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A MODEL FOR THE CENTAURUS X-3 PHENOMENON

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ABSTRACT

A model for Cen X-3 is presented according to which the X-ray emission comes from an atmosphere heated by shock waves produced by surface pulsations of a white dwarf. This model can explain the luminosity, period, and spectrum of Cen X-3. The way that these quantities vary with pulsation period and amplitude is discussed.

I. INTRODUCTION

As the data on X-ray stars accumulate, they form a growing body of evidence in support of the contention that these objects are compact and hence probably associated with the final phases of stellar evolution. A case in point is Cen X-3, which has a well-defined period of 4.84 s that is stable over at least half a year (Giacconi *et al.* 1971). Associating this periodicity with either rotation or vibration requires at least as compact an object as a degenerate dwarf.

Here we consider a model for Cen X-3 according to which the X-ray emission comes from an atmosphere heated by shock waves produced by surface pulsations of a degenerate dwarf. Before pulsating X-ray stars were discovered, such a model was proposed for Sco X-1 by Cameron (1966) and independently by Rose (1967). Subsequently Mock (1967) examined the model in some detail by performing numerical hydrodynamic calculations of the shock heating of an atmosphere surrounding a degenerate dwarf. The computed X-ray emission was periodic, and for the parameters considered, the luminosity was always less than 10^{35} ergs s^{-1}, facts which evidently discouraged further consideration of the model for Sco X-1. Here we attempt to revive the pulsating degenerate dwarf as a model for the Cen X-3 phenomenon.

The basic observational data on Cen X-3 are given by Giacconi *et al.* (1971). The average X-ray flux from this object is $\sim 10^{-8}$ ergs cm^{-2} s^{-1}, which corresponds to an average X-ray luminosity of $\sim 10^{36}$ ergs s^{-1} for an assumed distance of 1 kpc. Its spectrum resembles a thermal-bremsstrahlung spectrum with $T \sim 2 \times 10^8$ ° K and low-energy cutoff $E_a \sim 4$ keV. The mean fraction of radiation between 2 and 6 keV which is pulsed is ~70 percent, and the total luminosity has been observed to change by a factor of 10 in 1 hour. Finally, individual jumps in the period of relative magnitude $\sim \pm$ a few times 10^{-4} have been observed, but the net change in period is nearly zero over a timescale of months. Any model for Cen X-3 must provide the answers to the following basic questions: (1) What is the source of energy? (2) How is the energy transformed into X-radiation in such a way as to explain the intensity and the spectrum? (3) How can the time variations be accounted for?

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The model discussed here can provide plausible answers to all these questions. The energy source is nuclear burning in the shell of a degenerate dwarf. A shell of $\sim 10^{-5} M_\odot$ would suffice to power the source for $\sim 10^2$–10^3 years. Such a shell could be a result of stellar evolution, or it could be due to accretion onto the surface from a binary companion, which might imply a much longer lifetime. Whatever its origin, the shell energy source can induce pulsations whose frequency depends on the mass and chemical composition of the star. To get periods of the order of 5 s, one needs a star having a mass ~ 1.1–$1.2\ M_\odot$ and a radius ~ 5000 km for cold degenerate dwarfs (Thorne and Ipser 1968; Cohen, Lapidus, and Cameron 1969) and about 8000 km for hot degenerate dwarfs. Under the right conditions these pulsations will produce an X-ray emitting atmosphere above the surface of the degenerate dwarf. The X-ray emission will be produced by thermal bremsstrahlung in a plasma whose density is so large ($\sim 10^{16}$ cm^{-3}) that the cooling time of the plasma is less than a period, so that the pulse structure stands out. The absorption of X-rays below 4 keV is caused by matter flowing away from the star at a rate $\sim 10^{-8} M_\odot$ year^{-1}. On the other hand, if the pulsations are driven by accretion, then an accretion rate of $10^{-8} M_\odot$ year^{-1} is just sufficient to drive the oscillations at $\sim 10^{37}$ ergs s^{-1} and give a low-energy cutoff at 4 keV. However, the calculations presented here assume the mass flow to be directed outward. In either case, the existence of large-scale mass motion in the atmosphere of the star might imply that the atmosphere does not always relax back to the same configuration after each shock wave passes through it; this could account for variations in intensity from pulse to pulse. The sudden changes in period observed in Cen X-3 are more difficult to understand, although they may be due to the effects of a binary companion or to sudden changes in the boundary conditions of the star. The fact that the period fluctuates but seems always to return to the stable value of 4.84 s does seem to be more consistent with a pulsational than with a rotational origin of the period.

II. ORDERS OF MAGNITUDE

a) Shock-Wave Propagation

In order for a large fraction of the radiation to be pulsed, the cooling time of the plasma must be less than the pulsation period t_p, which is of the same order as the expansion time $t_{\exp}$. Thus, $t_c < t_p \approx t_{\exp}$. For temperatures of interest the cooling is predominantly due to bremsstrahlung, so the cooling time for a plasma of temperature T and density N is given by

$$t_c \approx 0.3 T_8^{1/2}/N_{16} \text{ s} , \tag{1}$$

where $T_8 = T/10^8\,°$ K, and $N_{16} = N/10^{16}$ cm^{-3}. For Cen X-3, $T_8 \approx 2$, $t_p \approx 5$, so that densities greater than about 10^{15} cm^{-3} are required. When $t_c < t_p$, the usual adiabatic shock relations are no longer valid. In the limit where $t_c \ll t_p$ all the thermal energy given the particles by the shock wave is radiated away and the motion of the shock is approximately that of a shell "snowplowing" matter in front of it.

If one assumes a plane geometry,[1] the equation of motion of this shell is

$$d(MU)/dt = -Mg , \tag{2}$$

where M is the mass of the shell, U its velocity, and g is the gravitational acceleration. In the early phases of the motion, gravity is unimportant, and the velocity of the shell $U \sim P_0/M$, where P_0 is the initial momentum. It can be estimated from the impulse given the atmosphere by the star:

$$P_0 \sim \pi R^2 \rho_* U_*^2 t_p , \tag{3}$$

[1] The assumption of plane geometry is not critical. The equation of motion in the spherically symmetric case can be easily solved, and yields essentially the same results.

where ρ_* and U_* are the density and maximum velocity of the surface of the star. The sound waves generated by this motion will travel upward through the atmosphere, increasing in velocity as the density decreases, so that ρU^2 remains approximately constant in the wave. Eventually the velocity of the sound waves becomes supersonic and a shock wave forms. If the density and velocity at this point z_0 are given by ρ_0 and U_0, respectively, then $P_0 \sim \pi R^2 \rho_0 U_0^2 t_p$, and the mass M behind the shock at a height z above the surface is

$$M = \pi R^2 \rho_0 U_0 t_p + 4\pi R^2 \int_{z_0}^{z} dz \rho(z) . \tag{4}$$

The equation of motion (2) then gives the velocity U as a function of height:

$$U^2 = -\frac{2g}{M^2} \int_{z_0}^{z} M^2 dz + \left(\frac{\pi R^2 \rho_0 U_0^2 t_p}{M}\right)^2 , \tag{5}$$

where $M(z)$ is given in equation (4). This solution for the velocity of the snowplowing shock in a plane-parallel atmosphere is valid for any initial density profile. Assuming the density in the atmosphere beyond z_0 to be constant until a height H where the velocity of the shell drops to zero yields for H

$$H \sim (\tfrac{1}{4} U_0 t_p)\{[1 + (6U_0/gt_p)]^{1/3} - 1\} . \tag{6}$$

The corresponding time t_m for the velocity to go to zero is given by an elliptic integral which may be approximated by

$$t_m \sim t_p(U_0/gt_p) \tag{7}$$

for all values of U_0/gt_p. Meanwhile the time for free-fall from the height H back to the surface of the star is

$$t_{ff} = (2H/g)^{1/2} \sim t_p(U_0/2gt_p)^{1/2}\{[1 + (6U_0/gt_p)]^{1/3} - 1\}^{1/2} . \tag{8}$$

In the case where $U_0 \lesssim$ the escape velocity, one can show that the deceleration due to the sweeping up of matter is important compared with gravity only for a short time $\lesssim t_p/10$. Therefore, for most of the motion $dU/dt \approx -g$, $U \approx U_0 - gt$, and $t_m \sim t_{ff} \sim U_0/g$.

Equations (7) and (8) show that the behavior of this class of model is determined by the ratio U_0/gt_p, which is a measure of the ability of gravity to overcome the initial momentum of the shock. If $\frac{1}{2}t_p$ is much greater than t_{ff} (or equivalently $U_0/gt_p \ll \frac{1}{2}$), the atmosphere will collapse in spite of the surface pulsations. On the other hand, if $\frac{1}{2}t_p$ is much less than t_m (or $U_0/gt_p \gg \frac{1}{2}$), the atmosphere will be blown off by successive pulses. Thus, only those degenerate dwarfs for which $U_0 \sim \frac{1}{2}gt_p$ can be pulsating X-ray sources. For Cen X-3, $t_p \sim 5$ s, implying $g \sim 2 \times 10^8$ cm s^{-2}, and the electron temperature $T \sim 2 \times 10^8$ ° K, implying $U_0 \sim 3 \times 10^8$ cm s^{-1}, so $U_0/gt_p \sim 0.3$.

Degenerate dwarfs with masses less than about 1 $M_\odot$ obey the mass-radius relationship $M \propto R^{-3}$ so $g \propto R^{-5}$. Furthermore, $t_p \propto R^3$ (cf. Ostriker 1971), so $gt_p \propto R^{-2}$. In general, U_0 must be proportional to the relative displacement $\delta R/R$; and if U_0 is proportional to U_* which varies as $R^{-2}\delta R/R$, then for a given relative surface displacement, $\delta R/R$, the ratio U_0/gt_p is independent of R. Therefore, for all degenerate dwarfs with $M < M_\odot$, if U_0/gt_p is to be of the order of 0.5, then U_0 must be approximately the escape velocity, U_{esc}. As the mass approaches the Chandrasekhar limit, U_{esc}/gt_p gradually decreases but only to $\sim$0.2 (see Thorne and Ipser 1968), so even the shortest-period degenerate dwarf might have an extended atmosphere.

b) Mass Loss

In reality the density in the atmosphere is not a step function but decreases on a scale $H \sim \rho/(d\rho/dz)$. When the density drops to the point where the radiative cooling time is greater than the expansion time, the snowplow approximation no longer holds. Because of the density gradient, the speed of the shock then increases and much of the atmosphere beyond this point is blown off. This critical point is given by the conditions $t_c \gtrsim \frac{1}{2}t_p$ and $z \sim H$, from which the mass loss per period can be estimated:

$$\Delta M \lesssim 5 \times 10^{-8} R^2 H T_8^{1/2}/t_p \text{ g} . \tag{9}$$

For Cen X-3, $\Delta M \lesssim 5 \times 10^{18}$ g per period.

c) Luminosity

When the shock is in the snowplowing phase, the swept-up matter radiates all its energy in a time t_c which is short compared with the expansion time, so the X-ray luminosity L is given to a good approximation by the rate at which the particles are accelerated to a velocity U by the shock; then

$$L \sim 2\pi R^2 \rho U^3 . \tag{10}$$

From equation (2) we can estimate what the pulse profile should be when $t_c \ll t_p$ and when electron scattering is negligible. Initially there is a rapid linear rise to maximum in a time of the order of the cooling time behind the shock (or the time required for the photons to escape from the source, whichever is longer), followed by a decay which is approximately proportional to $(1 - gt/U_0)^3$ in the plane-parallel, constant-density case. The time-averaged luminosity is obtained from equations (5) and (10):

$$\langle L \rangle \sim \tfrac{1}{2}\pi R^2 \rho_0 U_0^3 (U_0/gt_p) . \tag{11}$$

Since $U_0 \sim \frac{1}{2}gt_p$ if significant X-ray emission is to occur, $\langle L \rangle \sim \pi R^2 \rho_0 (gt_p)^3/32$.

For Cen X-3, $\langle L \rangle \sim 1 \times 10^{36} N_{16}$ ergs s^{-1}. The requirement that the cooling time be less than the rise time fixes a lower limit on N_{16} of 0.8. On the other hand the time for photons to escape from the atmosphere must be less than the rise time t_c:

$$t_{esc} \sim R(1 + \tau_{sc})/c < t_r , \tag{12}$$

where τ_{sc} is the electron-scattering optical depth. This fixes an upper limit on N_{16} of ~ 3. Therefore, $8 \times 10^{35} < \langle L \rangle < 3 \times 10^{36}$ ergs s^{-1}. Comparison with the observed flux of Cen X-3 then yields an estimate of the distance in kiloparsecs of $1 \lesssim d_{kpc} \lesssim 2$.

Note that $\langle L \rangle$ is restricted to a fairly narrow range, implying that the above inequalities are not strongly satisfied. In fact, any significant density gradient in the atmosphere implies a longer cooling time as the shock moves outward; this probably explains the large tail in the pulse shape of Cen X-3. Meanwhile, the observed flares may be due to an increase in N caused by sudden changes in mass loss or by changes in the accretion rate.

d) Spectrum

The low-energy spectrum will be greatly modified by absorption due to the matter flowing away from the star. At high energies, the average spectrum is given approximately by

$$\left\langle \frac{dL}{dE} \right\rangle \sim \frac{t_m}{Ht_p} \int_{z_0}^{z_0+H} R(U) \mathfrak{E}(E) dz \tag{13}$$

where $R(U) \propto U$ is the rate that particles are accelerated to U and $\mathfrak{E}(E)$ is the emitted energy spectrum integrated over the lifetime of a hot plasma. In general, $\mathfrak{E}(E) \propto E_1(E/$

kT), where $E_1(x)$ is an exponential integral of the first kind and $kT \sim \frac{1}{6}M_pU^2$. Then using the solution for U, the average spectrum becomes

$$\left\langle\frac{dL}{dE}\right\rangle \propto \frac{1}{x}[e^{-x} - (1+x)E_2(x)]\,, \tag{14}$$

where $x = 6E/m_pU_0^2$ and $E_2(x)$ is an exponential integral of the second kind. Then, asymptotically, for $x \gg 1$, $\langle dL/dE\rangle \propto e^{-x}/x^2$, steeper than both an exponential and the spectrum of a hot plasma integrated over time. As a function of phase the temperature varies approximately as $U^2 \propto (1 - gt/U_0)^2$ so that the spectrum at the end of the pulse should be softer than near the maximum.

The low-energy X-rays will be absorbed by the plasma flowing away from the star. Initially the plasma is too hot and collisions maintain a fully ionized state, but in about one expansion time it will have cooled sufficiently for recombination to occur for an ion of nuclear charge $Z \sim 10$. Recombination occurs at a density $N_c \sim 3 \times 10^{15}T_8^{1/2}/t_p$. The X-ray optical depth produced by this matter can then be estimated by using conservation of mass along with equation (9) for the mass loss:

$$\tau_x \lesssim 2 \times 10^{15}HT_8^{1/2}\sigma(E)/t_p\,, \tag{15}$$

where $\sigma(E)$ is the X-ray absorption cross-section. For energies greater than about 2 keV, this cross-section is given by $\sigma(E) \sim 3 \times 10^{-22}/E^3$ assuming "cosmic" abundances (Brown and Gould 1970). Therefore, the optical depth is unity at

$$E_a \lesssim (6 \times 10^{-7}HT_8^{1/2}/t_p)^{1/3} \sim 5\text{ keV}\,, \tag{16}$$

which is consistent with the observed $E_a \sim 4$ keV. The absorption cutoff will occur at lower energies if the atmosphere is deficient in medium-heavy elements, or if photoionization by the X-ray source keeps the gas ionized. As long as the luminosity is less than about 10^{36} ergs s^{-1}, this latter effect will not be too important.

Since free-free absorption by the hot envelope is negligible, the optical spectrum is a combination of the radiation from the hot plasma and the blackbody radiation from the surface of the degenerate dwarf. Which of these dominates depends on the parameters chosen. Optical radiation from the plasma will be pulsed only if the plasma behind the shock can cool to $T \lesssim 10^4$ ° K, while radiation from the stellar surface would not be significantly pulsed because of electron scattering and because the stellar surface is hottest when the surface area is smallest. For a surface temperature of 4×10^5 ° K, a radius of 8×10^8 cm and an X-ray luminosity of 10^{36} ergs s^{-1}, optical radiation from the surface is about twice that from the hot plasma, and they add up to produce a blue magnitude of ~16 plus a correction for extinction. This is consistent with the limiting blue magnitude of 16 for a blue pulsating object set by Elliot *et al.* (1971). Meanwhile, their limit of 19 mag for an unusual blue object within the Cen X-3 error box is consistent with this model if ~3 mag of absorption occur along the line of sight in the disk. Furthermore, since mass loss in stars is often associated with the formation of a dust shell (cf. Neugebauer, Becklin, and Hyland 1971), it is possible that much of the optical and ultraviolet radiation from Cen X-3 is absorbed and reradiated in the infrared.

III. NUMERICAL RESULTS

In order to check the validity of the order-of-magnitude estimates, numerical hydrodynamic calculations were carried out to determine the temperature and emission profiles caused by a shock wave propagating through the extended atmosphere of a hot white dwarf. The gas-dynamic equations were solved by using standard methods with the assumption that the atmosphere is optically thin; this simplifies the equation of energy conservation since radiative losses can be treated as a heat sink. A piston at the base of the atmosphere was assumed to move with constant velocity and with amplitude $\delta R/R = 0.1$. The shock wave forms when the piston has negative (downward) velocity, but it is strengthened and the highest temperatures are achieved during the subsequent compression phase. The shock-wave propagation depends upon the density profile toward which the atmosphere relaxes. This fact was ignored in the previous section, but the two approaches yield essentially the same results for the luminosity, mass loss, and average spectrum.

The numerical calculations were carried out for five consecutive pulsation periods, and results from the last two cycles are shown in figure 1. The temperature at the shock rises until after peak emission and then gradually falls off. The temperatures obtained here are about a factor of 2 less than observed for Cen X-3. Higher temperatures can be easily obtained by assuming a somewhat less massive atmosphere and therefore a somewhat lower density. Because of this fact, a cutoff energy of $E_a \sim 2$ keV was assumed in calculating the emission profile. The peak luminosity was found to be about 10^{36} ergs s^{-1}. Furthermore, the outer two mass shells seemed to be leaving the star corresponding to a mass loss $\sim 10^{-8}$ $M_\odot$ year^{-1}, in agreement with equation (9).

These calculations differ from those of Mock (1967) in that these calculations assume a more massive atmosphere and a pulsation amplitude higher by a factor of 10. We therefore obtain a higher luminosity and sharper pulses. The observed pulse width of Cen X-3 is somewhat broader than the result shown in figure 1, but broader pulses could also be produced with a less massive atmosphere since the cooling time would then be longer. A more complete discussion of this problem will be undertaken in a later paper.

IV. SCALING LAWS

Using a mass-radius relationship of the form $M \propto R^{-3}$, a period-radius relationship $t_p \propto R^3$, and the fact that $U_0 \propto g t_p \delta R/R$, we can scale the results of § III. From equation

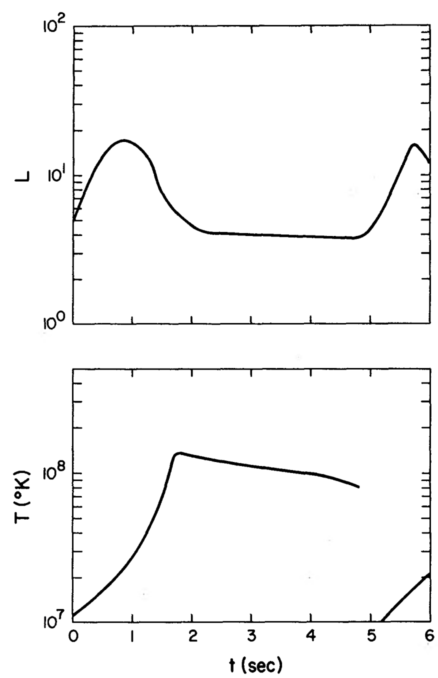

FIG. 1.—Plot of the total X-ray luminosity in arbitrary units and the temperature at the shock versus time for a white dwarf with an extended atmosphere pulsating with a period of 5 s.

(6) the scale height must vary as $H \propto t_p^{1/3}(\delta R/R)^{4/3}$, while from equation (9) the rate of mass loss $\Delta M/t_p \propto t_p^{-5/3}(\delta R/R)^{7/3}$. Therefore, for the directly observable quantities,

$$\langle L \rangle \propto N t_p^{-4/3}(\delta R/R)^4 \,, \tag{17}$$

$$T \propto t_p^{-4/3}(\delta R/R)^2 \,, \tag{18}$$

and

$$E_a \propto t_p^{-4/9}(\delta R/R)^{7/9} \,. \tag{19}$$

Thus, the strong dependence of $\langle L \rangle$ on $\delta R/R$ implies that $\delta R \sim 0.1\, R$ for significant X-ray emission to occur. In fact, normalizing to Cen X-3 (for which $\delta R/R \sim 0.1$), the conditions $t_c < t_r$ and $t_{esc} < t_r$ yield

$$8 \times 10^{35}(5/t_p)^3(\delta R/0.1R)^5 \lesssim \langle L \rangle \lesssim 3 \times 10^{36}(5/t_p)^{5/3}(\delta R/0.1R)^4 \,. \tag{20}$$

The longer-period objects should thus be less luminous, have a softer spectrum, and show less absorption for a fixed $\delta R/R$.

A pulsating white dwarf can give rise to quasi-steady X-ray emission if $t_c > t_p$ or $t_{esc} > t_p$. In the first case we arrive at a limiting luminosity $\sim 8 \times 10^{34}(5/t_p)^3$ ergs s^{-1}, in agreement with the numerical results of Mock (1967). In the case where $t_{esc} > t_p$, the densities must be greater than $\sim 3 \times 10^{17}(t_p/5)^{1/3}$ cm^{-3} and the luminosities greater than $3 \times 10^{37}\ (5/t_p)^{5/3}\ (\delta R/0.1\ R)^4$ ergs s^{-1}. When the luminosity exceeds 10^{38} ergs s^{-1}, radiation pressure would blow the atmosphere off the star, so $3 \times 10^{37}\ (5/t_p)^{2/3} \lesssim L \lesssim 10^{38}$ ergs s^{-1}. For these models the photoionization by the source might maintain a sufficiently high state of excitation in the outflowing plasma that no low-energy X-ray absorption occurs. Such a model with $t_p \sim 10$ s might just fit the Sco X-1 data, although the density seems a bit too high.

V. SUMMARY AND PREDICTIONS

We have shown that a pulsating degenerate-dwarf model provides a plausible explanation for many of the observed features of Cen X-3 such as the period, the luminosity, and the spectrum. This model entails the following predictions: (1) The X-ray spectrum above the cutoff averaged over a cycle should not be that of a uniform temperature hot plasma, i.e., a simple exponential, but should be better fitted by a steeper function as in equation (14). (2) The variation of the X-ray spectrum with phase should be such that the temperature continues to rise with emission until after peak emission, at which time the spectral shape gradually softens during the fall. (3) The optical spectrum is difficult to predict because of the possibility of a binary companion, and the lack of knowledge concerning the formation of an absorbing dust cloud by the outflowing matter. If there is no binary and no dust cloud formed, then the optical spectrum should be that of a blackbody with a temperature $\sim 5 \times 10^5$ ° K, with a B-magnitude $\sim$18–19. If a dust cloud is present, then $\sim 10^{36}$–10^{37} ergs s^{-1} are emitted in the infrared and Cen X-3 should be a conspicuous infrared object. Otherwise, the far-ultraviolet radiation would only partially ionize the hydrogen in a fairly large volume with rather poorly defined boundaries since the ionization cross-section goes as E^{-3} and since most of the emitted ultraviolet radiation has energy much greater than the ionization potential of hydrogen. Thus a rather large and nebulous partially ionized region would exist around this object. (4) If other objects of this type are found, then the luminosity, etc., should vary with period and amplitude approximately as described in § IV. This would be perhaps the strongest evidence bearing on the validity of this model. (5) If this model is to account for nonpulsating X-ray stars, then electron scattering must wash out the pulses. Then the density must be rather high ($\sim 10^{17}$ cm^{-3}), and no time variations on a timescale less than a few seconds should be observable.

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Note added in proof.—Further observations of Cen X-3 from *Uhuru* have shown the existence of 2-day periodic variations in X-ray intensity and correlated sinusoidal variations in the period of the $4^{s}.8$ pulsations (Schreier *et al.*, *Ap. J.* [*Letters*], **172** [in press]). They point out that the changes in intensity are consistent with occultation by a binary companion, with the changes in period being due to the Doppler effect. The existence of a binary companion suggests that accretion may be the source of fuel for the nuclear burning which drives the pulsations.