/4} E_{17}^{3/8}$, (6)

which gives $t_{\text{quench}}$ in terms of the thickness of the layer and energy release.

Figure 4 compares the predicted and observed quenching times. The observations of $t_{\text{quench}}$ and accretion rates (used to find $F_{\text{crit}}$ from eq. [5]) are taken from Kuulkers (2004; except for 4U 1636–53, which has a revised upper limit of 23 days; Kuulkers et al. 2004). The observations are upper or lower limits only; nonetheless, the general agreement is very good and supports the quenching picture suggested by CB01. There is much to learn from a careful comparison of superburst light curves and the corresponding quenching times, separately constraining both $E_{17}$ and $y_{12}$.

4. SUMMARY AND CONCLUSIONS

We have presented the first multizone models of the cooling phase of superbursts. The flux decay is not exponential but power law (eq. [5]). For $t < t_{\text{cool}}$, where $t_{\text{cool}}$ is the cooling time at the base of the layer, the flux depends mostly on the energy release $E_{17}$, and is insensitive to depth: the inward traveling cooling wave does not yet know that the layer has a finite thickness. For $t > t_{\text{cool}}$, the flux decays as a power law $F \propto r^{-4/3}$, independent of the initial temperature profile. The power-
Fig. 4.—Observed and predicted quenching timescales of normal type I bursts following a superburst. We show the predicted quenching time for $y_p = 10^3$ (solid lines) and $10^{13}$ g cm$^{-2}$ (dotted lines) as a function of both the critical flux needed to stabilize H/He burning and the accretion rate (given in terms of $F_{\text{crit}}$ by eq. [5]). For each value of $y_p$, we show curves for (bottom to top) $F_{\text{crit}} = 1$, 2, and 3. The observed upper or lower limits on $t_{\text{quench}}$ and $E_{\text{p}}$ are taken from Table 1 of Kuulkers (2004), with an updated $\dot{M}$ value for 4U 1636−53 (E. Kuulkers 2003, private communication).

law decay at late times gives predicted type I burst quenching times of weeks (eq. [6]), consistent with observational limits. Future comparisons of both superburst light curves and quenching times with observations will constrain both the thickness of the fuel layer and the energy deposited, particularly when combined with models of normal type I bursts from the same source (Cumming 2003, 2004).

There is still much to be done in terms of theory. Perhaps the most important issues are the physics of the rise (which sets the initial condition for our simulations) and production of the fuel. Important clues to the first are the observed precursors to superbursts, which may be normal type I bursts ignited by the superburst. Recent progress has been made on the second, with indications from both theory (Schatz et al. 2003b; Woosley et al. 2004) and observations (in ’t Zand et al. 2003) that stable burning may be required to produce enough carbon to power superbursts. A self-consistent model of H/He burning, followed by accumulation and ignition of the ashes, may require a better understanding of the transition from unstable to stable burning observed in normal type I bursting (e.g., Cornelisse et al. 2003).

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