EXPOSING THE NUCLEAR BURNING ASHES OF RADII EXPANSION TYPE I X-RAY BURSTS

NEVIN N. WEINBERG, 1,2 LARS BILDSTEN, 2 AND HENDRIK SCHATZ 3

Received 2005 September 3; accepted 2005 November 7

ABSTRACT

We solve for the evolution of the vertical extent of the convective region of a neutron star atmosphere during a type I X-ray burst. The convective region is well mixed with ashes of nuclear burning, and its extent determines the rise time of the burst light curve. Using a full nuclear reaction network, we show that the maximum vertical extent of the convective region during photospheric radius expansion (RE) bursts can be sufficiently great that (1) some ashes of burning are ejected by the radiation-driven radius expansion (RE) phase and (2) some ashes of burning are exposed at the neutron star surface following the RE phase. We find that ashes with mass numbers in the range \( A \sim 30–60 \) are mixed in with the ejected material. We calculate the expected column density of ejected and surface ashes in hydrogen-like states and determine the equivalent widths of the resulting photoionization edges from both the wind and the neutron star surface. We find that these can exceed 100 eV and are potentially detectable. A detection would probe the nuclear burning processes and might enable a measurement of the gravitational redshift of the neutron star. In addition, we find that in bursts with pure helium burning layers, protons from \((\alpha, p)\) reactions cause a rapid onset of the \(^{12}\text{C}(\alpha, \gamma)^{13}\text{N}\) reaction sequence. The sequence bypasses the relatively slow \(^{12}\text{C}(\alpha, \gamma)^{16}\text{O}\) reaction and leads to a sudden surge in energy production that is directly observable as a rapid (~millisecond) increase in flux during burst rise.

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

1. INTRODUCTION

Type I X-ray bursts are produced by the unstable nuclear burning of freshly accreted hydrogen- and/or helium-rich material on the surface of a neutron star (NS) in a low-mass X-ray binary (for reviews, see Lewin et al. 1995; Bildsten 1998; Strohmayer & Bildsten 2006). The burst energies \((10^{39}–10^{40} \text{ergs})\), durations \((\sim 10–100 \text{s})\), and recurrence times (hours to days) depend strongly on the composition of the accreted matter and on the accretion rate, \(\dot{M}\), which can range from \(10^{-11}\) to \(10^{-8} \text{M}_\odot \text{yr}^{-1}\). The burst properties are also sensitive to the composition of the ashes of burning from preceding bursts (Taam 2002), as underscored by recent burst simulations that implement large nuclear reaction networks for the energy generation (Schatz et al. 2001; Brown et al. 2002; Woosley et al. 2004).

The sensitivity of the nuclear energy generation rate to temperature and density concentrates the burning during a burst to a thin layer at the base of the accreted atmosphere (e.g., Fujimoto et al. 1981; Fushiki & Lamb 1987). Since the thermal timescale there is \(\sim 1–10 \text{s}\) while the dynamical timescale is \(\sim 10^{-6} \text{s}\), the temperature gradient near the burning layer is nearly adiabatic, resulting in a region of highly efficient convection. The short mixing timescale ensures that the ashes of burning are well mixed throughout the convective region.

The vertical extent of the convective region evolves during the burst, as demonstrated in time-dependent numerical simulations (e.g., Joss 1978; Taam 1980; Ayal & Joss 1982; Woosley & Weaver 1984; Woosley et al. 2004). In this paper, we carry out a thorough survey of the dependence of the convection region’s extent on \(\dot{M}\), the composition of the accreted material, and the preburst thermal state of the atmosphere. We also show how the evolution of the convective extent influences the observed burst rise times.

We demonstrate that for photospheric radius expansion (RE) bursts, in which the super-Eddington luminosity drives a radiation-driven wind, the convective region extends out to sufficiently low pressures that ashes can be ejected by the wind. Depending on the burst parameters, the wind can eject as much as \(\sim 1\%\) of the accreted mass (Paczynski & Proszynski 1986; Joss & Melia 1987; Nobili et al. 1994), corresponding to approximately \(\dot{E}_{\text{nuc}}/\dot{E}_{\text{grav}}\), the ratio of nuclear energy to gravitational binding energy. Sugimoto et al. (1984) pointed out that such mass ejection could expose the underlying helium-rich material and result in an Eddington-limited peak flux for helium rather than hydrogen. We show that the ejected ashes may be directly detectable with spectroscopy during the RE phase and afterward, when the photosphere, laced with heavy-element ashes, settles to the NS surface.

The convective region evolves during a burst in the following sequence of events. As the base temperature rises and the nuclear energy generation increases during the early stages of a burst, the entropy in the convective region increases. Initially, only a negligible amount of thermal energy is lost to radiation diffusing into the overlying radiative region and the underlying crust. Since the timescale for radiative diffusion across the convective-radiative interface is longer than the burning timescale during these early stages, the convective region extends vertically outward to lower pressures (Fujimoto et al. 1981). Eventually, the burning rate becomes sufficiently slow that the nuclear energy flux is carried most efficiently by radiation rather than by convection. At that point the convective region recedes back to higher pressures.

Regardless of whether the convective region is expanding outward or receding downward, its extent at a given time is set by the radial location in the atmosphere at which the constant
entropy of the convective region equals the radially increasing entropy of the overlying radiative region. On the basis of this argument, Joss (1977) showed that the convective region never acquires a high enough entropy to reach the photosphere located at column depth \( \sim 1 \text{ g cm}^{-2} \). Nonetheless, the convective region can reach pressures \( \leq 1\% \) of that at the base of the accreted layer (Joss 1978; Taam 1981; Ayasli & Joss 1982; Hanawa & Sugimoto 1982; Hanawa & Fujimoto 1984). In their simulations, Woosley et al. (2004) found that in a burst with a pure helium burning layer, the peak flux exceeded the Eddington limit and the convective region extended beyond their numerical surface (i.e., the resolution limit of their grid located at a pressure \( \approx 0.3\% \) of that at the base).

In this paper, we solve the time-dependent entropy equation that describes the evolving thermal structure of the atmosphere and growth of the convective region. We calculate the minimum pressure reached by the convective region (or, equivalently, the minimum column depth \( y_{c, min} \)) for a range of burst parameters. We show that \( y_{c, min} \) is sensitive to the burst ignition conditions and that, in general, the larger the burst peak flux and the smaller the entropy of the preburst atmosphere, the smaller the value of \( y_{c, min} \). Thus, \( y_{c, min} \) decreases with decreasing \( M \). We find that, in systems where the accreted material is helium-rich, such as 4U 1820–30 (see Cumming 2003 and references therein), or in systems accreting solar abundances at low \( M (\leq 10^{-3} M_\odot \text{ yr}^{-1}) \), \( y_{c, min} \) \( \ll y_{\text{wind}} \) during RE bursts, where \( y_{\text{wind}} \) is the column depth above which mass is ejected by the radiation-driven wind. As a result, the wind ejects some of the nuclear burning ashes. Furthermore, some of the ashes that remain bound to the NS are exposed at the photosphere.

In §2 we describe our analytic prescription for calculating the evolution of the thermal structure of the radiative and convective regions during an X-ray burst. We explain why the evolution during the early, convective stages is sensitive to \( M \) and accreting composition. In §3 we describe the full nuclear reaction network we use to calculate the nucleosynthesis and nuclear energy generation. In §4 we examine how \( y_{c, min} \) depends on the burst ignition conditions and we explore those conditions most conducive to ash ejection and exposure.

Since the rise time of the burst light curve is determined by the thermal diffusion time at the top of the convective zone, we also evaluate the rise time dependence on burst parameters such as \( M \) and accreting composition. We show these results in §4.

In §5 we show the ash composition profile and discuss the observational consequences of heavy-element ash ejection and surface exposure. Depending on the ignition conditions, nuclei as heavy as \( A \sim 40 \) are ejected by the wind. We calculate the column density of ejected and surface ashes in hydrogen-like states and discuss the prospects for detecting the resulting photoionization edge features during high spectral resolution observations of RE bursts. Such features probe the nuclear burning and may constrain the NS equation of state. If bursts ignite in the ashes of previous bursts, as Woosley et al. (2004) suggest, even heavier elements are ejected. These may include some light \( p \)-nuclei (Schatz et al. 1998, 2001), whose origins are not understood. We conclude in §6 with a summary of our work and mention the possibility of ash ejection during superbursts.

2. EVOLUTION OF THE ATMOSPHERE DURING A BURST

In this section we consider the temporal evolution of the atmosphere’s thermal structure during a burst. We start in §2.1 with a general description of the atmosphere’s structure and explain the boundary condition at the convective-radiative interface.

The evolution of the location of the convective-radiative interface is described in §2.2.

2.1. Thermal Structure of the Atmosphere

The NS atmosphere maintains hydrostatic equilibrium throughout the burst so that pressure varies with height as \( dP/dr = -\rho g \), where \( \rho \) is the density and \( g \) is the surface gravity. We assume a NS mass \( M = 1.4 M_\odot \) and radius \( R = 10 \text{ km} \), giving \( g = (1+z)GM/R^2 = 2.43 \times 10^{14} \text{ cm s}^{-2} \), where the gravitational redshift \( z = (1 - 2GM/R^2)^{-1/2} - 1 = 0.31 \). Since the atmosphere is thin compared with the NS radius, \( g \) is effectively constant throughout the accreted layer. Hydrostatic balance therefore yields \( P = gy \), where the column depth \( y \), defined by \( dy = -\rho dr \), is a convenient parameterization of the vertical spatial coordinate. We determine the extent of the convective region over a broad range of burst parameters and thus consider only one-dimensional models in our calculations. We do not account for the effect of a spreading burning front during burst rise, nor for the influence of rotation on the convective structure, although these effects may be important (see Spitkovsky et al. 2002).

The entire atmosphere is radiative before helium burning at the base triggers the burst. The thermal profile is then described by the diffusion equation

\[
\frac{dT}{dy} = \frac{3\kappa F}{4acT^3},
\]

where \( F \) is the outward heat flux and the opacity \( \kappa \) has contributions from electron scattering and free-free absorption and is calculated using the approximation given by Schatz et al. (1999). The preburst flux \( F_0 = F_{\text{HI}} + F_{\text{crust}} \), where \( F_{\text{HI}} \) is the flux from stable hydrogen burning via the hot CNO cycle and \( F_{\text{crust}} \) is the flux from heat released by electron captures and pycnonuclear reactions deep in the crust (Brown & Bildsten 1998; Brown 2000, 2004). Following the burst ignition calculations of Cumming (2003, hereafter C03), \( F_{\text{HI}} = \epsilon_{\text{min}} y_{\text{HI}} \), where \( \epsilon_{\text{min}} = (5.8 \times 10^{19}) (Z/0.01) \text{ ergs cm}^{-2} \text{ s}^{-1} \) is the hot CNO energy production rate for a CNO mass fraction \( Z \), \( y_{\text{HI}} \) is the column depth of the layer that is burning hydrogen, and \( y_{\text{HI}} \) is the column depth at the base. For typical ignition conditions, \( y_{\text{HI}} \approx 3 \times 10^9 \text{ g cm}^{-2} \). For a given local accretion rate \( \dot{m} \) (in units of \( \text{g cm}^{-2} \text{ s}^{-1} \)) and accreted hydrogen mass fraction \( X_0 \), the hydrogen burning depth is

\[
y_{\text{HI}} = (6.8 \times 10^8 \text{ g cm}^{-2}) \left( \frac{\dot{m}}{0.1 \dot{m}_{\text{Edd}}} \right) \left( \frac{0.01}{Z} \right) \left( \frac{X_0}{0.71} \right).
\]

Here \( \dot{m}_{\text{Edd}} = (8.8 \times 10^4) [1.71/(1 + X_0)] \text{ g cm}^{-2} \text{ s}^{-1} \) is the local Eddington accretion rate. For values of \( \dot{m} \) smaller than the critical accretion rate \( \dot{m}_{\text{crit}} \approx 0.04 \dot{m}_{\text{Edd}}, y_{\text{HI}} < y_{\text{th}}, \) and there is enough time to burn all the hydrogen at the base before the helium burning becomes unstable. The burst then ignites in a pure helium layer. As in C03, we assume \( F_{\text{crust}} = \dot{m} Q_{\text{crust}} = 10^{31} \dot{m}_4 Q_{0.1} \text{ ergs cm}^{-2} \text{ s}^{-1} \), where \( m_4 = \dot{m}/10^4 \text{ g cm}^{-2} \text{ s}^{-1} \) and \( Q_{\text{crust}} = 0.1 Q_{0.1} \text{ MeV nucleon}^{-1} \) is the energy per nucleon released in the crust from pycnonuclear and electron capture reactions that escape from the surface. The preburst flux \( F_{21} = F_0/10^{31} \text{ ergs cm}^{-2} \text{ s}^{-1} \) is thus

\[
F_{21} = m_4 Q_{0.1} + \min[37X_0(1 + X_0)m_4, 4Z_{0.01} y_8],
\]

where \( Z_{0.01} = Z/0.01 \) and \( y_8 = y_8/10^8 \text{ g cm}^{-2} \).
The thermal evolution of the NS atmosphere during a burst is described by the entropy equation

$$\frac{T}{\rho} \frac{ds}{dt} = \frac{dF}{dy} + \epsilon,$$  \hspace{1cm} (3)

where $\epsilon$ is the energy release rate from nuclear burning. During the burst the entropy grows with time due to nuclear burning, and we neglect the advective accretion flow. Therefore, $T \frac{ds}{dt} = C_p \frac{dT}{dt}$, where $C_p$ is the specific heat at constant pressure. Integrating equation (3) over column depth then gives

$$\int_{y_1}^{y_2} C_p \frac{dT}{dy} \, dy = F(y_2) - F(y_1) + \int_{y_1}^{y_2} \epsilon \, dy. \hspace{1cm} (4)$$

We assume that the atmosphere is composed of two regions: a completely convective region in the range $y_c < y < y_b$, and a completely radiative region for $y < y_c$. During the burst rise, $y_b$ is constant, while $y_c$ evolves from an initial value $y_c = y_{c,\text{min}}$ to a minimal value $y_c = y_{c,\text{max}}$ and finally back to $y_c = y_b$. Demarcating the atmosphere in this way is reasonable given that the convective eddies are highly subsonic over most of the convective zone; that is, near the base, $v_{\text{conv}} \approx (F/\rho)^{1/3} \ll c_s$. Near the top of the convective zone, convection becomes inefficient and $v_{\text{conv}} \approx c_s$.

The thermal profile in the radiative region satisfies equation (1). Since the radiative region is composed primarily of freshly accreted hydrogen and/or helium, the main opacity is Thomson scattering $\kappa \approx \kappa_{\text{th}} \approx \sigma_{\text{Th}}(1 + X)/2m_p$, where $m_p$ is the proton mass and $\sigma_{\text{Th}}$ is the Thomson scattering cross section (there are corrections to $\kappa_{\text{th}}$ due to relativistic electrons and degeneracy). The opacity varies only slightly with column depth, so over much of the radiative region, $\kappa \approx \kappa_{\text{th}}$. For mixed hydrogen/helium accretion at $\dot{m} < \dot{m}_{\text{crit}}$, a pure helium layer develops in the region $y_{\text{He}} < y < y_{c}$. In this region $F_{\text{He}} = 0$, and since $F_{\text{crit}}$ is small at low $\dot{m}$, the preburst profile there is nearly isothermal.

Although the entropy in the convective region grows with time, at a given instant it is nearly spatially constant. This, in addition to the subsonic motion of the convective eddies, suggests that the thermal profile in the convective region very nearly follows an adiabat, so that $(d \ln T/d \ln y)_{\text{conv}} \approx n(y)$. where the adiabatic index $n(y) = (d \ln T/d \ln y)_{\text{ad}}$ varies with column depth; that is, $T(y_c < y < y_b) = T_b(y/y_b)^n(y)$, where $T_b$ is the temperature at the base. We define the column depth of the convective-radiative interface $y_c$ as the location where the density of the radiative solution just exceeds that of the convective solution (i.e., the neutral buoyancy criterion).

For the equation of state we use the interpolation formulae of Paczynski (1983) to account for partially degenerate electrons. Using his notation, the specific heat and adiabatic index are given by

$$C_p = \frac{1}{\rho T} \left[ \frac{3}{2} P_e + 12 P_r + \frac{P_{\text{end}}^2}{(f - 1)P_e} + \frac{P_{\text{th}}^2}{\chi_{\rho}} \right],$$ \hspace{1cm} (5)

$$n = \frac{P}{C_p \rho T} \frac{\chi_{\rho}}{\chi_{\rho}},$$ \hspace{1cm} (6)

where $P_e$, $P_r$, and $P_e = (P_{\text{end}}^2 + P_{\text{th}}^2)^{1/2}$ are the pressure due to ions, radiation, and electrons, respectively, $P_{\text{end}}$ and $P_{\text{th}}$ are an approximation to the degenerate and nondegenerate components of the electron pressure, $P = P_e + P_r + P_r$, and $f = d \ln P_{\text{ad}}/d \ln \rho$, and

$$\chi_{\rho} = \left( \frac{\partial \ln P}{\partial \ln \rho} \right) \frac{1}{P} \left( P_1 + 4 P_r + \frac{P_{\text{end}}^2}{P_e} \right),$$ \hspace{1cm} (7)

$$\chi_{\rho} = \left( \frac{\partial \ln P}{\partial \ln \rho} \right) \frac{1}{P} \left( P_1 + \frac{P_{\text{end}}^2 + f P_{\text{th}}^2}{P_e} \right).$$ \hspace{1cm} (8)

At burst onset the pressure is nearly that of an ideal gas and $n \approx 2/5$, while at late times radiation pressure contributes significantly, which in the limit $P = P_r$ gives $n = 1/4$.

2.2. Temporal Evolution of the Thermal Structure

The evolution of the convective extent $y_c$ depends on the rate at which the base temperature rises, $dT_b/dt$, and at the rate at which the thermal energy of the overlying radiative region increases. The rate of temperature change in the convective region is

$$dT/ \rho = (y/y_b)^4 [dT_b/dt + T_b \ln (y/y_b) \, dn/dt].$$

The second term is negligible compared with the first term, so that by equation (4),

$$\frac{dT_b}{dt} = \int_{y_1}^{y_2} \frac{\epsilon \, dy + F_{\text{He}} + F_{\text{crit}} - F_{\text{loss}}(y_c)}{\int_{y_c}^{y_2} C_p (y/y_b)^n(y) \, dy}. \hspace{1cm} (9)$$

where $F_{\text{loss}}(y_c)$ is the radiative flux escaping from the convective region into the overlying radiative region. Physically, the rate at which $T_b$ changes is determined by the competition between the net energy input into the convective region (i.e., the energy generated by nuclear burning and crustal heating minus the energy lost to radiating) and the energy expended in heating up the growing convective region. We describe our full nuclear reaction network in § 3.

We use mixing-length theory to estimate $F_{\text{conv}}(y_c)$, which gives the fraction of the total flux transported by convection at column depth $y$ (Hansen & Kwaler 1994),

$$F_{\text{conv}} = \frac{\nabla_{\text{rad}} - \nabla_{\text{ad}}}{\nabla_{\text{rad}}} \left( 1 - \frac{1}{\chi_{\text{Nu}}} \right),$$ \hspace{1cm} (10)

where $\nabla_{\text{ad}} = n(y), \nabla_{\text{rad}} = 3 \kappa_{\text{Y}} F_y/4aT^4$, and $\chi_{\text{Nu}} = \chi_{\text{Nu}}(y)$ is the Nusselt number describing the efficiency of convection. During the burst, convection is very efficient over most of the convective region, and $\chi_{\text{Nu}} > 1$. Because mixing-length theory only provides an order-of-magnitude estimate of $F_{\text{conv}}(y_c)$, we introduce a scaling prefactor $\lambda$ to parameterize the uncertainty in its exact value. We thus have

$$F_{\text{loss}}(y_c) = F - F_{\text{conv}}(y_c) \approx \lambda F \nabla_{\text{ad}} / \nabla_{\text{rad}} = \lambda 4a T_c^4 \nabla_{\text{ad}} / 3 \kappa_{\text{Y}},$$

where $T_c \equiv T(y_c)$. We assume $\lambda = 1$ throughout, although as we show below, $y_{c,\text{min}}$ is sensitive to the value of $\lambda$.

In Figure 1 we show the evolution of the growth timescale of the convective zone, $t_g \equiv d \ln T_b$, and the thermal timescale at the base of the radiative zone, $t_{\text{fr}} = C_p T_{\text{fr}}/F_{\text{loss}}(y_c)$. During the early stages of a burst, the convective-radiative interface is located at large column depths and $t_{\text{fr}} > t_g$. Thus, the radiative region cannot thermally adjust to the growing convective region, and the thermal profile in the radiative region is unchanged from the preignition profile (see Hanawa & Sugimoto 1982). The initial
EXPOSING THE ASHES OF X-RAY BURSTS

FIG. 1.—Thermal timescale at the convective-radiative interface $t_{\text{th}}(\gamma_c)$ (solid line) and the growth timescale of the convective zone $t_{\text{growth}}$ (dotted line), both as a function of the base temperature $T_b$ for both pure helium accretion ($X = 0$, $Y = 0.99$, $Z = 0.01$) at $\dot{m}/\dot{m}_{\text{edd}} = 0.1$ and solar abundance accretion ($X = 0.71$, $Y = 0.28$, $Z = 0.01$) at $\dot{m}/\dot{m}_{\text{edd}} = 0.1$ and 0.01. The convective zone reaches a minimum pressure approximately when the equality $t_{\text{th}} = t_{\text{growth}}$ is first satisfied.

entropy of the atmosphere, which is set by $F_0 = F_{\text{He}} + F_{\text{crust}}$, is therefore important in determining the convective evolution.\(^4\)

Eventually the convective region reaches a low enough column depth that $\dot{m}_b = \dot{m}_{\text{gr}}$. The radiative flux $F_{\text{rad}}(\gamma_c)$ can then diffuse through the entire radiative region without being overtaken by the growing convective region. Some of this flux heats the radiative region, while the remainder escapes through the photosphere. The subsequent evolution of the convective-radiative interface is determined by the column depth at which $\dot{m}_b = \dot{m}_{\text{gr}}$; that is, the radiative region continuously adjusts to $F_{\text{rad}}(\gamma_c)$, which varies due to changes in the burning rate.

Figure 1 shows results for accretion of pure helium at a rate $\dot{m}/\dot{m}_{\text{edd}} = 0.1$ and solar abundance accretion at a rate $\dot{m}/\dot{m}_{\text{edd}} = 0.01$ and 0.1 (corresponding to models He0.1, HHe0.01, and HHe0.1 as described in Table 1). The pure helium model is similar to the ignition models found by C03 in fits to the burst properties of HHe0.1 as described in Table 1). The pure helium model is similar in its structure, in units of MeV nucleon\(^{-1}\) and thus also burns in a pure helium environment (see also Cumming & Bildsten 2000). This difference in convective zone evolution between the hydrogen and helium accretors influences $t_{\text{edd}}(\gamma_c)$ (see Fig. 1) and therefore affects the light curves during burst rise, as we show in \(\S\) 4.2.

To appreciate how the compositional contrast influences the convective evolution, first consider the very early stages of the evolution, when $\dot{m}_b \gg \dot{m}_{\text{gr}}$. As noted in \(\S\) 2.1, we define the convective-radiative interface as the location of neutral buoyancy, where the density of the radiative solution just equals that of the convective solution. The density is therefore continuous across the interface, even in the presence of a compositional contrast. Hydrostatic equilibrium ensures that the pressure is also continuous. Thus, to accommodate the jump in mean molecular weight across the interface, the temperature must increase by a fractional amount $(T_c - T_r)/T_r \approx \mu_c - \mu_r$, where $T_c$ and $T_r$ are the temperature and mean molecular weight on the convective side of the interface, and $\mu_c$ and $\mu_r$ are those on the radiative side, and the equality assumes that gas pressure dominates. A thermal wave propagates outward from the interface but only travels a distance $\approx h(\gamma_c/\gamma_r)^{1/2}$, where $h$ is the scale height at $\gamma_c$, before being outrun by the growing convective zone. Thus, during the very early stages of the evolution, when $\dot{m}_b \gg \dot{m}_{\text{gr}}$, the time-dependent solution yields a very sharp temperature gradient. We approximate the gradient as a jump as long as $\dot{m}_b \gg \dot{m}_{\text{gr}}$, as shown in Figure 2. Our approximation becomes inaccurate as the convective zone extends out to lower pressures and $\dot{m}_b \rightarrow \dot{m}_{\text{gr}}$, as then the thermal wave can "get ahead" of the growing convective zone. This approximation thus prevents us from resolving the very early light curve of the burst rise. We resolve the light curve well before the luminosity exceeds the accretion luminosity, however.

Once $\dot{m}_b = \dot{m}_{\text{gr}}$, there is enough time for heat diffusion (via radiation) to smooth out the temperature gradient over a scale

\[\dot{m}/\dot{m}_{\text{edd}} = 0.01\] for accretion of helium at a rate $\dot{m}/\dot{m}_{\text{edd}} = 0.1$ and solar abundance accretion at a rate $\dot{m}/\dot{m}_{\text{edd}} = 0.01$ and 0.1 (corresponding to models He0.1, HHe0.01, and HHe0.1 as described in Table 1). The pure helium model is similar to the ignition models found by C03 in fits to the burst properties of HHe0.1 as described in Table 1).

Model Burst Ignition Models
| Col. (1) | $y_b$ | $X_0$ | $Z_{\text{CNO}}$ | $Q_{\text{crit}}$ | Network |
|---------|-------|-------|-----------------|-----------------|---------|
| He[0.0] | 3.0   | 0.0   | 0.01            | 0.1             | Full    |
| He0.1   | 3.0   | 0.0   | 0.01            | 0.1             | $\alpha$ only |
| He0.1b  | 5.0   | 0.0   | 0.1             | Full            |
| He0.1Q  | 3.0   | 0.0   | 0.01            | 0.2             | Full    |
| HHe[0.0] | 3.0 | 0.71 | 0.1             | Full            |
| HHe0.01X | 3.0 | 0.01 | 0.1             | Full            |
| HHe0.01XZ | 3.0 | 0.0001 | 0.1 | 0.1 | Full |

Notes.—Col. (1): Model name, where $[\dot{m}]$ denotes the accretion rate in units of $\dot{m}_{\text{edd}}$. We consider models spanning the range $0.01 \dot{m}_{\text{edd}} - 0.2 \dot{m}_{\text{edd}}$. Col. (2): Ignition column depth $y_b = y_b/10^8$ g cm\(^{-2}\). Col. (3): Accreted hydrogen fraction $X_0$. Col. (4): CNO mass fraction $Z_{\text{CNO}}$. Col. (5): Energy per nucleon released in the crust, $Q_{\text{crit}}$ in units of MeV nucleon\(^{-1}\). Col. (6): Reaction network used.

\(^4\) When the density and temperature first get high enough for helium to ignite, the burning timescale may be longer than or comparable to the thermal timescale. The flux from this early burning may change the preburst profile of the radiative region slightly, although eventually the burning becomes nonlinear and the burning time becomes much shorter than the thermal time. To examine how this might affect the growth of the convective region, we computed $\dot{m}(t)$, assuming an artificially high preburst flux (e.g., $5F_{\text{gr}}$). We find that this effect may increase $\gamma_c$, min by as much as a factor of 3.

\(^5\) Since the particle diffusion timescale is much longer than $t_{\text{gr}}$, there is not adequate time for semiconvection to develop. We therefore approximate the composition gradient as a discontinuous step function.
height, and the radiative region heats up. During this stage the convective zone returns to higher pressure and the compositional contrast between burned and unburned matter is at a fixed location $y = y_{c, \text{min}}$ and is thus entirely within the radiative zone. The discontinuous change in opacity at $y_{c, \text{min}}$ results in a change of slope of the radiative temperature profile at that location, as can be seen in Figure 2.

We now obtain a rough estimate of $y_{c, \text{min}}$ by setting $t_{\text{fb}} = t_{\text{gr}}$ and solving for $y_c$, given the preburst thermal profile set by $T_0$. Using the mixing-length theory relation for $F_{\text{loss}}(y_c)$ and assuming an adiabatic profile in the convective zone so that $T_c = T_b(y_c/y_b)^{\frac{3}{3}}$ yields $t_{\text{fb}} = \frac{3c}{3}\frac{3}{\gamma_y} y_{c}^{3/4}/4\alpha c \lambda \nu T_{b}^{1/3}$. Then, since $t_{\text{gr}} \approx C_{P} T_{b}/\epsilon$, we have, upon setting $t_{\text{fb}} = t_{\text{gr}}$, that

$$y_c = \left(\frac{4\alpha c \lambda \nu T_{b}^{3/4}}{3\epsilon \nu \gamma_y b^{3}}\right)^{1/(2-3n)}.$$  \hfill (11)

At $y_c$ the density and pressure are continuous, and at $y_{c, \text{min}}$ the pressure is nearly that of an ideal gas. We can thus eliminate $T_b$ from the above expression, using $T_c = \mu_c T_c(y_c)/\mu_r$ and $T_c(y_c) \approx (3\nu \gamma_y, F_0/\alpha c)^{1/4}$. If we solve for $y_c$, we get

$$y_{c, \text{min}} \approx \frac{4\lambda \nu y_b^{4}}{\mu_r^{4} F_0 T_{b}^{1/3}} \epsilon^{(1+n)}$$

$$\approx (2 \times 10^4 \text{ g cm}^{-2}) F_{21}^{5/7} \left(\frac{\mu_c}{\mu_r}\right)^{20/7} y_b^{-15/7}$$

$$\times \left(\frac{y_b}{3 \times 10^8 \text{ g cm}^{-2}}\right)^{2/7} \left(\epsilon_{18}/Y_{3}^{3/2} \right)^{-5/7},$$  \hfill (12)

where the dependence on $\dot{m}$ is obtained by substituting equation (2) for $F_{21} = F_0/10^{21}$ ergs cm$^{-2}$ s$^{-1}$. In the lower expression, $y_b$ is the helium mass fraction in the burning layer, and we assumed $n = 2/5$, $\lambda = 1$, and an energy generation rate $\epsilon_{18} = (10^{18} \text{ ergs g}^{-1} \text{ s}^{-1}).$ For helium burning, $\epsilon_{18}/Y_{3}^{18} = (5.3 \times 10^2) \rho_b^{1/2} \exp\left(-4.4/T_{b}\right)$ (Hansen & Kawaler 1994), where $T_b = T_b/10^{10} \text{ K}$ and $\rho_b = \rho_b/10^{6} \text{ g cm}^{-3}$. The estimate and its scaling with burst parameters agree reasonably well with our full evolution calculations presented in § 4. Note the strong dependence on compositional contrast between the accreted material and the ashes.

As we show in § 4, the pure helium accretion models and the solar abundance accretion models with $\dot{m} \lesssim 0.05\dot{m}_{\text{Edd}}$ achieve a super-Eddington luminosity that drives a radiative wind capable of ejecting material located at column depths $y < y_{\text{wind}} \approx 0.01 y_b$. Since the convective zone reaches pressures $y_{c, \text{min}} \ll y_{\text{wind}}$ in these models, one expects ashes of burning to be among the wind ejecta.

3. NUCLEOSYNTHESIS FROM PURE HELIUM BURNING AND MIXED HYDROGEN AND HELIUM BURNING

The full nuclear reaction network used in this work contains 394 nuclei from H to Sr ranging from the proton drip line to a few mass units beyond stability and accounts for all relevant proton- and $\alpha$-particle-induced reactions and the corresponding inverse processes. The reaction rates have recently been updated, and a full discussion will be presented in a forthcoming paper. Here we are only concerned with nucleosynthesis in the early, convective stages of an X-ray burst. The relevant nuclear reactions involve only nuclei lighter than iron, and we restrict the following discussion to those nuclei. When no experimental information is available, reaction rates are taken from $sd$-shell model calculations (Herndl et al. 1995), $fp$-shell model calculations (Fisker et al. 2001), or the statistical model NON-SMOKER (Rauscher & Thielemann 2000). Experimental rates were taken from the NACRE (Angulo et al. 1999) and Iliadis et al. (2001) compilations, with some exceptions. As far as the reaction rates relevant to this work are concerned, these exceptions include...
The 14O(\alpha, \gamma) reaction plays an important role. For example, beyond Mg, \(3\alpha\) captures in some cases, such as for the flow from \(^{36}\text{Ar}\) to \(^{40}\text{Ca}\) or the \(^{28}\text{Si}\)\(-\text{chain}\) up to \(^{32}\text{S}\) for temperatures \(>1\) GK, by the much faster \(^{12}\text{C}(\alpha, \gamma)^{16}\text{O}\) reaction sequence. The \(3\alpha\) reaction, where we used the new experimental rate from Fynbo et al. (2005), although for the temperatures relevant here this rate is identical to the one in Caughlan & Fowler (1988). For \(^{12}\text{C}(\alpha, \gamma)^{16}\text{O}\) we use the reaction rate from Buchmann (1996). The \(^{14}\text{O}(p, \gamma)^{17}\text{F}\) reaction is from Hahn et al. (1996), and the \(^{13}\text{N}(p, \alpha)^{16}\text{O}\) and \(^{24}\text{Mg}(\alpha, \gamma)^{28}\text{Si}\) reactions are from Caughlan & Fowler (1988). Weak interaction rates were taken from Fuller et al. (1982) when available; otherwise, ground-state \(\beta\)-decay rates from NUBASE (Audi et al. 2003) were employed. Weak interactions, however, do not play a role in any of our models.

We calculate the nucleosynthesis and nuclear energy generation only at the base of the convective zone \(y_b\). To estimate the integral of energy generation over column depth in the calculation of \(dT_b/dt\) and \(t_{gr}\) (see eq. [9]), we first take the ratio of \(\int \epsilon(T_b, y) \, dy\) to \(\epsilon(T_b, y_b)\), assuming only \(\alpha\)-capture reactions. We find that, to a good approximation, the ratio is constant with temperature, and we use this constant factor to rescale the value of \(\epsilon(T_b, y_b)\) that is calculated using the full network. Our initial composition is the composition at \(y_b\) at ignition, and we neglect compositional changes due to convective mixing as the convection zone moves outward. Such mixing might be relevant for the hydrogen/helium accretion models in which the convection zone at later stages extends into hydrogen-rich regions. However, the effect is probably comparable to an increase in accretion rate, which also results in a larger hydrogen abundance. Therefore, our neglect of convective mixing might simply translate to a slight underestimate of the accretion rate in our sequence of calculated models.

Figure 3 shows a typical reaction flow during pure helium burning, and Figure 4 shows the composition as a function of time and base temperature \(T_b\). After a slow rise from 0.2 to 0.4 GK, temperatures rise quickly, typically within 10–100 ms, from 0.4 GK to peak temperatures set by \(y_b\) of 1.8 GK. The reaction flow is mainly characterized by a sequence of \(\alpha\)-captures into the \(^{44}\text{Ti}\) region. However, reactions off the \(\alpha\)-chain also play an important role. For example, beyond Mg, \((\alpha, p)\)-\((p, \gamma)\) sequences carry a substantial part of the flow and in fact dominate in some cases, such as for the flow from \(^{36}\text{Ar}\) to \(^{40}\text{Ca}\) or the flow beyond \(^{44}\text{Ti}\). The importance of \((\alpha, p)\)-\((p, \gamma)\) links and their inverses in helium burning has been pointed out before (Hashimoto et al. 1983; Timmes et al. 2000).

The most important effect of the reactions off the \(\alpha\)-chain, however, is the bypass of the \(^{12}\text{C}(\alpha, \gamma)^{16}\text{O}\) reaction, which is the slowest reaction in the \(\alpha\)-chain up to \(^{32}\text{S}\) for temperatures \(>1\) GK, by the much faster \(^{12}\text{C}(p, \gamma)^{13}\text{N}(\alpha, p)^{16}\text{O}\) reaction sequence. The bypass requires protons as catalysts, which are produced by \((p, \gamma)\) reactions on heavier \(\alpha\)-chain nuclei such as \(^{24}\text{Mg}\), \(^{32}\text{S}\), and \(^{36}\text{Ar}\). Initially, helium burning proceeds via the \(3\alpha\) reaction, building up \(^{12}\text{C}\) with some weak reaction flow through \(^{12}\text{C}(\alpha, \gamma)\) and successive \(\alpha\)-captures. At \(^{24}\text{Mg}\), \(^{32}\text{S}\), and \(^{36}\text{Ar}\), \((\alpha, p)\) branchings are significant and lead to the buildup of a small amount of protons. For example, in the pure helium \(\dot{m}/\dot{m}_{\text{Edd}} = 0.1\) model the proton abundance reaches \(3 \times 10^{-9}\) at a temperature of about 1 GK, which is sufficient for the \(^{12}\text{C}(p, \gamma)^{13}\text{N}(\alpha, p)^{16}\text{O}\) bypass to take over the reaction flow from \(^{12}\text{C}\) to \(^{16}\text{O}\). At this point a runaway effect sets in, as any increase in flow beyond \(^{12}\text{C}\) increases the production of protons, which further increases the efficiency of the \(^{12}\text{C}(p, \gamma)^{13}\text{N}(\alpha, p)^{16}\text{O}\) bypass. Only 2 ms later, at a temperature of 1.2 GK, the flow from \(^{12}\text{C}\) to \(^{16}\text{O}\) exceeds the flow via the \(3\alpha\) reaction, and \(^{12}\text{C}\) is quickly depleted. As a result, the \(^{12}\text{C}(p, \gamma)^{13}\text{N}(\alpha, p)^{16}\text{O}\) reaction sequence completely dominates the reaction flow toward heavy elements during the convective burning stage. This is a very robust effect, which is independent of the initial metallicity, and will always occur during helium burning under high-temperature conditions that lead to the production of \(^{22}\text{Mg}\) or heavier nuclei.

The rapid onset of the \(^{12}\text{C}(p, \gamma)^{13}\text{N}(\alpha, p)^{16}\text{O}\) bypass leads to a sudden spike in energy production, which causes the drop in \(t_{gr}\) and \(n_b\) to ~millisecond timescales and the second dip in \(y_b\), as seen in Figures 1 and 2, respectively. As we show in § 4, the bypass occurs in all models burning helium-rich material, regardless of \(\dot{m}\) (see Fig. 6), and is directly observable as a rapid increase in the rise time of the burst light curve (see Fig. 8). For comparison, we also calculated the pure helium \(\dot{m}/\dot{m}_{\text{Edd}} = 0.1\) model with a simple \(\alpha\)-chain network, the results of which are shown in Figures 6, 8, and 13. Another consequence of bypassing the \(^{12}\text{C}(\alpha, \gamma)\) reaction is the depletion of \(^{12}\text{C}\) and the enhancement of the production of heavier \(\alpha\)-chain nuclei in the
final composition. As Figure 13 shows, a pure helium \( \dot{m}/\dot{m}_{\text{Edd}} = 0.1 \) model with just an \( \alpha \)-chain network would predict the main product to be \( ^{12}\text{C} \), while with the full reaction network the main products are \( ^{28}\text{Si} \) and \( ^{32}\text{S} \), with negligible amounts of carbon. In addition, the presence of protons leads to proton-capture reactions producing small amounts of nuclei off the \( \alpha \)-chain. The very small amounts of nuclei around \( A = 60 \) are produced by proton capture of iron peak nuclei present in the accreted material.

For pure helium accretors, except for model He0.1b (see Table 1), the composition at the time that the convective zone has retreated back to the base is mainly doubly magic \( ^{40}\text{Ca} \) and unburned \( ^{4}\text{He} \), with some \( ^{36}\text{Ar} \) and \( ^{44}\text{Ti} \). The most abundant non-\( \alpha \)-chain nucleus is \( ^{39}\text{K} \). Model He0.1b ignites deeper, at a higher mass density, resulting in a higher peak temperature. Burning therefore proceeds beyond \( ^{44}\text{Ti} \) with a sequence of alternating \( (\alpha, p) \) and \( (p, \gamma) \) reactions up to \( ^{52}\text{Fe} \). Besides unburned helium, the composition is then dominated by \( ^{44}\text{Ti} \) and \( ^{48}\text{Mn} \). For mixed hydrogen/helium accretors that undergo pure helium burning at low accretion rates (e.g., models HHe0.01 and HHe0.01X), the composition is very similar to the pure helium accretion models, except for a somewhat lower \( ^{36}\text{Ar} \) and higher \( ^{44}\text{Ti} \) production. This difference is due to the faster rise in temperature in the hydrogen/helium models, a consequence of the smaller convective extent in the presence of hydrogen and thus a smaller convective mass to heat up for the same energy input (see the \( t_p \) curve for models He0.1 and HHe0.01 in Fig. 1).

Although some protons are produced during helium burning, the presence of large amounts of hydrogen at ignition in models such as HHe0.1 (\( X = 0.4 \)) and HHe0.01XZ (\( X = 0.08 \)) leads to a drastic change in the nuclear reaction sequence. Figure 3 shows a typical reaction flow, and Figure 5 shows the composition as a function of time and base temperature \( T_b \). The flow is

![Figure 4](image1.png)

**Fig. 4.**—Mass fraction of various isotopes as a function of time (bottom axis) and base temperature (top axis) for model He0.1. The left panel shows the full evolution, and the right panel zooms in on the time during which the \( ^{12}\text{C} (p, \gamma) ^{13}\text{N} (\alpha, p) ^{16}\text{O} \) bypass takes over the reaction flow from \( ^{12}\text{C} \) to \( ^{16}\text{O} \). Only those isotopes whose peak mass fraction exceeds 0.05 are shown.

![Figure 5](image2.png)

**Fig. 5.**—Mass fraction of various isotopes as a function of time (bottom axis) and base temperature (top axis) for model HHe0.1. Only those isotopes whose peak mass fraction exceeds 0.05 are shown.
EXPOSING THE ASHES OF X-RAY BURSTS

Fig. 6.—Evolution of the top of the convective zone \( y_c \) as a function of base temperature \( T_b \) for burst models with pure helium accreted material (left) and mixed hydrogen-helium accreted material (right), as described in Table 1. The two unlabeled helium models are, from bottom to top, He0.04 and He0.1. Model He0.1a (dash-dotted line) was calculated using a simple \( \alpha \)-chain reaction network. The horizontal dashed line denotes the column depth at which \( y = 0.01 y_\text{w} \), corresponding approximately to the column depth \( y_{\text{wind}} \).

4. RESULTS

The radiative wind generated by the super-Eddington luminosity of an RE burst will eject some ashes of burning if (1) the fraction of the accreted mass ejected during the wind satisfies \( \frac{\Delta M_p}{M_{\text{acc}}} = y_{\text{wind}}/y_h > y_{\text{c, min}}/y_h \), where \( M_{\text{acc}} \approx 4\pi R^2 y_h \), and (2) the wind is generated subsequent to the convective region reaching \( y_c < y_{\text{wind}} \). In § 4.1 we calculate \( y_{\text{c, min}} \) and demonstrate that condition 1 is satisfied for a wide range of burst ignition models. In § 4.2 we calculate the light curve during burst rise and show that condition 2 is satisfied for those same models.

We first note that since the gravitational binding energy at the surface is \( \approx 100 \) times greater than the helium burning energy release per unit mass, at most \( \approx 1\% \) of the atmospheric energy can be ejected by the wind. To obtain an estimate of \( \Delta M_p/M_{\text{acc}} \) suppose that the surface luminosity as measured by an observer at infinity \( L \) is super-Eddington so that the energy loss rate at the surface \( \dot{E} \approx L + M_p c^2 / (1+z) \). At the photospheric radius of the wind, \( z_p \approx R \) (see §5.3), the luminosity is very nearly the local Eddington value, \( \dot{E} \approx 4\pi G M c (1+z_p) / \kappa \approx 4\pi G M c / \kappa \). Equating the energy loss rate at the surface to that at the photosphere yields the mass-loss rate due to the wind (Paczynski & Proszynski 1986; see also Wallace et al. 1982; Yahel et al. 1984; Joss & Melia 1987; Nobili et al. 1994),

\[
\frac{\dot{M}_w}{M_{\text{acc}}} \approx \frac{L_{\text{Edd}}}{c^2} \left( \frac{L}{L_{\text{Edd}}} - 1 \right) \left( 1 - \frac{1}{1+z} \right)^{-1} \approx \left( 10^{18} \text{ g s}^{-1} \right) \left( \frac{0.2 \text{ cm}^2 \text{ g}^{-1}}{\kappa} \right) \left( \frac{L}{L_{\text{Edd}}} - 1 \right). \tag{13}
\]

The burst duration \( \Delta t \approx M_{\text{acc}} Q_{\text{nuc}} / L \), where \( Q_{\text{nuc}} = 1.6 + 4.0 / (X) \) MeV nucleon\(^{-1} \) is the nuclear energy release and \( X \) is the mass-weighted mean hydrogen fraction in the burning layer (C03), so that for a pure helium burst

\[
\frac{\Delta M_p}{M_{\text{acc}}} \approx \frac{\dot{M}_w}{M_{\text{acc}}} \Delta t \approx \left( 1 - \frac{L_{\text{Edd}}}{L} \right) \frac{Q_{\text{nuc}}}{\dot{L}_{\text{Edd}}} = 0.003 \left( 1 - \frac{L_{\text{Edd}}}{L} \right). \tag{14}
\]

4.1. Dependence of the Convective Extent on Burst Parameters

In Figure 6 we show the time variation of the convective extent \( y_c \) during a burst, with the progression marked by the temperature at the base of the burning layer, \( T_b \), rather than by time. We consider burst ignition models assuming both mixed hydrogen/helium accretion and pure helium accretion (see Table 1); the latter models are similar to those obtained by C03 in fits to the burst properties of 4U 1820–30. We consider models over a range of accretion rates, ignition column depths, metallicities, and crustal heat fluxes.

The evolution of \( y_c \) is sensitive to the assumed ignition conditions and to the energy production from non-\( \alpha \)-capture reactions (see §3). Ignition conditions that maximize the peak flux of a burst and minimize the entropy of the atmosphere during the fuel accumulation stage yield bursts with the smallest value of \( y_{\text{c, min}} \). This is because the larger the peak flux, the faster \( T_b \) rises,
and hence the shorter the growth timescale $t_{gr}$ at a given $T_b$. The
convective region must therefore reach out to lower column depths
before $t_{hb}(y_c) = t_{gr}$. As illustrated in Figure 7, such ignition
conditions are best satisfied at low $\dot{m}$. The reason is twofold.
First, for $\dot{m} \lesssim 0.04\dot{m}_{Edd}$ and solar abundance accreted material,
a pure helium layer develops at the base of the atmosphere, and
the peak luminosity is considerably greater than that of bursts
with mixed hydrogen/helium burning layers. Second, because the
preburst flux $F_0 \propto \dot{m}$, the lower the value of $\dot{m}$, the smaller
the initial atmospheric entropy.

For the models considered, the fraction of accreted mass lying above the convective region at its maximum extent is in the range
$10^{-4} \lesssim \Delta M/M_{\text{acc}} \lesssim 10^{-3}$. Ashes are therefore likely to be ejected in the wind.

### 4.2. Light Curve During Burst Rise and Radiative Winds

We determine the rising portion of the burst light curve by calculating the radiative flux loss at the photosphere $F_{\text{ph}}$, as $T_b$ increases. To obtain $F_{\text{ph}}$, we integrate equation (1) inward (assuming the radiative zero solution at the outer boundary), varying the flux at the top until the radiative solution intersects the convective solution at the column depth $y_c$. Initially, $t_{hb} > t_{gr}$ and $F_{\text{ph}} = 0$. Eventually $t_{hb} = t_{gr}$ and $F_{\text{ph}}$ rises as the radiative region heats up due to the rise in $F_{\text{low}}(y_c)$. In Figure 8 we show the rising portion of the light curve for the same models as Figure 1 and also for a pure helium model using only the $\alpha$-chain network. We plot the luminosity $L = 4\pi R^2 F_{\text{ph}}$ as a function of $t - t_0$, where $t_0$ corresponds to the time when the equality $t_{hb} = t_{gr}$ is first satisfied.

The rise time at a given time $t$ is set by the thermal time at the convective-radiative interface $t_{fr}(y_c(t))$. As the differences between the He0.1 and He0.1$\alpha$ models illustrate, the shape and rise time of the light curves are sensitive to the energy generation rate. The sharp submillisecond rise in models HHe0.01 and He0.1 occurs at the onset of the $^{12}\text{C}(p, \gamma)^{13}\text{N}(\alpha, P)^{16}\text{O}$ bypass ($\lesssim 3$). The bypass leads to a spike in energy production that in turn drives the convective zone to lower pressures. Since

$$T(y_c) = T_e$$ is very high by this time, $t_{hb} = 3\pi G \rho y^2_{c}/4ac\Delta T_c$ becomes very small (see Fig. 1). In the He0.1 model, the light curve initially rises relatively slowly ($\approx 10$ ms) and then, once the temperature is high enough for the bypass to take over, it sharply rises. This is distinct from the HHe0.01 model, which rises sharply at the outset because $t = t_0$ nearly coincides with the onset of the bypass.

Depending on the ignition conditions, the base temperature can get sufficiently high that the flux becomes super-Eddington. In Figure 9 we show the local Eddington luminosity $L_{\text{Edd}}(y)$ as a function of column depth for the models HHe0.01, HHe0.1, and He0.1 at the end of the calculation, when the atmosphere is completely radiative (i.e., $y_c = y_0$). The increase in $L_{\text{Edd}}(y)$ with depth is the result of the decrease in $\kappa_{\text{es}}$ with temperature due to relativistic corrections. The discontinuity in $L_{\text{Edd}}(y)$ for the mixed hydrogen/helium models occurs at $y_c_{\text{min}}$ and is due to the discontinuity in hydrogen abundance between the burned and unburned material. Such a discontinuity could explain the possible bimodal distribution of the observed burst peak luminosities (see Kuulkers et al. 2003 for a recent discussion), as first pointed out by Sugimoto et al. (1984) and investigated in greater detail by Ebisuzaki & Nakamura (1988). They proposed that the ejection of the hydrogen-rich envelope during an RE burst occasionally exposes the underlying helium-rich region and results in the Eddington limit transitioning from that of hydrogen-rich material to that of hydrogen-poor material.

We find that the peak luminosity exceeds $L_{\text{Edd}}(y)$ for $\dot{m} \lesssim 0.1\dot{m}_{Edd}$ if the accreted material has solar abundance ($\dot{m}$ can be higher if the material is helium-rich) and that the Eddington value is first exceeded near the top of the atmosphere. Such systems are expected to develop a radiative wind. Eventually, the flux loss exceeds the flux from nuclear burning and the atmosphere cools, although we do not calculate this portion of the light.
curve. Note that since the convective zone reaches a minimum pressure at \( t \approx t_0 \), bursts that become super-Eddington achieve \( y_c < y_{\text{wind}} \) well before the wind is generated.

In Figure 10 we show the dependence of the rise time on \( \dot{m} \) for the pure helium models and the mixed hydrogen/helium models. We calculate rise times during the early portion of the rise, when the surface flux just exceeds the accretion flux (\( \tau \text{detect} \)), and during the later portion of the rise, taken to be the time for the burst to rise from 10% to 50% of its peak luminosity (\( \tau_{1/2} \)). Since \( \tau_{\text{detect}}(y_c, \text{min}) \) sets the initial rise time, models with smaller values of \( \dot{m}/c_{\text{min}}^2 T_{c, \text{min}}^2 \) tend to have smaller values of \( \tau \text{detect} \), which, to a first approximation, corresponds to models with lower values of \( \dot{m} \). For helium accretors, \( \tau_{1/2} \) is nearly independent of \( \dot{m} \) because the convective zone evolution at late times is very similar among models that differ only in \( \dot{m} \) (see Fig. 6). By contrast, for accretion of hydrogen-rich material, \( \tau_{1/2} \) increases by a factor of \( \approx 50 \) between \( \dot{m}/c_{\text{Edd}}^2 = 0.02 \) and 0.1. In their multizone numerical simulations of X-ray bursts, Woosley et al. (2004) consider models that accrete mixed hydrogen/helium with solar abundance at \( \dot{m}/c_{\text{Edd}} = 0.02 \) and 0.1. They follow several sequences of bursts for each model and obtain values of \( \tau_{1/2} \) in the range \( 0.51 \times 10^{-3} - 32.1 \times 10^{-3} \) s for \( \dot{m}/c_{\text{Edd}} = 0.02 \) and \( 0.51 - 0.66 \) s for \( \dot{m}/c_{\text{Edd}} = 0.1 \), both of which are in good agreement with our estimates. Our rise times are also in broad agreement with the relevant models of earlier numerical simulations (e.g., Taam 1981; Ayles & Joss 1982; Wallace et al. 1982).

5. DETECTING THE NUCLEAR BURNING ASHES

In § 4 we showed that the convective zone reaches sufficiently low pressures during RE bursts that \( y_{c, \text{min}} < y_{\text{wind}} \) over a broad range of burst parameters. Thus, ashes of nuclear burning can be ejected by the radiative wind of RE bursts. Furthermore, when the photosphere settles back down to the NS surface following the RE phase, it is laced with heavy-element ashes. In this section we describe the composition of the ejected and exposed ashes (§ 5.1) and address whether they can be detected. In § 5.2 we discuss the possibility that some of the \( p \)-nuclei found in the solar system might owe their origin to ash ejection during RE bursts. In §§ 5.3 and 5.4 we determine the expected strength of photoionization edge features from ashes ejected in the wind and those exposed at the NS surface.

5.1. Composition of Ejected and Exposed Ashes

Just prior to the onset of the wind, the convective zone has receded to the base and the atmosphere has a stratified compositional structure. Throughout the region \( y < y_{c, \text{min}} \), the composition is that of the unprocessed accreted material. As we show in Figures 11 and 12, for \( y > y_{c, \text{min}} \) the composition is determined by the burning stage at the moment \( y_c(t) = y \) during the convective zone’s retreat to the base. Results are shown for models He0.1 and HHe0.1. The discontinuity at \( y = 3 \times 10^5 \) g cm\(^{-2} \) in model He0.1 is due to the rapid nucleosynthesis that occurs once the \( ^{12}\text{C} + (p, \gamma)^{13}\text{N}(\alpha, p)^{16}\text{O} \) bypass takes over the reaction flow from \( ^{12}\text{C} \) to \( ^{16}\text{O} \). The composition at the base of the wind is approximately that at \( y = 0.01y_{\text{Edd}} = 3 \times 10^8 \) g cm\(^{-2} \). In Figure 13 we show the ash composition at \( y/y_{\text{Edd}} = 0.01 \approx \Delta M_{\text{acc}}/M_{\text{acc}} \) for models He0.1, He0.1, and HHe0.1. While helium comprises the largest fraction of the mass in all three models, the models’ overall abundance distributions differ significantly from one another due to the impact of the \( ^{12}\text{C} + (p, \gamma)^{13}\text{N}(\alpha, p)^{16}\text{O} \) bypass and differences in the initial composition at the base of the burning layer. Some of the proton-rich isotopes shown in the above figures have half-lives comparable to the duration of the atmosphere’s convective phase. Although we do not show the decay products in these figures, we do account for such decays in our calculation of photoionization edges below.

5.2. Ejection of \( p \)-Nuclei

In their numerical simulations of X-ray bursts, Woosley et al. (2004) find that bursts ignite in the ashes of previous bursts. If so, these endpoint ashes, which are heavier than those processed during the ongoing burst at the time \( y_c/y_{\text{Edd}} = 0.01 \), are mixed
throughout the convective region (we do not show such ashes in Figs. 11, 12, or 13). Using a reaction network that extends up to Xe, Schatz et al. (2001) find that in models where solar abundance material is accreted at $\dot{m} = 0.3\dot{m}_{\text{Edd}}$ (i.e., burning in a hydrogen-rich environment), the endpoint of rp-process burning is a closed SnSbTe cycle that naturally limits rp-process nucleosynthesis to light $p$-nuclei. They find overproduction factors (relative to solar abundance) of $\sim 10^8$ for the $p$-nuclei $^{92}\text{Mo}$ and $^{96}\text{Ru}$ and $\sim 10^9$ for the $p$-nucleus $^{98}\text{Ru}$ (see also Schatz et al. 1998). Standard $p$-process scenarios are unable to adequately explain the observed solar system abundances of these $p$-nuclei (for a review, see Wallerstein et al. 1997).

Whether ash ejection during RE bursts can account for the observed solar system abundances depends on the amount of $p$-process material ejected into the interstellar medium by an RE burst and the event rate of such bursts over the Galactic lifetime. In order to produce large amounts of $p$-nuclei, the burning layer must be hydrogen-rich ($\dot{m} \approx 0.3\dot{m}_{\text{Edd}}$), while RE bursts require helium-rich burning layers and thus $\dot{m} \leq 0.05\dot{m}_{\text{Edd}}$. However, this does not preclude RE bursts from ejecting $p$-nuclei, as accretion rates in bursting low-mass X-ray binaries are observed to vary by factors of a few, with individual systems undergoing transitions from hydrogen-rich to helium-rich burning over timescales of a year (as evidenced by variations in burst duration and peak fluxes;
Once matter reaches the sonic point, larger at earlier times in our Galaxy. The convective zone is located at \( \nu_c \approx \nu_{\text{wind}} \) and is receding to higher pressures. Results are shown for models He0.1, He0.1, and He0.1α.

To determine the fractional amount \( \eta \) of ashes that must be ejected in order to account for the observed solar system abundance of \( p \)-nuclei, we assume a \( p \)-nuclei overproduction factor \( \xi = 10^6 \) and a Galactic disk mass \( M_{\text{disk}} = 4 \times 10^{10} \) (Klypin et al. 2002). Currently, there are \( \approx 10 \) active X-ray burst systems at \( M \sim 10^{-9} M_\odot \) yr\(^{-1} \) and \( \approx 100 \) at \( M \sim 10^{-10} M_\odot \) yr\(^{-1} \) (Lewin et al. 1995). If we assume that this is representative of the population count over the Galactic lifetime, then the total amount of mass accreted by all RE burst systems over \( 10^{15} \) yr is \( M_{\text{acc, tot}} \approx 100 M_\odot \). A fraction

\[
\eta = 0.4 \left( \frac{M_{\text{disk}}}{4 \times 10^{10} M_\odot} \right) \left( \frac{\xi}{10^6} \right)^{-1} \left( \frac{M_{\text{acc, tot}}}{100 M_\odot} \right)^{-1}
\]

of all accreted material must therefore be ejected. Thus, \( \eta \) is a factor of \( \sim 100 \) too high, given that only \( \approx 1\% \) of all accreted matter is ejected. The discrepancy can be overcome if, for example, the Galactic distribution of \( p \)-nuclei is inhomogeneous (i.e., the solar abundance is higher than the Galactic mean by a factor of \( \sim 100 \)) or the population of bursting systems was much larger at earlier times in our Galaxy.

### 5.3. Spectral Edges in Wind Outflow

During the RE phase, an optically thick, transonic, radiation-driven wind forms. The sonic point of the wind lies \( 10 \)–\( 100 \) km above the NS surface and the photosphere, defined as the location where the effective optical depth \( \tau_e \approx \kappa r_p \) is near unity, is a factor of \( \sim 10 \) farther out (Paczynski & Proszynski 1986; Joss & Melia 1987; Nobili et al. 1994). Once matter reaches the sonic point, it is essentially unbound from the NS and is ejected to infinity.

The ejected matter is at a sufficiently low temperature during the RE phase that some heavy elements bind with one or more electrons. The resulting column density of hydrogen-like ions above the photosphere is thus

\[
N_{\text{wind}} \sim f(A, Z) \zeta(T, \rho) N_c \sim (10^{20} \text{ cm}^{-2}) \left( \frac{f}{10^{-4}} \right),
\]

where \( f(A, Z) \) is the abundance by number of element \( Z \) with mass number \( A \), \( \zeta(T, \rho) \) is the fraction in the hydrogen-like state at a given temperature and density from the Saha equilibrium, and \( N_c \approx \sigma_{\text{fs}}^{-1} \) is the electron column density. This corresponds to an optical depth to a photoionization edge of \( \tau \approx N_{\text{wind}} \sigma_{\text{fs}} \sim 1 \), where the bound-free cross section \( \sigma_{\text{fs}}(E) = (6 \times 10^{-18})(E_e/E)^{3/2} \text{ cm}^2 \) and \( E_e \approx 13.6Z^2 \) eV is the edge energy (Rybicki & Lightman 1979). The effective temperature is low enough at the photosphere that even for metals at solar abundances, \( \tau \sim 1 \). That the ejected ashes have abundances much larger than solar improves the likelihood of detection.

We now calculate \( N_{\text{wind}} \) more exactly and determine the resulting equivalent width (EW) of the photoionization edge. The column density of hydrogen-like ions for a given species is

\[
N_{\text{wind}} = \frac{f(A, Z)}{m_p} \int_{r_p}^{\infty} \zeta(T(r), \rho(r)) \rho(r) dr.
\]

The recombinations are nearly instantaneous on the timescale of the wind flow, i.e., \( t_{\text{rec}} \ll t_{\text{flow}} \sim r/v(r) \approx 0.1 \text{ s} \), where \( v(r) \) is the fluid velocity at \( r \) and \( t_{\text{rec}} \equiv 1/n_e(\sigma_{\text{fs}}) \), where \( \sigma_{\text{fs}} \) is related to \( \sigma_{\text{fs}} \) through the Milne relation (Rybicki & Lightman 1979). We use the results of Paczynski & Proszynski (1986), who calculate general relativistic models of radiation-driven winds from NSs to obtain values for \( r_p \), \( T_{\text{ph}} \), and \( \rho(r_{\text{ph}}) \) as a function of the mass outflow rate \( M_w \). For example, at \( M_w = 10^{18} \text{ g s}^{-1} \), \( r_p = 3 \times 10^3 \text{ cm} \), \( T_{\text{ph}} = 5 \times 10^6 \text{ K} \), and \( \rho_{\text{ph}} = 5 \times 10^{-7} \text{ g cm}^{-3} \). The wind duration is longer than \( t_{\text{flow}} \), so we assume that \( M_w \approx 4\pi r_p^2 \rho_{\text{ph}} \) (see, e.g., Joss & Melia 1987). Paczynski & Proszynski (1986) show that the velocity is nearly constant beyond the photosphere; that is, \( v(r > r_{\text{ph}}) \approx v_{\text{ph}} \approx 10^8 \text{ cm s}^{-1} \), so \( \rho(r > r_{\text{ph}}) \approx M_w/(4\pi r_{\text{ph}}^2) \).

The gas above the photosphere is Compton-heated by the hot photons originating in the photosphere. Since the Compton heating timescale \( t_{\text{c}} \sim kT_{\text{ph}}c^2/E \), \( T_{\text{ph}} \sim 10^7 \text{ s} \) \( \ll t_{\text{flow}} \), where \( E \) is the photon energy and \( F \) is the flux, the gas temperature is nearly constant out to radii well above the photosphere (Joss & Melia 1987; Nobili et al. 1994). For \( r \gg r_{\text{ph}} \), \( T_c \approx t_{\text{flow}} \) and the gas cools adiabatically, although the density in this region is so low that in calculating \( N_{\text{wind}} \) we assume \( T(r > r_{\text{ph}}) = T_{\text{ph}} \). For the more abundant species shown in Figure 13, we obtain column densities in the range \( N_{\text{wind}} \sim 10^{16} \text{–} 10^{21} \text{ cm}^{-2} \).

To determine the EW of the photoionization edge, we integrate over the optical depth above the edge, \( \tau_e = N_{\text{wind}}\sigma_{\text{fs}}(E) \). Thus, if we assume an effectively cold atmosphere, \( \text{EW}_{\text{c}} = \left\{ \int \left[ 1 - \exp(-N_{\text{wind}}\sigma_{\text{fs}}(E)) \right] dE \right\} \) (Bildsten et al. 2003). In Figure 14 we show \( \text{EW}_{\text{c}} \) as a function of \( M_w \) for some of the ejected ashes of models He0.1 and HHe0.1. In calculating the EWs we accounted for the decay of those isotopes whose half-lives are shorter than \( 1 \text{ s} \) and therefore shorter than the wind duration. Thus, \( ^{25}\text{Si} \to ^{24}\text{Mg} \), \( ^{28}\text{S} \to ^{28}\text{Si} \), \( ^{49}\text{Fe} \to ^{49}\text{Cr} \), and \( ^{62}\text{Ge} \to ^{62}\text{Zn} \), all from model HHe0.1. As \( M_w \) increases, \( T_{\text{ph}} \) decreases and the hydrogen-like fraction \( \zeta \) increases. For the heavier ashes (e.g., \( ^{60}\text{Zn} \) and \( ^{62}\text{Zn} \)) \( \zeta \approx 1 \) even at low \( M_w \), so \( N_{\text{wind}}(Z \approx 30) \) and \( \text{EW}_{\text{c}}(Z \approx 30) \) are essentially set by \( f(A, Z) \), as noted above. The
Several observational studies report the possible presence of spectral lines and edges during the RE phase of bursts. Indications of such features were found in Rossi X-Ray Timing Explorer (RXTE) Proportional Counter Array (PCA) observations of GX 17+2 (Kuulkers et al. 2002), 4U 1722−30, 4U 2129+12, XB 1745−25 (Kuulkers et al. 2003), and 4U 1820−30 (Strohmayer & Brown 2002; Kuulkers et al. 2003), in Ginga Large Area Counter (LAC) observations of 4U 2129+12 (van Paradijs et al. 1990), in EXOSAT observations of EXO 1747−21 (Maggi et al. 1989), and Tenma observations of 4U 1636−53 (Waki et al. 1984). In general, these studies report significant residuals in blackbody spectral fits that cannot be explained by spectral hardening alone. Higher resolution observations with Chandra and XMM-Newton are clearly needed.

### 5.4. Spectral Edges from NS Surface

Following the burst peak, the atmosphere cools, the flux becomes sub-Eddington, and the wind turns off. The photosphere, still laced with heavy-element ashes, settles back down to the NS surface. As in the case of spectral edges in the wind outflow, whether these surface ashes can be detected depends on the column density, $N_{\text{surf}}$, of ashes that are not fully ionized. Here $N_{\text{surf}} \approx f(A, Z)(T, \rho)\gamma_{\text{phot}}/m_p$, where $\gamma_{\text{phot}} \approx 1$ g cm$^{-2}$. As the surface temperature $T$ decreases, $N_{\text{surf}}$ increases, although edges cannot be detected once $L = 4\pi R^2 \sigma T^4 \leq L_{\text{acc}}$, where $L_{\text{acc}}$ is the accretion luminosity. The decay timescale of RE bursts (i.e., the time during which $L_{\text{acc}} < L < L_{\text{Edd}}$) is typically around $10s$. The downward drift speed of a nucleus with $Z \sim 30$ in a pure hydrogen atmosphere is $v \approx (1 \text{ cm s}^{-1})T_7^{3/2}/\rho$, so it takes $t_d \approx (1 \text{ s})(y/1 \text{ g cm}^{-2})T_7^{3/2}$ for such a nucleus to fall a column depth $y$ (Bildsten et al. 1992, 2003). Thus, the residence time of ashes in the photosphere also limits the detectability of edges, as will high NS rotation rates (Özel & Psaltis 2003; Chang et al. 2005).

In Figure 15 we plot $E_{\text{w}} = f(1 - \exp(-N_{\text{surf}}\gamma_{\text{phot}}(E)))$ as a function of $L/L_{\text{acc}}$ for the same models as Figure 14. For reference, $L_{\text{Edd}}/L_{\text{acc}} \approx 10$ and 30 for models He0.1 and HHe0.1, respectively. The values of $E_{\text{w}}$ are again within the range accessible by current X-ray telescopes, and a measurement of the NS gravitational redshift may be possible.

### 6. SUMMARY AND CONCLUSIONS

We have shown that during a radius expansion type I X-ray burst the ashes of nuclear burning can be ejected by the burst’s radiative wind. Specifically, we solved for the evolution of the atmosphere’s thermal structure and found that in systems accreting pure helium, such as $4U 1820−30$, and in systems accreting mixed hydrogen and helium at $m \lesssim 0.05m_{\text{Edd}}$, the convective zone reaches sufficiently low pressures that it lies within the wind region located at pressures $\lesssim 1%$ of that at the base. Hence, ashes of burning, mixed throughout the convective zone, can be among the wind ejecta. Previous studies have also found that the convective zone can reach $\lesssim 1%$ of the base pressure for low-$m$ bursts (Joss 1978; Taam 1981; Ayasli & Joss 1982; Hanawa & Sugimoto 1982; Hanawa & Fujimoto 1984; Woosley et al. 2004). However, they focused on numerical simulations of bursts and typically did not resolve the low-pressure zones that convection reaches. Furthermore, such studies only explored a limited range of burst parameter spaces. Our analytic treatment enabled a survey of the dependence of the minimum pressure reached by the convective zone on a variety of burst parameters such as $m$, the composition of accreted matter, and the preburst thermal state of the atmosphere. We have compared some of the results of our analysis, such as the burst rise times, to those of numerical simulations and have found good agreement.

To calculate the nuclear energy generation rate and ash composition, we used an updated full reaction network that contains 394 nuclei and all relevant proton- and $\alpha$-particle-induced reactions and the corresponding inverse processes. We have found that in models where burning initiates in a pure helium layer (i.e., those systems accreting mixed hydrogen and helium at $m \lesssim 0.04m_{\text{Edd}}$ or systems accreting pure helium), protons are...
produced by \((\alpha, p)\) reactions on heavier \(\alpha\)-chain nuclei such as \(^{24}\text{Mg}, \ ^{35}\text{S}, \) and \(^{36}\text{Ar}\). The proton abundances achieved are high enough to enable the bypass of the relatively slow \(^{12}\text{C}(\alpha, \gamma)\) reaction by the much faster \(^{12}\text{C}(p, \gamma)\) reaction sequence. The bypass leads to a burst in energy production that pushes the convective zone to lower pressures and results in a rapid, \(~10^{-3}\) s, rise in the burst light curve (see Fig. 10). The bypass also enhances the production of heavier \(\alpha\)-chain nuclei, resulting in heavier ejected ashes.

For specific burst models, we determined the composition of ejected ashes and calculated the expected column density of hydrogen-like nuclei using models of relativistic radiation-driven winds. We then computed the EW of the photoionization edge for the more abundant hydrogen-like nuclei in the wind. We carried out a similar procedure to determine the EW for those ashes that remain bound to the NS and thus reside in the photosphere after it settles back down to the NS surface. We found EWs in the range 10–1000 eV (see Figs. 14 and 15), which suggests that the edges can be detected with current high-resolution X-ray telescopes. Detecting them would directly probe the nuclear burning. Those edges formed at the surface may also provide a measurement of the NS gravitational redshift and thus help constrain the NS equation of state.

If bursts ignite in the ashes of previous bursts, then some of these processed ashes, which are thought to contain large overabundances of \(\alpha\)-nuclei relative to solar, are also ejected in the wind. We showed that at least \(~1\)% of the \(\alpha\)-nuclei observed in the solar system may originate in X-ray bursts.

We did not account for the effect of a laterally spreading burning front during burst rise nor for the influence of rotation, although these may significantly alter the convective structure and therefore the conditions under which ashes are ejected. Bursting systems have magnetic fields \(B \approx 10^9\) G, so the magnetic energy density, \(B^2/8\pi\), is considerably smaller than even the minimum convective energy density \(\rho V_c^2 \approx g V_c, \text{min} \approx 5 \times 10^{16}\) erg cm\(^{-3}\). Magnetic fields are therefore unlikely to significantly affect the convective evolution.

Although low-\(n\) systems accreting a mix of hydrogen and helium yield X-ray bursts with convective zones that reach low pressures, their burst recurrence times are long and often irregular, which makes them difficult targets to monitor, given a narrow window of observing time. More promising are systems in which the neutron star accretes helium-rich material \((X_{\text{He}} \approx 0.9–1)\) from an evolved companion, such as a cold helium white dwarf. The binary 4U 1820–30 is thought to reside in such a system, as evidenced by its ultracompact nature \((P_{\text{orb}} = 11.4\) minutes; Stella et al. 1987) and the fast rise times, decay times, and \(\alpha\)-values observed during bursts (C03 and references therein). The system also exhibits radius expansion bursts with fairly regular burst recurrence times of only a few hours (see, e.g., Cornelisse et al. 2003; Kuulkers et al. 2003).

Observations of four candidate ultracompact binaries have shown an unusual Ne/O abundance ratio in the absorption along the line of sight, with ratios several times the interstellar medium (ISM) value (Juett et al. 2001; Juett & Chakrabarty 2003). Two of the systems, 4U 1543–624 and 4U 1850–087, have shown variations in the Ne/O ratio in follow-up observations (Juett & Chakrabarty 2003, 2005), which suggests either variations in the ionization state of the Ne and O or variations in the intrinsic abundances. Both possibilities imply the absorption is due to material local to the binaries.

One explanation for the unusual ratios is that the intrinsic abundance of Ne and O is the same as the ISM, but the O is in a higher ionization state than the Ne, leading to an apparent enhancement of the Ne/O ratio. Another possibility is that the degenerate donors in these ultracompact binaries have Ne/O abundances above the solar value. In particular, the donors may be the chemically fractionated cores of C-O-Ne or O-Ne-Mg white dwarfs that have previously crystallized (Schulz et al. 2001; Juett & Chakrabarty 2003; Sidoli et al. 2005), or, in some cases, the central mass of a helium white dwarf. In ‘t Zand et al. (2005) have shown that the latter donor type is consistent with both the properties of the long X-ray burst observed from the ultracompact binary 2S 0918–549 and that system’s enhanced (and constant) Ne/O ratio.

A third possibility we now propose is that the donors in these ultracompact systems are helium white dwarfs, as in 4U 1820–30, and that the accretion of helium-rich matter results in RE bursts and ejection of ashes with highly nonsolar abundances of Ne and O. Radius expansion bursts have indeed been seen in two bursts from 2S 0918–549 (Cornelisse et al. 2002; in ‘t Zand et al. 2005). Although it is only a single realization, the Ne/O ratio by number at \(y = y_{\text{w}}\) for model HeA1 shown in Figure 11 is \(~1.0\), which is comparable to that seen in the four systems. In order for this explanation to work, the ejected ashes must either somehow remain in the environment of the binary or be continuously replenished by periodic RE bursts. The variations in the observed Ne/O ratio in observations of 4U 1543–624 and 4U 1850–087 could just reflect the time since the last RE burst.

Finally, we note that ash ejection may also occur during superbursts that undergo photospheric radius expansion. The amount of mass ejected by superbursts may be much larger than those from ordinary RE bursts, suggesting even higher column densities of ejected ashes. Furthermore, the ejected ashes are likely to be heavier, since the carbon that fuels the superbursts lies underneath the heavy \(\alpha\)-ashes of normal bursts.

We thank A. Heger for sharing calculations that encouraged the start of this study, F.-K. Thielemann for providing the network solver, A. Sakharuk for fitting and implementing the nuclear reaction rates, and E. Kuulkers for bringing to our attention earlier observations. This research was supported by NASA grant NNG05GF69G and NSF grants PHY 99-07949, AST 02-05956, and PHY 02-16783 (Joint Institute for Nuclear Astrophysics). H. S. is also supported by NSF grant PHY 01-10253.

**REFERENCES**

Angulo, C., et al. 1999, Nucl. Phys. A, 656, 3

Audi, G., Bersillon, O., Blachot, J., & Wapstra, A. H. 2003, Nucl. Phys. A, 729, 3

Ayaisi, S., & Joss, P. C. 1992, ApJ, 396, 637

Beardmore, A. P., et al. 2004, MNRAS, 348, 101

Bildsten, L. 1998, in The Many Faces of Neutron Stars, ed. R. Buccheri, J. van Paradijs, & M. A. Alpar (Dordrecht: Kluwer), 419

Bildsten, L., Chang, P., & Paerels, F. 2003, ApJ, 591, L29

Bildsten, L., Salpeter, E. E., & Wasserman, I. 1992, ApJ, 384, 143

Brown, B. A., Clement, R. R. C., Schatz, H., Volya, A., & Richter, W. A. 2002, Phys. Rev. C, 65, 045802

Brown, E. F. 2000, ApJ, 531, 988

Brown, E. F. 2004, ApJ, 614, L57

Brown, E. F., & Bildsten, L. 1998, ApJ, 496, 915

Buchmann, L. 1996, ApJ, 468, L127

Caughlan, G. R., & Fowler, W. A. 1988, At. Data Nucl. Data Tables, 40, 283

Chang, P., Bildsten, L., & Wasserman, I. 2005, ApJ, 629, 998

Cornelisse, R., et al. 2002, A&A, 392, 885

———. 2003, A&A, 405, 1033

Cumming, A. 2003, ApJ, 595, 1077 (C03)

Cumming, A., & Bildsten, L. 2000, ApJ, 544, 453

Ebisuzaki, T., & Nakamura, N. 1988, ApJ, 328, 251
Fisker, J. L., Barnard, V., Górres, J., Langanke, K., Martínez-Pinedo, G., & Wiescher, M. C. 2001, At. Data Nucl. Data Tables, 79, 241
Fujimoto, M. Y., Hanawa, T., & Miyaji, S. 1981, ApJ, 247, 267
Fuller, G. M., Fowler, W. A., & Newman, M. J. 1982, ApJS, 48, 279
Fushiki, I., & Lamb, D. Q. 1987, ApJ, 323, L55
Fynbo, H. O. U., et al. 2005, Nature, 433, 136
Galloway, D. K., Cumming, A., Kuulkers, E., Bildsten, L., Chakrabarty, D., & Rothschild, R. E. 2004, ApJ, 601, 466
Hahn, K. I., et al. 1996, Phys. Rev. C, 54, 1999
Hanawa, T., & Fujimoto, M. Y. 1981, ApJ, 247, 267
Hanawa, T., & Sugimoto, D. 1982, PASJ, 34, 1
Hansen, C. J., & Kawaler, S. D. 1994, Stellar Interiors: Physical Principles, Structure, and Evolution (Berlin: Springer)
Hashimoto, M.-A., Hanawa, T., & Sugimoto, D. 1983, PASJ, 35, 491
Herndl, H., Górres, J., Wiescher, M., Brown, B. A., & van Wormer, L. 1995, Phys. Rev. C, 52, 1078
Iliadis, C., D’Auria, J. M., Starrfield, S., Thompson, W. J., & Wiescher, M. 2001, ApJ, 560, L59
in ’t Zand, J. J. M., Cumming, A., van der Sluys, M. V., Verbunt, F., & Pols, O. R. 2005, A&A, 441, 675
Joss, P. C. 1977, Nature, 270, 310
———. 1978, ApJ, 225, L123
Joss, P. C., & Melia, F. 1987, ApJ, 312, 700
Juett, A. M., & Chakrabarty, D. 2003, ApJ, 599, 498
———. 2005, ApJ, 627, 926
Juett, A. M., Psaltis, D., & Chakrabarty, D. 2001, ApJ, 560, L59
Klypin, A., Zhao, H., & Somerville, R. S. 2002, ApJ, 573, 597
Kuulkers, E., den Hartog, P. R., in ’t Zand, J. J. M., Verbunt, F. W. M., Harris, W. E., & Cocchi, M. 2003, A&A, 399, 663
Kuulkers, E., Homan, J., van der Klis, M., Lewin, W. H. G., & Méndez, M. 2002, A&A, 382, 947
Lewin, W. H. G., van Paradijs, J., & Taam, R. E. 1995, in X-Ray Binaries, ed. W. H. G. Lewin, J. van Paradijs, & E. P. J. van den Heuvel (Cambridge: Cambridge Univ. Press), 175
Maguire, E., Lewin, W. H. G., van Paradijs, J., Tan, J., Penninx, W., & Damen, E. 1989, MNRAS, 237, 729
Nobili, L., Tsurulla, R., & Lapidus, I. 1994, ApJ, 433, 276
Ozel, F., & Psaltis, D. 2003, ApJ, 582, L31
Paczyński, B. 1983, ApJ, 267, 315
Paczyński, B., & Proszynski, M. 1986, ApJ, 302, 519
Rauscher, T., & Thielemann, F.-K. 2000, At. Data Nucl. Data Tables, 75, 1
Rybicki, G. B., & Lightman, A. P. 1979, Radiative Processes in Astrophysics (New York: Wiley)
Schatz, H., Bildsten, L., Cumming, A., & Wiescher, M. 1999, ApJ, 524, 1014
Schatz, H., et al. 1998, Phys. Rep., 294, 167
———. 2001, Phys. Rev. Lett., 86, 3471
Schulz, N. S., Chakrabarty, D., Marshall, H. L., Canizares, C. R., Lee, J. C., & Houck, J. 2001, ApJ, 563, 941
Sidoli, L., La Palombara, N., Oosterbroek, T., & Parmar, A. N. 2005, A&A, 443, 223
Spitkovsky, A., Levin, Y., & Ushomirsky, G. 2002, ApJ, 566, 1018
Stella, L., Friedl, W., & White, N. E. 1987, ApJ, 312, L17
Strohmayer, T., & Bildsten, L. 2006, in Compact Stellar X-Ray Sources, ed. W. H. G. Lewin & M. van der Klis (Cambridge: Cambridge Univ. Press), in press (astro-ph/0301544)
Strohmayer, T. E., & Brown, E. F. 2002, ApJ, 566, 1045
Sugimoto, D., Ebisuzaki, T., & Hanawa, T. 1984, PASJ, 36, 839
Taam, R. E. 1980, ApJ, 241, 358
———. 1981, ApJ, 247, 257
Timmes, F. X., Hoffman, R. D., & Woosley, S. E. 2000, ApJS, 129, 377
van Paradijs, J., Dotani, T., Tanaka, Y., & Tsuru, T. 1990, PASJ, 42, 633
Waki, I., et al. 1984, PASJ, 36, 819
Wallace, R. K., & Woosley, S. E. 1981, ApJS, 45, 389
Wallace, R. K., Woosley, S. E., & Weaver, T. A. 1982, ApJ, 258, 696
Wallnerstein, G., et al. 1997, Rev. Mod. Phys., 69, 995
Woosley, S. E., & Weaver, T. A. 1984, in AIP Conf. Proc. 115, High Energy Transients in Astrophysics, ed. S. E. Woosley (New York: AIP), 273
Woosley, S. E., et al. 2004, ApJS, 151, 75
Yahel, R. Z., Brinkmann, W., & Braun, A. 1984, A&A, 139, 359