DYNAMICAL FRAGMENTATION OF THE T PYXIDIS NOVA SHELL DURING RECURRENT ERUPTIONS

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ABSTRACT

Hubble Space Telescope images of the ejecta surrounding the nova T Pyxidis resolve the emission into more than 2000 bright knots. We simulate the dynamical evolution of the ejecta from T Pyxidis during its multiple eruptions over the last 150 years using the adaptive mesh refinement code Ramses. We demonstrate that the observed knots are the result of Richtmyer–Meshkov gas dynamical instabilities (the equivalent of Rayleigh–Taylor instabilities in an accelerated medium). These instabilities are caused by the overrunning of the ejecta from the classical nova of 1866 by fast-moving ejecta from the six subsequent recurrent nova outbursts. Magnetic fields may play a role in determining knot scale and preventing their conductive evaporation. The model correctly predicts the observed expansion and dimming of the T Pyx ejecta as well as the knotty morphology. The model also predicts that deeper, high-resolution imagery will show filamentary structure connecting the knots. We show reprocessed Hubble Space Telescope imagery that shows the first hints of such a structure.

Key words: hydrodynamics – instabilities – novae, cataclysmic variables – shock waves

Online-only material: color figures

1. INTRODUCTION

Recurrent novae (RNe) are cataclysmic variables that display multiple eruptions on a timescale of decades (Warner 1995). The same mechanism powers both RNe and classical novae: thermonuclear runaways in the hydrogen-rich envelopes of white dwarfs (WDs). The erupting envelopes have been accreted by the WDs from the non-degenerate companions—red dwarfs or red giants—in these binary systems. The observed short recurrence timescales of RNe are only possible for massive WDs (∼1.2–1.4 M⊙) with high-mass accretion rates (∼10−8 M⊙ yr−1; Starrfield et al. 1985; Yaron et al. 2005). Such high accretion rates suggest that RNe might be part of the population of Type Ia supernova progenitors (Kovetz & Prialnik 1994; della Valle & Livio 1996). Ten RNe are known in the Milky Way; Schaefer (2010) gives an extensive review of their observational properties.

The first RN discovered (and the prototype of such systems) was T Pyxidis (T Pyx), detected by H. Leavitt in 1913 (Pickering 1913). With six identical (Mayall 1967; Schaefer 2010) recorded eruptions (1890, 1902, 1920, 1944, 1966, and 2011) to its credit, T Pyx remains the RN with the longest track record, eruptions (1890, 1902, 1920, 1944, 1966, and 2011) to its.

The ejecta of T Pyx, displaying a circular ring with a radius of 5′′, were first detected by Duerbeck & Seitter (1979). More sensitive observations have since revealed fainter outer material, reaching a radius of 6′5″ (Schaefer et al. 2010). This morphology was initially interpreted as an edge-brightened shell. Williams (1982) used spectroscopy of the ejecta to show that the material has roughly solar abundances and displays emission lines similar to those seen in planetary nebulae. Shara et al. (1989) demonstrated that the apparent shell was expanding too slowly to have originated in the 1944 or 1966 eruptions, and that it was at least twice as large as previously suspected. Contini & Prialnik (1997) showed that the spectral line fluxes can be successfully explained by shock heating when the fast ejecta from one nova eruption run into the slower ejecta from an earlier eruption. Hubble Space Telescope (HST) images from 1994 to 1995 resolved the previously observed ring into over 2000 knots, some of which fade or brighten significantly on a timescale of order one year (Shara et al. 1997). Expansion of the knots was not detected, though the time baseline used (1.7 yr) was quite short.

Thirteen years later, Schaefer et al. (2010) re-observed T Pyx with HST, detecting knot expansion velocities of 500–700 km s−1. The observed fractional expansion of the knots is constant (hence there is little deceleration from the interstellar medium). Dozens of the brightest knots are observed to have survived intact during these 13 years. The HST observations constrain the knots to have originated in an explosion close to the year 1866, and to possess a total mass of ∼1045 M⊙. The distance to T Pyx is poorly constrained. Schaefer et al. (2010) estimate it to be D = 3.5 ± 1 kpc. The 1866 event must have been a classical nova eruption, preceded by a low accretion epoch lasting of order 1 Myr (Schaefer et al. 2010). The 1866 event would have triggered a supersoft X-ray source on the T Pyx WD, driving a much higher mass transfer rate after that event (Knigge et al. 2000). Mass transfer has been steadily declining in the T Pyx binary (Schaefer 2005), and in 2009 was a factor of 30 less than in 1890. No more RN outbursts are expected before T Pyx enters a state of hibernation, expected to last for 2–3 Myr. There is consensus that the WD in T Pyx is unlikely to exceed the Chandrasekhar mass and become a Type Ia supernova (Selvelli et al. 2008; Schaefer et al. 2010).
One critical test of the scenario in which the ejecta are shaped by the history of a nova followed by six RN eruptions is to reproduce the morphology of the ejecta. We hypothesize that the observed knots are produced by Richtmyer–Meshkov instabilities excited in the thin, swept-up shell of circumstellar gas surrounding the nova ejecta when subsequent explosions accelerate it. This instability is the equivalent of the Rayleigh–Taylor instability that occurs when acceleration rather than gravity drives overturn of a dense fluid supporting a more rarefied fluid (Richtmyer 1960; Meshkov 1969). Outward acceleration of the dense shell by more rarefied ejecta results in an effective gravity pointing inward, from the shell into the ejecta, and thus results in overturn and fragmentation of the dense shell.

We test our hypothesis using three-dimensional, gas dynamical simulations that include radiative cooling, though not thermal conduction, and are sufficiently well resolved numerically to follow the instabilities that can cause the observed fragmentation. García-Segura et al. (2004) previously reported a preliminary two-dimensional model of this problem. In Section 2 we outline the simulation techniques, and describe the initial velocities, time intervals, and ejected masses we assumed. These were taken from the observations of T Pyx by Adams & Joy (1920), Joy (1945), Catchpole (1969), and Schaefer et al. (2010). Our numerical results are presented and compared with the HST observations of T Pyx in Section 3. Our results and conclusions are summarized in Section 4.

2. SIMULATIONS

2.1. Code Used

We compute a fully three-dimensional model of the original T Pyx classical nova of 1866, and the subsequent RNe up to, though not including, the 2011 event, using a slightly modified version of the gas dynamical code RAMSES, version 3.0 (Teyssier 2002). This code uses adaptive mesh refinement, with a tree-based structure that allows recursive grid refinement on a cell-by-cell basis. The hydrodynamical solver is based on the second-order Godunov method. The conservative variables $\rho, \rho v, \rho e$, where $\rho$ is the mass density, $v$ is the velocity, and $e$ is the specific energy density, are taken to be piecewise constant over the mesh cells at each time step. The time solution is determined by the exact solution of the Riemann problem at the intercell boundaries (use of the exact Riemann solver is one of the choices in the default code).

We refined the grid wherever $\nabla \cdot v > 0.1|v|$ across a single cell, so long as $|v| > 10 \text{ cm s}^{-1}$. We used a monotonized central (MonCen; Toro 1997) slope limiter in the interpolation scheme for newly refined cells. All other algorithmic parameters were set to default values.

2.2. Problem Setup

A cubical grid with an edge length of $6.0 \times 10^{17} \text{ cm}$ was used for our simulations. We assumed reflecting boundary conditions at the surfaces of the cube, which never participated in the solution, as our runs always terminated before interacting with the boundary.

We implemented repeated instantaneous energy inputs to simulate nova and RN explosions. A driver routine was added that set up nova explosions at given times, which were then allowed to evolve over time without any further changes until the next explosion. The nova explosions were located at the center of the cubical grid. The spherical source was assumed to have a radius $r_s = 1.25 \times 10^{16} \text{ cm}$, which is about 11 pixels on a $512^3$ cube, in order for the adaptive mesh refinement code.
Figure 2. Log density distribution for six epochs as shown on the figure, each just before the next RN is about to explode, except for the final epoch which is 1995. Shown are two-dimensional cuts through the log density field in the $xz$-plane, showing the imposed asymmetry. The figure demonstrates the fragmentation of the nova shell by instabilities as it is repeatedly accelerated. These results are from our canonical $512^3$ zone simulation with 50% background and 30% source noise.

Figure 3. (a) Square root and (b) linear display of the emission measure of the ejecta in 2007, as viewed along the direction of observation (the $z$ axis of the simulation). The T Pyx source is at the center of the box. The color bars give the values in cgs units. The square-root scaling emphasizes faint structure, while the linear scaling reproduces the scaling used in Figure 1.

to be able to resolve the initial evolution of the spherical blast wave. The density and the velocity of the source material were derived from the observed mass and velocity of the ejecta as we explain below.

To decide on the background density, we determined the density that would be consistent with the size of the observed shell and the energy of the 1866 nova explosion. Inverting the Sedov (1959) similarity solution for the radius of a blast wave to solve for the density,

$$\rho = \frac{2.2E\tau^2}{r^5},$$

where $r$ is the radius of the expanding shell, $E = 7.46 \times 10^{53}$ g cm$^{-2}$ s$^{-2}$ is the kinetic energy of the 1866 nova, and the age of the shell $\tau = 129$ yr. We take the distance to T Pyx to be 4.5 kpc (Shore et al. 2011), although this is a value uncertain to a factor of two. If we then take the angular radius of the shell to be 6$'$5 (Schaefer et al. 2010), we can use Equation (Sedov 1959) to derive a background density $\rho_0 = 1.689 \times 10^{-25}$ g cm$^{-3}$.

The background pressure was set to $1.38 \times 10^{-12}$ erg cm$^{-3}$, yielding a temperature of $6.8 \times 10^4$ K. This rather high background temperature was set by accident (initially we had set the pressure to be appropriate for $n = 1$ cm$^{-3}$ pure hydrogen gas at $T = 10^4$ K, and neglected to update when we moved to the lower density), but that does not appear to markedly influence the shell dynamics of the nova that we study, as the internal pressure of the nova blast wave is orders of magnitude higher throughout its evolution. It could certainly be consistent with a location of the nova within a pressurized superbubble.

In order to trigger the shell instabilities, in our initial simulation with $512^3$ we set up random zone-to-zone noise for background at a level of 50% and zone-to-zone source noise at 30% for the 1866 nova. For all other novae, the background was unchanged, while the source noise was 50% of the source density. Later, we simulated the $512^3$ setup again with only the source and background noise changed to study its effect.

We assumed the gas to be monatomic, with adiabatic index $\gamma = 5/3$. We used the equilibrium ionization cooling
The density of the explosion source $\rho_s$ was found by setting the value of the kinetic energy $E$ and the expansion velocity $u_0$ of each nova outburst, so that

$$\rho_s = \left( \frac{2E(\Delta x)^3}{\Sigma v_{ijk}^2} \right).$$

where $\Delta x$ is the length of a single zone, $v_{i,j,k}$ is the velocity in the zone with index $(i, j, k)$, and the sum runs over all the zones in the source and transition regions. The kinetic energy $E$ is derived from the mass and velocity measured by the observations.

The 1866 eruption was a classical nova, occurring after several hundred thousand years of very low mass accretion rate. Its observed ejecta mass was $\sim3 \times 10^{-5} M_\odot$, and velocity was $u_0 = 5 \times 10^2$ km s$^{-1}$ (Schafer et al. 2010). The post-1866 eruptions were RN eruptions, driven by the initially very high, but declining, observed mass transfer rate, which started at a value of roughly $10^{-8} M_\odot$ yr$^{-1}$. As the rate declined, the interval between RN eruptions extended. The post-1866 events eject much less mass than the 1866 classical nova. We use parameters derived from the observed RNs, with ejecta masses of $1 \times 10^{-7} M_\odot$ and ejection velocities of $2 \times 10^3$ km s$^{-1}$.

The ejecta from each of the eruptions was allowed to expand in the simulations for 145 years after the first (classical nova) eruption, reaching the year 2011 to enable a comparison with the HST observations. Our simulation stops just before the 2011 eruption, whose effects on the ejecta will not become significant for several decades.

The $T$ Pyx ejecta are observed to be elongated in the polar directions (Shara et al. 1997). We assume that this asymmetry is caused by an asymmetric nova explosion or by interaction of the blast with the accretion disk within several stellar radii of the white dwarf (Walder et al. 2008; Drake & Orlando 2010). In either case, the region where the elongation occurs falls within our source region, so we simply give the ejecta velocity a dependence on the polar angle

$$u_1(\phi) = u_0(1.0 + 0.1 \cos \phi).$$

The subsequent explosions were assumed to be spherical. We initialize each zone in the source region with the Cartesian components of the derived velocity, $u_1$, or for the subsequent explosions, $u_0$.

The source region was extended by 50% as a transition region between the high-velocity material at the edge of the source region and the material at rest outside. We found this necessary to maintain the stability of the Riemann solver. To do this, we multiplied velocities by a radial profile function, so that $v_{ijk} = u_{0,1} \xi(r)$ as appropriate. We also chose the profile to dip to zero at the center of the source to avoid Riemann solver failures caused by a strong rarefaction wave at the center of the source region, and have a smooth transition to zero across the outer transition region. This also qualitatively reproduces the near linear decline in velocity toward the center in the actual

$$\xi(r) = \begin{cases} 
(tanh[10(r/r_s - 0.5)] + 1.0)/2, & \text{if } r < r_s \\
(tanh[10(r/r_s - 0.5)] + 1.0)(1.5r_s/r - 1), & \text{if } r_s < r < 1.5r_s \\
0, & \text{if } r > 1.5r_s.
\end{cases}$$

The density within the transition region was set to the source density $\rho_s$. We do not expect the details of the source region to influence the solution after the swept-up material exceeds the ejecta in mass by a factor of a few.

We performed a resolution study ranging from $256^3$ to $1024^3$ zones. Most results we display here are shown for either our canonical resolution of $512^3$ zones, or for our highest resolution, and we discuss the convergence properties of our results. We only ran one octant of the $1024^3$ zone grid (with reflecting boundary conditions on the inner faces) because of computational expense. We used an adaptive mesh with a base grid resolution of $128^3$ zones at seven levels of refinement and maximum refinement level of $g = 8, 9, \text{ or } 10$ for the three resolutions that we studied. The number of zones along a side of the cubic grid is given by $2^g$. 

![Figure 4](image-url)

Figure 4. Square root of the emission measure for year 2007, as viewed along the x- and y-axes (perpendicular to the line of observation). The color bar gives the values in cgs units. The morphology strongly resembles that in Figure 3(a), and so seems not to be sensitive to the displayed line of sight.

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3. OBSERVATIONS

We extracted the deepest set of HST images of T Pyx, taken with the Wide Field and Planetary Camera 2 through the F658N filter, which isolates the strongest optical [N II] emission line. Details of the images are given in Table 2 of Shara et al. (1997). These were drizzled (Fruchter & Hook 2002) to produce Figure 1. Figure 1 reveals faint, emitting blobs as much as 0.5 mag fainter than the faintest blobs seen in the (undrizzled) F658N image that is shown in Figure 2 of Shara et al. (1997). Our previous impression of many knots (Shara et al. 1997) is reinforced.
4. RESULTS

Figure 2 displays polar cuts through the computed log density distribution of the ejecta of the first six eruptions, seen just before the next RN eruption, except for the last frame, which shows the distribution in the year 1995. The figure shows the model prediction that the knots expand away from the center as time progresses, as recently demonstrated observationally by Schaefer et al. (2010). The outer shell (i.e., the outermost part of the ejecta) appears to be the remnants of the original 1866 nova eruption, moving slowly away from the central star.

5. COMPARISON TO OBSERVATION

We use the discretized emission measure from the simulations $EM = \Sigma_{ijk} n_{ijk} \Delta x$ to compare to the HST observations, where $n_{ijk} = \rho_{ijk}/\mu$ is the number density in the zone with index $(i, j, k)$, and we use a value of the mean mass per particle $\mu = (14/21)m_H$ to account for ionized hydrogen and singly ionized helium. The emission measure gives a reasonable estimate of the [N II] emission due to collisional excitation as the cool, dense, recombining gas completely dominates the sum, because of the nonlinear dependence on density. The hot $10^6$ K gas on the grid is uniformly at low density and thus does not contribute significantly to the emission measure, as we discuss further when we show the temperature distribution of our models below.

Emission measure maps were produced by interpolating the adaptive grid onto a uniform grid at the finest level and then summing. Comparison of Figures 3(a) and 4 suggests that the overall shape of the ejecta distribution is neither an artifact of nor sensitive to the direction of observation despite the asymmetry of expansion velocity observed.

We first compare observations with our canonical simulations done on a $512^3$ grid with initial background and source noise levels of 50% and 30% respectively. We then discuss the effects of changing spatial resolution and initial noise levels.

Figures 3(a) and 4 display images from the year 2007 model of the square root of the simulated emission measure along all three axes, where $n$ is the number density in the model. The same filamentary structures seen in the density distributions of Figure 2 are clearly visible in these emission measure images. They are emphasized in Figure 3(a) by the square-root scaling adopted. Figure 3(b) shows a clipped display of the emission measure to again emphasize the weaker emission. In this case, we take the top of the gray scale to be $4 \times 10^{21}$ cm$^{-5}$. The fingers that are so prominent in the two-dimensional cuts shown in Figure 2 are not seen in the observations shown in Figure 1. We predict that deeper images of the T Pyx ejecta will show such elongated features; however, we also note that the resolution study that we describe in the next subsection shows that the length and strength of these features have not yet converged, with higher resolution yielding shorter and more numerous features.

Beyond morphology, another comparison that we can make between the model and the observations is the radial profile of azimuthally averaged surface brightness. This comparison is complicated because the distance to T Pyx is only known to within a factor of two, so neither the physical scale nor the conversion between luminosity and observed flux can be well determined. In Figure 5, we assumed a distance of 4.5 kpc (as assumed in the background density calculation above) in order to set the physical scale. As the conversion between H$\alpha$ surface brightness in the simulated observations and the actual [N II] surface brightness is also uncertain, we then simply scaled the model profile vertically to fit the intensity at the first complete peak of the observed profile in our highest resolution model, to allow comparison of the measured, azimuthally averaged radial surface brightness distribution observed in 1995 with the model prediction. There is general, overall agreement in the trend and shapes of the curves. The $512^3$ zone model has somewhat fewer peaks and troughs, but the $1024^3$ zone model appears more similar to the observations, though not reproducing the exact observed radial distribution.

The only HST observations of T Pyx available are for the epochs 1994, 1995, and 2007; using these data to directly compare the positions of selected knots, Schaefer et al. (2010) were able to demonstrate the knots’ expansion. Our models succeed in demonstrating the same expansion: Figure 6 is a prediction (based on Figure 3) of the azimuthally averaged radial surface brightness profile of T Pyx expected in 1995, 2007, and 2011. The simulation predicts both a dimming and an expansion of the ejecta, in good agreement with the observations (Schaefer et al. 2010).

Our model allows us to predict the changes in the nebula over the time span covered by the observations. In Figure 7, we show the predicted behavior of the square root of emission measure between 1995 and 2007. Slow expansion of the ejecta is predicted, while the overall shape remains very similar.
5.1. Resolution Studies

Thousands of knots were found in the T Pyx ejecta by Shara et al. (1997), while only around a hundred are seen in Figure 3. This can be shown to be due to the limited numerical resolution we used in our simulation. As a result, increasingly fine numerical grids, with decreasingly low numerical viscosity, resolve increasingly fine wavelengths. We demonstrate this effect through a comparison in Figure 8 which demonstrates how the fragmentation increases with the grid resolution for 256\(^3\), 512\(^3\), and 1024\(^3\) zone grids. Only one octant of the 1024\(^3\) zone grid was run because of the large computational expense (that is, we ran a 512\(^3\) grid covering only one octant of the problem).

Figure 9 shows a comparison of square-root emission measure for the 512\(^3\) and 1024\(^3\) cases in the \(y\)-direction. This shows many more knots for the higher spatial resolution case, which approaches (though does not yet reach) the observed level of fragmentation.

5.2. Effect of Source and Background Noise

To investigate how our assumption of source and background noise can affect the solution, we repeated the same problem with only these parameters changed. The background noise was reduced to 5\% (from 50\%) and source noise was reduced to 10\% from 30\%. Figure 10 shows the square root of emission measure in the \(y\)-direction for year 2007 for the two cases. The morphology and magnitude of the emission appear almost identical, although the model with lower noise can just be discerned to be slightly more symmetric.

5.3. X-Ray Emission

While the nova ejecta sweep up the dense shell, a strong reverse shock travels back through the much lower density ejecta. This reverse shock heats the ejecta to temperatures well over 10\(^8\) K, as can be seen in Figure 11, which shows the temperature distribution of the ejecta at the end of 2007. The relatively low density of the shocked gas results in less...
efficient radiative cooling, so that the gas remains at these high temperatures. (Because of these low densities, these regions also do not contribute to the emission measures evaluated above for optical emission.)

This hot gas should emit soft X-rays. To calculate their luminosity and spectrum, we used the 1993 October version of the Raymond & Smith (1977) code, with the default set of cosmic abundances from Allen (1973). We calculated the emission over the energy range 100 eV–10 keV. We included no absorption from neutral hydrogen, no enrichment of metals in ejecta, and no contribution from the central object or from photoionization. The total luminosity over this waveband at the end of 2007 was calculated to be $2.48 \times 10^{39} \text{erg s}^{-1}$. Figure 12(a) shows an image of the calculated X-ray emission for the $512^3$ zone grid, while Figure 12(b) shows the $1024^3$ grid. Compared to Figure 12(a) this shows more detailed structure, similar to the optical image.

We note that a great deal of the extended luminosity actually comes from a small number of hot spots that may well represent numerical artifacts caused by diffusion from cold, dense knots into immediately adjacent hot gas. We discuss diffusion further below in Section 6.2, where we argue that the luminosities derived here are likely upper limits to the actual values.

Although the total luminosity is uncertain because of these issues of diffusion, it is orders of magnitude less than the value found by Balman (2010) of $6 \times 10^{32} \text{erg s}^{-1}$, suggesting that the emission they report comes from the central object, not from extended emission from the ejecta. Montez et al. (2012) report only emission from the central object in the energy range 0.5–8 keV. We find an order of magnitude less emission in this harder waveband, because of the relatively low effective emission temperature. Although they do not give an explicit lower limit to the extended emission they can measure, it is likely higher than our predicted value.

We also use the same technique to predict the X-ray spectrum over this waveband at the same time, as shown in Figure 13. We compare this spectrum to a blackbody spectrum to estimate how well or poorly a single temperature model reproduces our spectrum. In Figure 13 we overplot a blackbody spectrum with an arbitrary, so the shape is the actual point of comparison rather than the absolute normalization. The comparison reveals that the actual non-isothermal temperature distribution shown in Figure 11 produces a substantial soft X-ray excess compared to a single-temperature model.

6. MAGNETIC FIELDS

Although our simulations do not include magnetic fields, there are two reasons to think that they must actually play a role in determining the properties of the observed knots in T Pyx and other nova shells. First, magnetic fields probably set a length scale for fragmentation, and second, they may suppress thermal conduction sufficiently for the knots to survive for decades as observed.

6.1. Fragmentation Scale

The Richtmeyr–Meshkov instability that forms the knots has no intrinsic minimum wavelength, similar to the Rayleigh–Taylor instability. The only mechanism that limits fragmentation is thus viscosity. If the scale at which physical viscosity dominates could be resolved, then its inclusion might limit this process to the correct physical result. The effective resolution of our $1024^3$ zone simulation is $6 \times 10^{14} \text{cm}$. The physical viscosity will dominate when the Reynolds number $Re = L V / \eta \sim 1$, where $\eta$ is the kinematic viscosity, and $L$ and $V$ are the characteristic length and velocity scales. For dilute gas, the kinematic viscosity is (e.g., Zwicky 1941)

$$\eta \sim (6 \times 10^{16} \text{cm}^2 \text{s}^{-1}) \left( \frac{T}{3000 \text{K}} \right)^{1/2} \left( \frac{\rho}{10^{-22} \text{g cm}^{-3}} \right),$$

where we have used typical values for the shell density and temperature for scaling. The requirement that $Re \sim 1$ then gives a length scale of

$$L \sim \eta / V = (6 \times 10^{11} \text{cm}) \left( \frac{T}{3000 \text{K}} \right)^{1/2} \left( \frac{\rho}{10^{-22} \text{g cm}^{-3}} \right) \times \left( \frac{V}{1 \text{km s}^{-1}} \right)^{-1},$$

where we have roughly scaled the velocity by the sound speed in the cold gas. This is clearly far shorter than we can numerically resolve, and indeed far shorter than the scale at which the knots...
are physically observed, which Schaefer et al. (2010) suggest is about 0.2, or $10^{16}$ cm if we assume a distance to T Pyx of $D = 3.5$ kpc.

This suggests that some other process is actually limiting fragmentation, with the most obvious candidate being magnetic fields in the swept-up shell. Chandrasekhar (1961) uses linear analysis to demonstrate that a magnetic field suppresses growth of the magnetized Rayleigh–Taylor instability at scales

$$L < \frac{B^2}{g(\rho_1 - \rho_0)},$$

where $B$ is the field strength, $g$ is the gravitational acceleration, and $\rho_0$ and $\rho_1$ are the mass densities of the light and heavy fluids. The nonlinear development of the magnetized Rayleigh–Taylor instability has been studied, for example, by Stone & Gardiner (2007). Magnetized Richtmyer–Meshkov instability would be expected to behave essentially identically. We can check this explanation for plausibility by making order of magnitude estimates of the fragmentation scale, shell density, and acceleration. As above, we take $L = 10^{16}$ cm, while the acceleration acts over less than a month, so that we can estimate $g = (100 \text{ km s}^{-1})/(10^8 \text{ s}) = 0.1 \text{ cm s}^{-2}$. The shell density reaches at least $10^{-21} \text{ g cm}^{-3}$ in our models, in the absence of magnetic field, although the field, if present, will act to limit its compression, so this is not a fully self-consistent estimate. The density of the rarefied gas can be neglected for our purposes. These estimates would then yield a required magnetic field strength of

$$B \sim (Lg(\rho_1 - \rho_0))^{1/2} \approx (1 \text{ mG}) \left( \frac{L}{10^{16} \text{ cm}} \right) \left( \frac{g}{0.1 \text{ cm s}^{-2}} \right) \times \left( \frac{\rho_1}{10^{-21} \text{ g cm}^{-3}} \right).$$

This is easily consistent with amplification by compression within the shell of swept-up interstellar magnetic field from the ambient value of 3–5 $\mu$G, given that the density has increased by well over two orders of magnitude from the ambient medium to the cold, compressed shell.
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Figure 12. Predicted X-ray flux in erg cm$^{-2}$ s$^{-1}$ at the end of year 2007, (a) for 512$^3$ grid, and (b) a single octant of the 1024$^3$ grid. Shown is the line-of-sight (xy-plane) X-ray flux distribution at the source. The color scale at the top indicates flux values. These are likely upper limits as discussed in the text.

Figure 13. Predicted X-ray spectrum in units of erg s$^{-1}$ eV$^{-1}$, calculated from the model temperature distribution using the Raymond & Smith (1977) code for the X-ray spectrum of a hot plasma. A blackbody spectrum with a temperature of $T = 4 \times 10^6$ K and an area of 10$^3$ zones (dashed) is overlaid for comparison, demonstrating that a single temperature fit will not capture the behavior of this spectrum.

6.2. Thermal Conduction

Hot electrons can travel significant distances across temperature gradients, diffusing heat from hot to cooler gas. We can estimate the importance of thermal conduction in the absence of magnetic fields (Ferrara & Shchekinov 1993) by comparing the length and timescales of the knots $\ell$ and $t_d$, where $c_s$ is the internal sound speed of the cold knots, to the dimensionless Field length

$$\ell_F^2 = \kappa T/n^2 \Lambda,$$

and the conduction timescale

$$t_c = P \ell^2/\kappa T(n - 1),$$

where $P = nkT$ is the pressure, $T$ is the external temperature, $n$ is the number density, $k$ is the Boltzmann constant, $\gamma$ is the adiabatic index, $\kappa = 6.7 \times 10^{-6} T^{5/2}$ (in cgs units) is the electron thermal conductivity, and $\Lambda = 10^{-22} (T/10^6 \text{ K})^{-0.7}$ erg cm$^{-3}$ s$^{-1}$ is the cooling coefficient, such that $n^2 \Lambda$ is the volume cooling rate. We use the functional approximation to $\Lambda$ good from 10$^5$ to 10$^7$ K given by Mac Low & McCray (1988). We can also consider the cooling timescale

$$t_c = P/\kappa T(n - 1),$$

for a complete picture of the thermodynamics.

We take the material immediately around the dense knots to give the relevant values of density $n$ and temperature $T$. (This is the material within the fingers but outside the dense, cold knots.) As can be seen in Figure 6, typical values for this material are $n = 10$ cm$^{-3}$ and $T = 10^6$ K. We also take the typical knot scale to be $\ell = 5 \times 10^{15}$ cm and internal sound speed to be 30 km s$^{-1}$. With these assumptions, we find that

$$t_d = 53 \text{ yr}; \quad t_c = 0.24 \text{ yr}; \quad t_c = 5100 \text{ yr};$$

and the Field length $\ell_F = 8.2 \times 10^{17}$ cm, larger than our entire simulation box. Although all these values could easily vary by a factor of a few, the conclusion that the knots, with $t_c \ll t_d \ll t_r$, would lie in the regime of rapid evaporation (Ferrara & Shchekinov 1993) in the pure hydrodynamical case is inescapable. This contradicts the observed persistence of single knots over a decade.

However, magnetic fields strongly suppress thermal conduction across them, with the ratio of perpendicular to parallel thermal conduction (Spitzer 1962):

$$\kappa_{lr}/\kappa_{\parallel} = (\omega_{ci} t_d)^{-2} = 2.3 \times 10^{-12} B_{-6}^{-2},$$

where $B_{-6} = B/(1 \mu \text{G})$, $\omega_{ci} = 9.58 \times 10^3 (Z/\mu) B$ s$^{-1}$ is the ion cyclotron frequency (Huba 2011), and $t_d = 11.4 T^{3/2} (n \ln \Lambda)$ is the self-collision time of the ions (Spitzer 1962). The Coulomb
logarithm \( \ln \Lambda = 23 \) for our case (Huba 2011), and we take the ionic charge \( Z = 1 \). Thus, if the strong fields suggested to be present by the arguments in the previous subsection are indeed present, they will virtually prevent thermal conduction across them. With a plausible geometry in which the field threads the swept-up shell, we might expect the fields to wrap around the clumps, protecting them against evaporation by the surrounding gas. This possibility will ultimately require simulation to demonstrate, though.

Although magnetic shielding will ensure the survival of the knots, it does call into question the reliability of our models of X-ray emission, as the density and temperature gradients at the edges of the knots are entirely determined by numerical diffusion rather than the physical effects we have discussed here. If thermal conduction is suppressed as strongly as we expect from Equation (13), though, then numerical diffusion represents a strong upper limit on the width of the interfaces, suggesting that our simulated X-ray luminosity represents an upper limit on the actual values.

7. SUMMARY

We have carried out fully three-dimensional, gas dynamical simulations of the evolution of the ejecta of T Pyx, starting with its classical nova outburst in 1866 and continuing up to the time of the latest HST observations in 2007. We are able to predict not only the observed expansion of the ejecta, but also the extremely knotty morphology. Our simulations demonstrate that the knots form when ejecta from the later outbursts collide with the swept-up, cold, dense shell from the nova explosion and drive Richtmyer–Meshkov instabilities in it. The resulting overturn and fragmentation leave dense knots connected by filamentary structures and surrounded by hot gas. The fragmentation scale observed and the apparent lack of strong thermal conduction around the knots can perhaps be explained by a swept-up interstellar magnetic field limiting fragmentation at the smallest scales. We predict that deeper optical observations of T Pyx’s ejecta using HST or other instruments will show filamentary structures connecting the knots, while X-ray observations may be able to find evidence of the interaction between the hot gas and the dense cold knots, although quantitative predictions of its strength await simulations including magnetized thermal conduction. These conclusions should be generalizable to other classical novae followed by RNe.

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