CARBON FLASHES IN THE HEAVY ELEMENT OCEAN ON ACCRETING NEUTRON STARS

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Submitted to The Astrophysical Journal

ABSTRACT

We show that the burning of a small mass fraction $X_{12}$ of $^{12}$C in a neutron star ocean is thermally unstable at low accumulated masses when the ocean contains heavy ashes from the hydrogen burning rapid proton (rp) process. The key to early unstable ignition is the decreased thermal conductivity of a heavy element ocean. The instability requires accretion rates, $\dot{M}$, in excess of one-tenth the Eddington limit when $X_{12} < 0.1$. Lower $\dot{M}$’s will stably burn a small mass fraction of $^{12}$C. The unstable flashes release $\sim 10^{42} - 10^{43}$ergs over hours to days and are likely the cause of the recently discovered large Type I bursts (so-called “superbursts”) from six Galactic low-mass X-ray binaries. In addition to explaining the energetics, recurrence times and durations of the superbursts, these mixed $^{12}$C flashes also have an $M$ dependence of unstable burning similar to that observed. Though the instability is present at accretion rates $\approx M_{\text{Edd}}$, the flashes provide less of a contrast with the accretion luminosity there, thus explaining why most detections are made at $M \approx (0.1 - 0.3)M_{\text{Edd}}$. Future comparisons of time dependent theoretical calculations with observations will provide new insights on the rp-process.

Subject headings: accretion, accretion disks – nuclear reactions, nucleosynthesis, abundances – stars: neutron – X-rays: bursts

1. INTRODUCTION

The monitoring of galactic sources by the Wide-Field Camera on BeppoSAX and the All-Sky Monitor on the Rossi X-Ray Timing Explorer have detected large Type I bursts (hereafter referred to as “superbursts”) from six accreting neutron stars (Cornelisse et al. 2000; Strohmaier 2000; Heise et al. 2000; Wijnands 2001; Kuulkers 2001). These bursts share a number of characteristics: burst energies $\approx 10^{42}$ erg (roughly 1000 times larger than a normal Type I burst); durations of a few hours and accretion rates $\approx (0.1 - 0.3)M_{\text{Edd}}$ (Wijnands 2001). The superburst recurrence time is not well known, though one recurrent burst has been seen 4.7 years later (Wijnands 2001).

The light curves are similar to a Type I burst in spectral evolution, though the immediate decay timescale is about two hours; 1000 times longer than a typical Type I X-ray burst. The superbursts appear to be thermonuclear flashes from fuel at much larger depths than a typical Type I burst. For example, if the nuclear energy release is 1 MeV per accreted nucleon ($\dot{E}_{\text{nuc}} = 10^{18}$erg g$^{-1}$), the accumulated mass would be $10^{24}$ g, implying a recurrence time of $\approx 4$ months for an accretion rate of $M \approx 0.1M_{\text{Edd}} \approx 10^{17}$ g s$^{-1}$, much longer than a typical Type I burst.

Flashes from pure $^{12}$C layers have been discussed previously (Woosley & Taam 1976; Taam & Picklum 1978; Brown & Bildsten 1998) and might well apply to the pure helium accretor 4U 1820-30 (Strohmaier 2000). However, there are two difficulties with such a model for the superbursts from hydrogen/helium accretors: (1) the recurrence times are very long, and the burst energies correspondingly large, unless a large heat flux from the crust of the star is present deep in the ocean; and, (2) our current understanding of the ashes from hydrogen/helium burning does not point to accumulation of pure $^{12}$C. Instead, Schatz et al. (1999, 2001) have shown that only a small amount of $^{12}$C remains after all of the hydrogen and helium has burned via the rapid proton (rp) process (Wallace & Woosley 1981); either from steady state burning or unstable type I X-ray burning. Even though a small fraction by mass (typically $X_{12} \approx 0.05 - 0.1$), we show here that it is enough to trigger a thermonuclear runaway with energy comparable to the superbursts. Most important, however, is the role played by the heavy ashes from the rp process. These reduce the thermal conductivity enough to force a large temperature gradient in the ocean. This ignites the $^{12}$C at much lower masses than previously expected and reduces the energy and recurrence times to levels consistent with the superbursts. In addition, the conduction cooling times from these flashes are consistent with the long decay times observed.

2. CARBON IGNITION IN A HEAVY ELEMENT OCEAN

We presume that the ashes of hydrogen/helium burning in the upper atmosphere primarily consist (by mass) of a single dominant nucleus of mass $A_{\text{np}}$ and charge $Z_{\text{e}}$ and a small amount of $^{12}$C with mass fraction $X_{12}$. Schatz et al. (2001) recently showed that the onset of a closed cycle (called the SnSbTe cycle) in the rp-process naturally stops the increasing pressures, the equation of state is that of degenerate, relativistic electrons with Fermi energy $E_F = 1.9$ MeV$(2Z/A)^{1/3}/\rho_8^{1/3}$, where $\rho_8 = \rho/10^8$ g cm$^{-3}$ and $E_F \approx 4.3$ MeV.
Though they exert little pressure, the nuclei set the thermal conductivity via electron-ion scattering. For a classical one component plasma with ion separation, a, defined by \(a^3 = 3/4\pi n_i\), where \(n_i = \rho/A m_p\), the importance of Coulomb physics for the ions is measured by

\[
\Gamma = \frac{(Ze)^2}{akT} = 434 \frac{\rho^{1/3}}{T_8} \left(\frac{Z}{44}\right)^2 \left(\frac{104}{A}\right)^{1/3},
\]

where \(T_8 = T/10^8 K\). For this initial calculation we use the microphysics of the liquid state; \(\Gamma < 173\) (Farouki & Hanaguchi 1993 and references therein), even when, in some instances (mostly when \(m < 0.3 m_{\text{Edd}}\)) \(\Gamma > 173\).

Electron conduction dominates the heat transport in the liquid ocean. The heat flux is given by \(F = -KdT/dz\), where \(K\) is the thermal conductivity (see Yakovlev & Urpin 1980; Itoh et al. 1983, Potekhin et al. 1999). For \(X_{12} \ll 1\), we find

\[
K \approx 1.1 \times 10^{18} \text{ergs cm}^{-1}\text{s}^{-1}\text{K}^{-1},
\]

where we assume the electrons are relativistic, and take the Coulomb logarithm to be unity (a good approximation in the high densities, depth. This is critical, as it allows for ignition at lower column densities, as is shown in Figure 1 for two cases. The lower dotted line is for a pure \(^{12}\)C ocean and matches that in Brown & Bildsten (1998). The upper dotted line shows the ignition curve for a \(^{12}\)C–\(^{104}\)Ru mixture with \(X_{12} = 0.1\). Because of the lower abundance, a larger column density is needed to ignite \(^{12}\)C in a \(^{104}\)Ru ocean for a given temperature. However, ignition occurs at smaller column depths in this case, because \(^{12}\)C–\(^{104}\)Ru ocean is hotter.

For the pure \(^{12}\)C models of Figure 1, ignition occurs at \(y = 2.1 \times 10^{12} \text{ g cm}^{-2}\) (\(y = 2.4 \times 10^{13} \text{ g cm}^{-2}\) for \(m = m_{\text{Edd}}\) (\(m = 0.3 m_{\text{Edd}}\))). Writing the burst energy \(E_{\text{burst}} = 4\pi R^2 \nu \epsilon_{\text{nuc}} X_{12}\), \(E_{\text{burst}} = 7 \times 10^{42}\) \(\text{ergs per y} 12\text{C}/(R/10\text{ km})^2\), and recurrence time \(t_{\text{rec}} = y/m = 0.36 \text{ yr} X_{12}(m_{\text{Edd}}/m)\), we find that pure \(^{12}\)C ignitions have recurrence times 0.8 years (29 years) and burst energies 1.5 \(\times 10^{43}\) \(\text{ergs} (1.7 \times 10^{44}\) ergs). These burst energies exceed the energy of the superbursts by at least a factor of ten. For the \(^{12}\)C–\(^{104}\)Ru models, ignition occurs at \(y = 1.0 \times 10^{11} \text{ g cm}^{-2}\) (\(y = 1.8 \times 10^{12} \text{ g cm}^{-2}\)) for \(m = m_{\text{Edd}}\) (\(m = 0.3 m_{\text{Edd}}\)), with recurrence time 13 days (2.2 years), and burst energy 6.9 \(\times 10^{40}\) \(\text{ergs} (1.3 \times 10^{42}\) ergs).

3. ACCRETION RATE DEPENDENCES OF THE CARBON FLASHES

Not only must the thermal settling solution reach the ignition curve, but the \(^{12}\)C must also have survived to that depth. Thus, in order to obtain a flash, the \(^{12}\)C lifetime to the fusion reaction, \(t_{\text{rec}} \approx E_{\text{nuc}} X_{12}/\nu_{\text{nuc}}\), must exceed the accumulation time, \(y/m\), at the ignition point. Combining this condition with equation (5) implies that \(m\) must exceed the critical value

\[
\dot{m}_c = \frac{2\nu K T}{E_{\text{nuc}} X_{12} \nu^2}.
\]

When we evaluate this using \(K\) and the \(\rho, y\) relations for degenerate relativistic electrons, the density cancels out, giving

\[
\dot{m}_c \approx \frac{0.1}{X_{12}} \left(\frac{T_8}{6}\right)^2 \left(\frac{26}{\nu}\right) \left(\frac{g_{14}}{2}\right) \left(A_{104}\right) \left(\frac{44}{Z}\right)^2.
\]

This relation is robust and shows that large accretion rates are required for unstable \(^{12}\)C ignition. In order to eliminate the temperature, we neglect the outer temperature in equation (4) and substitute for the temperature at the base in equation (7). We then find that the explicit accretion rate dependences cancel and the requirement for an instability is just

\[
X_{12} > \left(\frac{4Q}{\nu \epsilon_{\text{nuc}}} \ln \left(\frac{P_2}{P_1}\right)\right)_{\text{approx}} \approx 0.2 Q_{17} \left(\frac{26}{\nu}\right) \ln \left(\frac{P_2}{P_1}\right)/8.
\]

At smaller \(X_{12}\), the \(^{12}\)C depletes and burns steadily before reaching ignition. Figure 2 shows the \(^{12}\)C mass fraction required to
achieve an explosion as a function of \( \dot{m} \) and \( Q_{17} \). The accretion rate dependence is mostly from the ignition pressure dependence on \( \dot{m} \). The top two panels in figure 3 show the burst energies and recurrence times for unstable models. The curves are for \( X_{12} = 0.1 \) and (top to bottom), \( Q_{17} = 0.5, 1.0 \) and 2.0.

4. TIME DEPENDENT BEHAVIOR OF THE FLASH

Having derived the conditions for an unstable ignition of a small mass fraction of \( ^{12}\text{C} \), we now calculate the time dependent nature of the flash and estimate the time it takes for the energy to escape from the burning layer. The “one-zone” version of the time dependent heat equation is

\[
C_V \frac{dT}{dt} = \epsilon_{\text{nucl}} - \epsilon_{\text{cool}},
\]

in which we write the specific heat at constant volume \( C_V \) rather than constant pressure (an excellent approximation since \( E_F \gg k_BT \) throughout the flash).

Once ignited, the \(^{12}\text{C} \) burning is extremely rapid and there is no time for cooling to occur while the \(^{12}\text{C} \) burns. The fluid then evolves to a final temperature set by the total energy released

\[
\int_{T_1}^{T_f} C_V dT = \int \epsilon_{\text{nucl}} dt = X_{12} L_{\text{nucl}}.
\]

Using \( C_V \) for the electrons and taking \( T_i \gg T_f \) gives

\[
T_f = 1.92 \times 10^9 \text{K} \left( \frac{E_F}{\text{MeV}} \right)^{1/2} \left( \frac{X_{12}}{0.1} \right)^{1/2} \left( \frac{A}{2.36Z} \right)^{1/2}.
\]

Thus, \( X_{12} > 0.01 \) is adequate to substantially change the temperature from that at ignition and trigger a thermal instability. If \( X_{12} > 0.2-0.3 \), the peak temperatures would exceed \( 5 \times 10^9 \text{K} \) and neutrino cooling would reduce the observed energetics of the event.

After all the \(^{12}\text{C} \) has burned and the fluid has reached \( T_i \), the layer cools. Integrating equation (9) starting from \( T_i \) with only the \( \epsilon_{\text{cool}} \) term, we find
\[
d\ln T = -dt/2t_{\text{cool}},
\]

where
\[
t_{\text{cool}} = 4.15 \text{ hrs} \left( \frac{\dot{m}}{8\times 10^{-12}} \right)^4 \left( \frac{Z}{44} \right)^4 \left( \frac{104}{A} \right)^3.
\]

depends only on the depth in the star and the composition, and is temperature independent. The luminosity during the conductive cooling phase is thus set by the temperature at the base of the layer; \( L = 4\pi R^2 \epsilon_{\text{cool}} \propto T^2 \propto T_f^2 \exp(-t/t_{\text{cool}}) \). Thus, the observer should see a timescale of luminosity decay equal to equation (11), as shown in Figure 3. If we use \( T_i \) from equation (10) for an upper limit, we find the maximum conductive cooling luminosity

\[
\frac{L_c}{L_{\text{Edd}}} = 0.77 \left( \frac{X_{12}}{0.1} \right)^{1/3} \left( \frac{A}{104} \right)^{5/3} \left( \frac{44}{Z} \right)^{8/3}.
\]

5. CONCLUSIONS AND FUTURE WORK

We have shown that the ignition of a small mass fraction of \(^{12}\text{C} \) in a neutron star ocean consisting mostly of heavier elements is thermally unstable and leads to flashes whose energy, recurrence times and durations are similar to the superbursts. The energy from these flashes takes a long time to escape the star, possibly explaining the persistent “offset” in the flux nearly a day after the superburst in 4U 1735-44 (see Figure 1 of Cornellise et al. 2000). We also predict flashes at high local accretion rates which include the Z sources and accreting X-ray pulsars. In the pulsar case, some of the difficulties associated with explaining flashes from LMC X-4 via pure \(^{12}\text{C} \) burning (see Brown & Bildsten 1998) might well be alleviated in the mixed case we present here.

Cornellise et al. (2000) reported an absence of regular type I bursts for seven days after the superburst. This might be due to thermal stabilization of the hydrogen/helium burning layers by the large heat flux from the cooling ashes of the \(^{12}\text{C} \) burning. Paczynski (1983) and Bildsten (1995) have shown that luminosities in excess of the helium burning flux (or \( L > L_{\text{accr}}/100 \)) will stabilize the burning, as then the thermal state of the burning layer is independent of the local burning rates. This appears most likely at low accretion rates, where the conductive cooling is slowest, thus halting Type I bursts for \( \approx 5t_{\text{cool}} \sim \) days after the superburst.

There is still much to be done, including a time dependent study of the onset of the instability. Our initial integrations of the time-dependent code from Bildsten (1995) demonstrate a transition from steady-state burning at low \( \dot{m} \) to unstable flashes at higher rates. In future work, this code will be used to more accurately calculate the transition accretion rate, \( \dot{m}_c \), with realistic mixtures of \(^{12}\text{C} \) and rp-process ashes. Only then can we make the comparisons to observations that will allow us to constrain the rp-process ashes.

We thank Erik Kuulkers and Marten van Kerckwijk for conversations about the properties of the superbursts and Hendrik Schatz for insights on the rp-process. We also thank Deepto Chakrabarty and Bob Rutledge for comments on the manuscript. This research was supported by NASA via grant NAG 5-8658 and by the National Science Foundation under Grants PHY99-07949 and AY97-31632. L. B. is a Cottrell Scholar of the Research Corporation.
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**Fig. 1.** Thermal profiles and ignition curves in the deep ocean of an accreting neutron star. The ignition curves are shown as dotted lines for (bottom to top) $X_{12} = 1.0$ and 0.1. The long-dashed curve shows where $\Gamma = 173$ for $^{104}$Ru (the equivalent curve for $^{12}$C lies off the right of the plot). The temperature as a function of column depth for a $^{12}$C–$^{104}$Ru mixture with $X_{12} = 0.1$ is shown in the two upper curves, while the lower two curves are for a pure $^{12}$C ocean. We show two accretion rates, $\dot{m} = 0.3 \dot{m}_{Edd}$ (solid), and $\dot{m} = \dot{m}_{Edd}$ (dashed). We take $Q_{17} = 1$. These profiles end at the depth where $^{12}$C ignites, as defined by equation (5).
FIG. 2.— The minimum mass fraction of $^{12}$C, $X_{12}$, needed to reach unstable ignition before depletion, as a function of accretion rate for $Q_{17} = 0.5, 1.0$ and $2.0$ in a $^{104}$Ru ocean.
Fig. 3.—Burst energy, recurrence time, luminosity decay time (equation [11]) and $L_c$ (equation [12]) as a function of $\dot{m}$ for $X_{12} = 0.1$ in a $^{104}$Ru ocean for (top to bottom) $Q_{17} = 0.5, 1.0$ and 2.0. The curves begin at the lowest accretion rate at which the thermal instability occurs in the $^{104}$Ru ocean.