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LATE EVOLUTION OF VERY LOW MASS X-RAY BINARIES SUSTAINED BY RADIATION FROM THEIR PRIMARIES

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ABSTRACT

Under suitable conditions, it is possible for part of the radiation originating in the vicinity of weakly magnetized neutron stars in low-mass X-ray binaries (LMXBs) to illuminate the low-mass secondary and power a wind from it. This evaporative wind can produce a Roche lobe overflow of the secondary with a mass loss close to the neutron star's Eddington limit. This can occur even when gravitational radiation or other causes of loss of angular momentum from the LMXB are weak. A bootstrapping mechanism can keep the LMXB X-ray luminosity near the Eddington limit until the secondary's mass falls below a few hundredths of a solar mass. Thereafter it would be expected that the LMXB X-ray luminosity would either drop abruptly by several orders of magnitude or cease completely, if the accretion spun-up neutron star period is less than several milliseconds. As the LMXB secondary mass traverses the range 10^{-1} to a few times 10^{-2} $M_{\odot}$, the primary radiation can substantially shorten the lifetime of an LMXB compared with that predicted by standard gravitational radiation scenarios. However, it will not greatly affect those earlier stages of evolution when the companion is heavier. Because the LMXB evolution speeds up dramatically through this very low mass secondary range, few LMXBs in this phase would be expected to be observed. Several features of LMXB observations can be plausibly explained by this model.

Subject headings: gamma rays: general — pulsars — stars: accretion — stars: evolution — stars: neutron — X-rays: binaries

I. INTRODUCTION

The evolution of many X-ray binaries containing neutron stars has been thought to be driven by Roche lobe overflow of the companion caused by one of two possible mechanisms depending on the mass of the companion star. (1) If the companion is sufficiently heavy ($m > 0.8\ M_{\odot}$), it may become a giant star in less than a Hubble time and then overflow its Roche lobe because of nuclear core evolution. Two major subclasses of systems belong to this class. For heavier companions this Roche lobe overflow phase may be quite short (10^5 yr or so) and the neutron stars relatively young with large magnetic fields. For lighter companions ($m \simeq 0.8\ M_{\odot}$, say) the overflow may last 5×10^7 yr and the neutron star magnetic field may be substantially weaker. (2) If the companion is too light to evolve off the main sequence in a Hubble time and the binary is tight enough for gravitational radiation to bring the stars close to each other, then the secondary's Roche lobe contracts to its stellar surface. Additional losses of angular momentum may also drive the evolution of these low-mass X-ray binaries (e.g., winds which contribute to magnetic braking). This second class of binaries will usually evolve into systems containing neutron stars with light ($m < 0.1\ M_{\odot}$) white dwarf companions. It is the evolution of this family of very low mass X-ray binaries (LMXBs) that we shall consider below.[4]

There are a number of unanswered questions raised by models for these LMXBs:

1. Why is the neutron star accretion rate relatively near the Eddington limit, $\dot{M} \simeq 10^{18}$ g s^{-1} $\simeq 1.6 \times 10^{-8}\ M_{\odot}$ yr^{-1}, for a large number of LMXBs? The observed sources have luminosities near the accretion-supported X-ray luminosity of $L_X \simeq 10^{38}$ ergs s^{-1}, and their number decreases rapidly down to 10^{36} ergs s^{-1} (Long and van Speybroek 1983). Sources in the range $L_X \sim 10^{35}$–10^{36} ergs s^{-1} may not have been observed at all, considering distance and anisotropy uncertainties. Sources with $L_X < 10^{35}$ ergs s^{-1} appear to be cataclysmic variables in which a white dwarf replaces the neutron star as the compact accreting object. If gravitational radiation does indeed drive the evolution of a good fraction of LMXBs (at least those with very short orbital periods), one should see more rather than fewer sources with accretion rates and X-ray luminosities far below the Eddington limit, since mass transfer driven by gravitational radiation becomes increasingly small as the mass-losing secondary becomes lighter. Such an effect is, possibly, observed in cataclysmic variables, which are also thought to evolve because of gravitational radiation; but it is not

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[4] Typical sources belonging to this class of LMXBs are 4U 1820−30 (Stella, White, and Priedhorsky 1987) with a He companion of mass $m \simeq 0.05\ M_{\odot}$, and XB 1916−053 (White and Swank 1982; Smale *et al.* 1988) with a companion of mass $m \simeq 0.1\ M_{\odot}$.

observed in LMXBs. For degenerate dwarf companions, the mass transfer rate driven by gravitational radiation is (Faulkner 1971; Li *et al.* 1980; Rappaport *et al.* 1987)

$$\dot{m} \simeq 10^{-4}\left(\frac{y}{1+y}\right)^{14/3} M_{\odot}\ \mathrm{yr}^{-1}\ , \tag{1.1}$$

where $y = m/M$ is the ratio of the mass of the degenerate light secondary, m, to that of the (much heavier) neutron star, M. Thus $|\dot{m}|$ begins to fall below the Eddington limit as m drops below about 0.1 $M_{\odot}$. In earlier evolutionary stages, when the companion is more massive and still a main-sequence star, it is also sufficiently close to the compact primary when Roche lobe overflow occurs that gravitational radiation easily powers mass accretion rates near the neutron star's Eddington limit. The underabundance of low-luminosity LMXBs is a problem only when one is confronted with theoretical scenarios of late binary evolution driven by gravitational radiation.

2. Wide consensus exists presently that binary millisecond pulsars have been spun up by accretion in LMXBs (Alpar *et al.* 1982). It is also widely accepted that the two isolated millisecond pulsars 1937+214 and 1821−24 were similarly spun up because of many similarities to their binary siblings (see, e.g., Taylor 1986; Taylor and Stinebring 1986). However, standard scenarios for a typical LMXB that might produce a millisecond pulsar severely limit the number of candidates compared with those needed to account for the other observed pulsars. The mass accretion rates in the gravitationally driven LMXBs are predicted to drop, in less than the Hubble time, to such low values that in almost all cases the proposed ancestral neutron star would be spun down to rotational periods of order 1 s (Jeffrey 1986). How do these neutron stars avoid the expected accretion spin-down phase, i.e., why, as discussed under question 1, does $|\dot{m}|$ remain near the Eddington limit for a long interval ($> 10^8$ yr) and then, relatively quickly, turn off almost completely (one way of explaining the statistical discrepancy)?

3. Why does the secondary star often disappear in binary systems containing weakly magnetized neutron stars which have been maximally spun up by accretion? The spin-up scenarios are quite successful in explaining quantitatively the observed relationship between the neutron star's dipole moment and its spin period (see, e.g., van den Heuvel 1986). But the two very rapidly spinning neutron stars, PSR 1937+21 and PSR 1821−24, although their periods and dipole moments are similarly related, do not have companions. How has the companion disappeared? Some possible scenarios for the disappearance of the companion have been suggested by Bonsema and van den Heuvel (1984), which include the coalescence of a red dwarf with a massive white dwarf or with a neutron star. Nevertheless, these models are not free of problems, and there is no evidence for any of their intermediate stages.

4. What is the origin of the coronae which seem to exist around neutron stars in several LMXBs (White and Holt 1982; White, Stella, and Parmar 1988; Stollman *et al.* 1987; Mason 1986)?

In this paper it is proposed that the answer to all these questions may lie in the nature and consequences of physical processes caused by energetic radiation produced in the inner region of the accretion disk close to the neutron star and intercepted by the secondary. For LMXB neutron stars the relevant components are, predominantly, the soft X-ray flux, which is similar (albeit relatively more intense) to the radiation from massive X-ray binaries, and a possible intense flux of MeV γ-radiation and/or a comparable $e^{\pm}$ wind. This radiation proves to be effective in sustaining winds because it can deposit most of its energy after passing through only a relatively small depth (up to a few grams per square centimeter) of stellar atmosphere. Thus it deposits much of its energy in a region where radiative cooling is less effective than outward bulk motion of the heated gas in balancing the incoming energy deposition. Possible effects on the fate of the secondary in the postaccretion phase of LMXB evolution by various kinds of radiation from a companion pulsar are considered in another paper (Ruderman, Shaham, and Tavani 1989, hereafter RST). We proposed, along the lines suggested in previous discussions of the Her X-1 system (Arons 1973; Basko and Sunyaev 1973), that the radiation intercepted by the LMXB secondary drives a strong stellar wind from the secondary. This wind may be able to remove angular momentum from the system and, at the late evolutionary stages with a very light secondary, can dominate gravitational radiation in controlling the LMXB's evolution. Mass and angular momentum loss driven by the evaporative wind leads to the secondary's Roche lobe overflow during late stages of binary evolution if a substantial fraction of wind material escapes from the system. Otherwise, if the gas outflowing the companion is not able to escape the system, the wind itself feeds the accretion disk as the secondary contracts into its expanding Roche lobe (Kluźniak *et al.* 1988*a*). Accretion on the surface of the neutron star, in turn, is the energy source of the radiation which powers the wind, and, under some conditions, a bootstrapping mechanism can be established. When the secondary's diminishing mass becomes sufficiently small (e.g., $m \sim 2 \times 10^{-2}\ M_{\odot}$) the Roche lobe filling secondary has moved so far away from the neutron star that it no longer intercepts enough of the accretion-powered radiation flux to maintain this bootstrap symbiosis. The accretion rate then rapidly falls by several orders of magnitude to the value which would have existed if gravitational radiation alone had been the cause of angular momentum loss from the LMXB. If, at that time, the neutron star's spin is sufficiently rapid that its light cylinder lies within its new Alfvén radius, accretion stops completely and the consequent spin-down energy loss is that of an isolated pulsar. Because the radiation induced mass loss of the secondary remains of order $10^{-8}\ M_{\odot}\ \mathrm{yr}^{-1}$ as the secondary mass falls from $10^{-1}\ M_{\odot}$ to a few times $10^{-2}\ M_{\odot}$, the binary evolution time scale in this phase ($\lesssim 10^7$ yr) is much less than typical LMXB lifetimes. Thus we would expect to observe only a small fraction of LMXBs in the late stages of evolution discussed above.

In § II we consider the accretion-powered radiation from the best studied prototype of "massive" X-ray binaries, the X-ray pulsar system Her X-1 (McCray *et al.* 1982) for which the possibility of self-excited winds had already been considered in 1973 (Arons 1973; Basko and Sunyaev 1973). In the case of Her X-1 the X-rays overwhelm any other radiation. The much less intense, deeply penetrating, 10^{12} eV γ-rays intercepted by the secondary do nothing beyond a negligible warming of the secondary's surface, while the soft X-rays can actually lead to an evaporative wind from the secondary's surface. However, if the dipole moment of the neutron star is strongly diminished and the secondary ($m \simeq 2.2\ M_{\odot}$) is replaced by a light dwarf ($m < 10^{-1}\ M_{\odot}$), the nature and consequences of the intercepted radiation are different. The changes in the soft X-ray and γ-ray flux and in the accompanying $e^{\pm}$ wind are discussed in § VI. These are associated with the moving inward of the inner edge of the accretion disk which corresponds to the boundary with the neutron star's corotating magnetosphere (Alfvén radius).

The 10^{12} eV γ-ray emission of Her X-1 seems to be powered by particle acceleration processes accuring near the Alfvén radius R_A, where differentially rotating parts of the Keplerian accretion disk and the neutron star are connected by a strong magnetic field. The produced luminosity is comparable to the power available from accretion in the region close to the Alfvén radius either when this dynamo sustains a steady particle acceleration or when it produces currents which drive magnetic field line reconnection in that region. For an LMXB with a surface magnetic dipole field $B_s \simeq 10^9$ G and Eddington limit accretion rate, the Alfvén radius R_A is almost 2 orders of magnitude smaller than in the case of Her X-1, where $B_s \simeq 4 \times 10^{12}$ G and $R_A \simeq 2 \times 10^8$ cm (for $\dot{M} \simeq 10^{17}$ g s^{-1}). One consequence is that almost 100 times more power is available for the accelerator (which is the ultimate γ-ray and $e^{\pm}$ wind source) in LMXBs than in Her X-1.

Furthermore, the magnetic field at the LMXB Alfvén radius is sufficiently higher than that at the Alfvén radii of massive X-ray pulsars, and therefore 10^{12} eV γ-rays would materialize there into $e^{\pm}$ pairs before they can escape. Lower energy multi-GeV γ-rays convert to $e^{\pm}$ pairs in collisions with the accretion-powered X-rays whose local flux is much greater at the smaller R_A. A further process that can degrade very high energy γ-ray photons are collisions with both electrons (and possibly positrons) and with X-rays within the X-ray corona thought to exist around bright LMXBs (and, indeed, possibly even created and energized in part by these very pairs and photons). The energy released in acceleration processes near the Alfvén radius ultimately escapes from that region as γ-rays and X-rays with possible energies ranging from around the threshold for pair creation, $\gamma + \gamma \rightarrow e^+ + e^-$, down to typical coronal energies near 5 keV. The expected spectrum of this emerging radiation is discussed elsewhere (Kluźniak *et al.* 1988*b*, *c*). A comparable flux of $e^{\pm}$ pairs may also escape. For large accretion rates, however, the X-ray flux will cool the $e^{\pm}$ pairs down to the X-ray temperatures. The optical depth of these pairs for scattering by X-rays emitted close to the neutron star surface may be large enough to prevent observation of strongly modulated X-ray beams from accretion-heated regions on the LMXB's spinning neutron star surface. In § VII we will discuss how strong γ-ray beams, especially in the regime $L_\gamma > 10^{37}$ ergs s^{-1}, may bore and sustain holes through LMXB coronae. The soft X-rays could also use these holes to reach the companion. Even if channel boring is ignored, a subsidence of coronae for sub-Eddington accretion rates could also allow primary radiation to reach the secondary. In the light-companion regime ($m < 0.1\ M_\odot$), where winds excited by illumination of the secondary may control LMXB Roche lobe overflow, a falling intensity reaching the secondary would automatically act to reduce an absorbing corona until a compatible symbiosis is reached. Soft X-rays might also reach the companion after multiple scattering in a very large hot corona (possibly associated with the secondary's induced wind).

In § III we consider the wind expected from a very light secondary intercepting the radiation emitted near an LMXB neutron star. X-rays powered by accretion, if intercepted by the secondary, may be less effective in sustaining winds from the secondary of LMXBs than in the case of Her X-1, even if the X-ray spectrum of LMXBs relative to that of Her X-1 may be substantially softer. For typical values of the relevant quantities govening the evolution of LMXBs with degenerate companions the temperature of the outer coronal atmosphere heated by X-rays may turn out to be too small for the formation of an evaporative wind. The presence of heavy elements in the companion's atmosphere with abundance larger than average may be relevant in absorbing the X-rays more efficiently and for producing a relatively hotter corona. For cosmic abundance companions soft γ-rays may be needed to drive a wind out of a basically static corona heated by X-rays. They could also power a wind on their own.

In § IV, which describes the main result of this paper, we consider LMXB evolution which is self-sustained by the secondary winds of § III. The neutron star accretion rate is determined, essentially, by a bootstrapping condition that keeps $\dot{M}$ close to 10^{17}–10^{18} g s^{-1} until the secondary's mass falls below a few hundredths of a solar mass. Accretion then suddenly falls by several orders of magnitude or ceases completely, depending upon the neutron star spin rate when that secondary critical mass is reached. In the latter case, which occurs for millisecond pulsar spin rates, the companion might be evaporated by the incident pulsar X-ray, γ-ray, and $e^{\pm}$ winds as discussed elsewhere (RST).

II. HER X-1 RADIATION AND INDUCED SECONDARY WINDS

Her X-1 is probably the best studied of all accreting neutron stars in binaries, and it can be considered a prototype for such systems. In this section we consider two phenomena occurring in Her X-1 and relevant for LMXBs: the self-excited wind and the emission of TeV γ-rays.

a) The Self-excited Wind

In the Her X-1 binary system the secondary star, HZ Her, is relatively massive (2.2 $M_\odot$) and probably overflows its Roche lobe because of continuing core nuclear evolution. The maximum resulting mass-loss rate of the secondary, which also takes into account effects of magnetic braking, has been estimated to be (Rappaport, Verbunt, and Joss 1983)

$$\dot{m} \simeq -2 \times 10^{-11} \left(\frac{P_{\rm orb}}{1\ {\rm hr}}\right)^{3.2} M_\odot\ {\rm yr}^{-1}\ , \tag{2.1}$$

where $P_{\rm orb}$ is the binary orbital period. For Her X-1, $P_{\rm orb} = 41$ hr and $|\dot{m}|$ is large enough to sustain a mass accretion rate which is a considerable fraction of the neutron star's Eddington limit. One result of this accretion is strong X-ray emission (McCray *et al.* 1982) with an integrated observed energy flux detected at the Earth of $F_X \simeq 6.7 \times 10^{-10}$ ergs $cm^{-2}\ s^{-1}$ for X-ray photons with energy E_X in the range 0.1 keV $\leq E_X \leq$ 0.8 keV and $F_X \simeq 40 \times 10^{-10}$ ergs $cm^{-2}\ s^{-1}$ for 2 keV $\lesssim E_X \lesssim$ 20 keV corresponding to a total (0.1–20 keV) X-ray luminosity larger than 10^{37} ergs s^{-1}. Interception of an X-ray flux *with an assumed flat spectrum* up to 30 keV by the HZ Her secondary has been estimated to be able to sustain a stellar wind of magnitude (London, McCray, and Auer 1981; London and Flannery 1982)

$$\dot{m}_w \simeq -8 \times 10^{-12} \varphi^{-1/2} \hat{L}_X\ {\rm g\ s}^{-1}\ , \tag{2.2}$$

with φ the secondary's surface gravitational potential per unit mass,

$$\varphi = Gm/R , \tag{2.3}$$

and $\hat{L}_X$ the soft X-ray flux intercepted by the secondary of radius R. The physical mechanism yielding equation (2.2) can be summarized as follows: as long as radiative cooling at relatively large depth in the stellar atmosphere is able to balance heating from external irradiation, only a minor readjustment of the stellar atmosphere occurs. However, as the pressure and density decrease with increasing radius, a density is reached at which cooling is only marginally able to balance energy deposition. The value of the pressure corresponding to this particular density, which depends linearly on the heating rate per atom and inversely on the corresponding cooling rate, determines the outward momentum flux of the wind. From this follows (§ III) the linear dependence of $\dot{m}_w$ on $\hat{L}_X$ and the inverse dependence on the escape velocity.

The main mechanism for energy deposition in equation (2.2) is photoelectric absorption of soft X-rays by the bound electrons left on heavy elements (C, N, O) present in small abundances. Much harder X-rays are mainly scattered; they deposit in the outer atmosphere of the companion only a small fraction of their energy in each scattering, and their energy deposition per atom is too small to be effective in wind formation.

With a distance a from the secondary to the neutron star and an isotropic luminosity L_X,

$$\hat{L}_X = \frac{R^2}{4a^2} L_X \tag{2.4}$$

if disk absorption is neglected. Then for Her X-1 with $a = 4 \times 10^{11}$ cm, equation (2.2) gives

$$\dot{m}_w \simeq -2 \times 10^{-21}\left(\frac{R}{10^{11}\ \mathrm{cm}}\right)^{5/2} L_X \ \mathrm{g\ s^{-1}} . \tag{2.5}$$

With $L_X \simeq 10^{37}$ ergs s^{-1} for Her X-1 and an assumed secondary radius $R = 10^{11}$ cm, the mass transfer rate $|\dot{m}| \sim 10^{17}$ g s^{-1} needed to explain the observed luminosity might possibly be partly a consequence of this wind (London, McCray, and Auer 1981). The angular momentum loss caused the wind could sustain a Roche lobe overflow of the secondary even in the absence of stellar expansion driven by nuclear evolution.

b) *Emission of* TeV *γ-Rays*

Several groups claimed to have detected 10^{12} eV γ-rays from the binary X-ray pulsars Her X-1, Vela X-1, and 4U 0115+63 (Lamb 1986; Weekes 1988). The data from Her X-1, observed as a sporadic source a dozen times by several groups, are the most compelling. These very high energy (VHE) γ-rays, unlike the soft γ-rays considered in §§ III and V, are not at all effective in sustaining mass-loss winds from the secondaries, for two reasons. First, the γ-ray intensity ($L_\gamma \simeq 10^{-2} L_X$) is too small to be important compared with the effects of L_X. Second, most of the energy of these γ-rays is deposited far below the secondary's surface, typically after passage through almost 10^3 g cm^{-2}, where it is quickly shared by so many particles, and especially photons, that its net result is only an unimportant warming of the outer part of the stellar atmosphere. However, as discussed in § III, the γ-ray flux of a strongly accreting neutron star should be very different when the neutron star's dipole magnetic field is 4 orders of magnitude less than the 4×10^{12} G of Her X-1. One expects at least some neutron stars in LMXBs to have such a low field ($B_s \simeq 10^9$ G) if, as discussed above, some have evolved into observed millisecond pulsars.

There is not yet a consensus on the details of the accelerator which powers the charged particles which produce the VHE γ-rays observed in some X-ray pulsars. Because of the large neutron star dipole moments in these pulsars, acceleration by a stand-off shock, supported by accretion onto the neutron star surface, does not seem to be a viable mechanism, since this works only when the magnetic field behind the shock is much less than 10^{10} G (Eichler and Vestrand 1984; Gaisser 1987); nor would the 10^{12} eV γ-rays be able to escape from the region of a strong local magnetic field. Several possibilities have been proposed which utilize the potential and power from the difference in angular rotational speed between the corotating magnetosphere of a neutron star and the accreting matter of a differentially rotating Keplerian accretion disk. In these models, the maximum power is the accretion power available at the neutron star's magnetospheric radius R_A. If this power is efficiently converted into a γ-ray flux, then

$$L_\gamma \sim \frac{R_{ns}}{R_A} L_X , \tag{2.6}$$

with R_{ns} the neutron star radius. For Her X-1, $R_A \simeq 2 \times 10^8$ cm, and equation (2.6) gives $L_\gamma \simeq 10^{-2} L_X$, about that observed. The local Her X-1 neutron star magnetic field at $r \sim R_A$, $B(R_A) = B_s(R_{ns}/R_A)^3 \simeq 5 \times 10^5$ G, will not convert most 10^{12} eV γ-rays into $e^\pm$ pairs even if these γ-rays move normally across the local field lines. The R_A region is about the closest to the neutron star in which the magnetic field cannot prevent the escaping of 10^{12} eV photons. Thus, even if these γ-rays are locally produced, almost all will escape.[5] Substantial potential drops may be produced in the region near R_A because of the possibility large difference between disk and magnetospheric rotational speeds. The potential drop is

$$V \simeq R_A^2 \Delta\Omega\, B/c \lesssim 10^{15} \dot{m}_{18}^{5/7} \tilde{\mu}_{30}^{-3/7}\ \mathrm{V} , \tag{2.7}$$

where $\Delta\Omega$ is the beat frequency between the neutron star magnetosphere and the inner region of the accretion disk around it (Alpar

[5] The location of the accelerator is not constrained in this way in other models which assume that the limb of the secondary is the source of VHE radiation produced from the decay $\pi^0 \to \gamma + \gamma$, with π^0's produced by collisions of protons accelerated near the R_A region. See, e.g., Kazanas and Ellison (1986).

and Shaham 1985), and $\dot{m}_{18}$ and $\tilde{\mu}_{30}$ are the accretion rate (in units of 10^{18} g s^{-1}) and the neutron star's magnetic dipole (in units of 10^{30} G cm^3), respectively.

III. WINDS IN LOW-MASS X-RAY BINARIES

Winds can form in outer regions of an externally irradiated stellar atmosphere when the net heating produced by incoming radiation is able to raise the gas temperature to a value sufficiently high that the gas kinetic energy per unit mass is larger than the local gravitational potential per unit mass. This may happen in a relatively dilute region of the stellar atmosphere where heating caused by Compton scattering and photoionization of heavy elements is substantially larger than radiative cooling. The condition necessary for wind formation, i.e., the occurrence of substantial net heating at a depth characterized by relatively small radiative cooling, is satisfied only for some particular kinds of incident radiation. In this section we describe quite schematically the mechanism of wind formation under X-ray and γ-ray illumination, and discuss the general conditions concerning the incident spectrum of radiation which make possible the formation of a substantial evaporative wind from the surface of a degenerate companion in an LMXB.

Under the effect of an external illumination the entropy per unit mass S of the irradiated matter increases at a rate

$$\frac{dS}{dt} = \frac{\Lambda}{T}, \tag{3.1}$$

where Λ is the rate of net heating (external heating minus radiative cooling) per unit mass (in ergs s^{-1} g^{-1}), and T is the gas temperature. The wind dynamics is determined, under steady state conditions, by the hydrodynamics equations governing the conservation of mass, energy, and momentum of the outflow. The equation giving the conservation of momentum is

$$\frac{d}{dr}\left(\frac{1}{2}v^2 - \frac{Gm}{r}\right) = -\frac{1}{\rho}\frac{dP}{dr}, \tag{3.2}$$

where r is the distance from the stellar center, v is the radial outflow velocity ($dt \equiv dr/v$), m is the stellar mass, and ρ and P are the density and the pressure, respectively. As an atmospheric mass moves outward, its pressure decreases; if cooling balances heating, $\Lambda = 0$, the entropy remains constant, and it can be shown that $P \propto T^{5/2}$. This decrease of pressure, however, would almost always result in a decrease of cooling as long as $n < 5/2$ where the rate of cooling is proportional to the effective cooling (bremsstrahlung and recombination) coefficient multiplied by the local density, i.e., to $PT^{-n} \propto P^{1-2n/5}$. In general, for a given radiation temperature T_R established at the stellar photosphere, Λ will become significantly larger than zero only when the local density is smaller than some critical value ρ_c. The pressure P_0 corresponding to this particular density ρ_c for which external heating dominates cooling is of fundamental importance for the estimation of the gas outflow. The atmosphere can be considered as static at relatively large depths with densities corresponding to pressures much larger than P_0 for which $\Lambda \simeq 0$. P_0 is proportional to $P_0 = \alpha^{-1}T_R^n\Lambda$, where α is the constant in the cooling coefficient per hydrogen atom for a cosmic abundance medium with H completely ionized (at high temperatures the cooling coefficient is $\sim\alpha/T^{1/2}$ cm^3 s^{-1}). For sufficiently high temperatures ($T_R \geq 10^4$ K) the index n is equal to $\frac{1}{2}$ and the relevant pressure corresponding to a region where a substantial temperature gradient is established is $P_{min} = P_0$ (at $T = 3T_R$) (London, McCray, and Auer 1981), i.e.,

$$P_{min} = 3^{3/2}\alpha^{-1}T_R^{1/2}\Lambda\,. \tag{3.3}$$

Because of the nonzero net heating occurring in the illuminated outer atmosphere, local thermal balance requires additional cooling that can be accomplished only by gas expansion, i.e., by the formation of an outflowing stellar wind. When most of the external radiation is deposited at a depth associated with a value of the pressure larger than P_{min}, that radiation contributes only to an overall heating of the star determining T_R, and has low efficiency for energy conversion of radiation into wind kinetic energy. Relatively high efficiency can be obtained only when substantial energy deposition occurs at densities corresponding to pressures equal to or smaller than P_{min}.

Let $-\dot{m}_w$ be the mass outflow rate and v_s the value of the outflowing gas velocity at the sonic point which satisfies the relation $v_s^2 = (5/3)(1/\mu)kT_s$, with T_s the local temperature and μ the mean molecular weight per particle ($\mu = \frac{1}{2}m_p$ for a completely ionized hydrogen gas; $\mu = 0.58$ for a solar composition gas with H and He completely ionized). Equations (3.1) and (3.2) can be integrated, and one can determine the wind velocity at the sonic point. From the equations of hydrodynamics it can be shown that v_s satisfies, for spherical geometry,

$$v_s^2 = \frac{Gm}{2R_s} + \frac{R_s\Lambda_s}{3v_s}, \tag{3.4}$$

where R_s is the sonic radius and Λ_s is the net heating rate per unit mass at the sonic radius. Incident γ-rays, whose heating rate per unit mass $\Lambda_{H,\gamma}$ is independent of the temperature, provide a net heating of the outer atmosphere to a high temperature, and in this case equation (3.4) may be satisfied with $R \simeq R_s$. We will consider below the role played by γ-rays in this context.

On the other hand, X-rays yield a temperature-dependent heating rate per unit mass $\Lambda_{H,X} = \Lambda_X(1 - T/T_X)$ (with Λ_X the X-ray heating rate per unit mass at low temperature) which tends to zero as the temperature approaches a critical value T_X at which photoabsorbtion of X-rays is suppressed because of complete ionization of heavy elements. For X-ray illuminations expected in LMXBs we will show below that equation (3.4) may be satisfied almost always with $R_s > R$, a condition which tends to suppress exponentially the evaporative wind.

A relevant dimensionless parameter may be defined as

$$\eta = \frac{8}{3}\left(\frac{R_s}{R}\right)^{5/2} \frac{R\Lambda_s}{v_e^3}, \tag{3.5}$$

where v_e is the escape velocity from the star, $v_e = (2Gm/R)^{1/2}$. For

$$x = \left(\frac{R_s}{R}\right)^{1/2} \frac{2v_s}{v_e} \tag{3.6}$$

equation (3.4) can be rewritten as

$$x^2 = 1 + \frac{\eta}{x}, \tag{3.7}$$

so that

$$x = \begin{cases} 1 & \text{if } \eta \ll 1, \\ \eta^{1/3} & \text{if } \eta \gg 1. \end{cases} \tag{3.8}$$

For soft X-ray heating (in contrast to, say, soft γ-ray heating) η depends on x. In this case we may write the net heating in dimensionless form (London, McCray, and Auer 1981; London and Flannery 1982) as

$$\eta = \eta_0(1 - \xi x^2) - \eta'. \tag{3.9}$$

Equation (3.9) is meaningful as long as $\xi x^2 \leq 1$, where η_0 is defined in equation (3.13), η' stands for the radiative cooling, and ξ is related to the temperature T_X at which the X-ray absorption drops because of almost complete ionization of the absorbing metals,

$$\xi = \frac{3}{20}\left(\frac{R}{R_s}\right)\frac{\mu v_e^2}{kT_X} \equiv \left(\frac{R}{R_s}\right)\frac{v_e^2}{4v_X^2}. \tag{3.10}$$

Here v_X is defined by equation (3.10) (k is the Boltzmann constant). The temperature T_X depends on the particular absorbing heavy element as well as on its abundance and on the spectrum of the intercepted X-ray radiation. Then equation (3.7) becomes

$$x(x + \xi\eta_0) = 1 + \frac{\eta_0 - \eta'}{x}. \tag{3.11}$$

For pressures below $P_{\min}$ we can neglect η' and find

$$x = \begin{cases} \xi^{-1/2} - \dfrac{1-\xi}{2\xi^2\eta_0}, & \eta_0 \gg 1, \quad \eta_0^{-2/3} \ll \xi \leq 1, \\ \eta_0^{1/3}, & \eta_0 \gg 1, \quad \xi \ll \eta_0^{-2/3}, \\ 1 + \dfrac{1-\xi}{2}\eta_0, & \eta_0 \ll 1, \quad \xi \leq 1, \end{cases} \tag{3.12}$$

with

$$\eta_0 = \frac{8}{3}\left(\frac{R_s}{R}\right)^{5/2} \frac{R\Lambda_0}{v_e^3} \tag{3.13}$$

and Λ_0 the total net heating at zero temperature. Equation (3.12) may be written as

$$v_s = \begin{cases} v_X - \dfrac{1-\xi}{4\xi^2\eta_0}\left(\dfrac{R}{R_s}\right)^{1/2} v_e, & \eta_0 \gg 1, \quad \eta_0^{-2/3} \ll \xi \leq 1, \\ \left(\dfrac{R_s\Lambda_0}{3}\right)^{1/3}, & \eta_0 \gg 1, \quad \xi \ll \eta_0^{-2/3}, \\ \dfrac{1}{2} v_e\left(\dfrac{R}{R_s}\right)^{1/2}, & \eta_0 \ll 1, \quad \xi \leq 1. \end{cases} \tag{3.14}$$

Winds form only if the kinetic energy per particle of the heated gas is sufficiently large to overcome the local gravitational potential; for X-ray illumination the latter condition becomes, for spherical geometry,

$$kT_X \geq \frac{3}{20}\left(\frac{R}{R_s}\right)\mu v_e^2. \tag{3.15}$$

For relatively low values of T_X, equation (3.15) means that the sonic point is far from the star, i.e., $R/R_s \ll 1$ and $v_s \sim v_X < v_e$. For low heating rates, $\eta_0 < 1$ and the wind velocity at the sonic point is about half the escape velocity if $R_s \simeq R$. For strong heating and for a v_X sufficiently large, the wind reaches the sonic point with velocity v_X and generally escapes to infinity with higher energy. If $2v_X$ is close to v_e, the sonic point velocity remains essentially unchanged for all illuminations. (However, one should always remember that R_s may vary with illumination.)

The total mass outflow rate $\dot{m}_w$ is given by

$$-\dot{m}_w = 2\pi\rho_s v_s R_s^2 \equiv \frac{10\pi}{3}\frac{P_s R_s^2}{v_s}\,, \tag{3.16}$$

where ρ_s and P_s are the sonic point values of the density and of the pressure, respectively. If the sonic point is very close to the stellar radius, $R_s - R \ll R$, as in the case for high heating rates, we may substitute in equation (3.16)

$$\lambda \equiv \rho_s v_s = \frac{1}{2v_s^2}\int_R^{R_s} \rho\Lambda\, dr\,. \tag{3.17}$$

With the definition

$$l_s = \frac{\int_R^{R_s} \rho\Lambda\, dr}{\int_R^{R_s} \rho\, dr} \tag{3.18}$$

for the average heating l_s in the acceleration zone, we have

$$\lambda = \frac{l_s \hat{\sigma}}{2v_s^2}\,, \tag{3.19}$$

where $\hat{\sigma}$ is the column density at which the acceleration of the outflowing gas occurs. The substitution in equation (3.16) gives

$$-\dot{m}_w v_s^2 = \pi R^2 l_s \hat{\sigma}\,. \tag{3.20}$$

The pressure difference $\Delta P = P(R) - P(R_s)$ is, in this case,

$$\Delta P = \frac{1}{2}\hat{\sigma}\left(\frac{l_s}{v_s} + \frac{v_e^2}{R}\right) \leq P_{\min}\,. \tag{3.21}$$

Hence equation (3.16) becomes, for $\eta_0 \gg 1$,

$$-\dot{m}_w v_s = 2\pi R^2 \Delta P\,. \tag{3.22}$$

A numerical integration of equations (3.1) and (3.2) will give ΔP and hence $\dot{m}_w$. However, some results can be stated analytically (Basko *et al.* 1977). The most important one is that $|\dot{m}_w|$ decreases rapidly as R_s/R increases, as $\exp[-b(R_s/R)]$ with constant b. Therefore, for the maximum wind mass flux, one would require, from equation (3.15) with $R_s \simeq R$,

$$kT_X \gtrsim 0.087 m_p v_e^2 \tag{3.23}$$

or

$$T_X \gtrsim (10^6\ \mathrm{K})\left(\frac{v_e}{3\times 10^7\ \mathrm{cm\ s^{-1}}}\right)^2 \tag{3.24}$$

for a solar composition star with H and He completely ionized. Otherwise, the atmosphere behaves like a static corona of temperature T_X with density exponentially decreasing with increasing distance from the companion's photosphere. In this case, outward motion begins only high enough above the stellar surface where the thermal energy is sufficient to escape the star; hence the exponential dependence of the mass outflow on R_s/R. Escape velocities for LMXB companions considered in this paper ($m \leq 0.1\ M_\odot$) require $T_X > 10^6$ K. Such high temperatures may not be easy to obtain in a cosmic abundance companion (even with some heavy-element enrichment), particularly when the X-ray flux contains a large fraction of soft photons, as observed in several LMXBs (Vrtilek and Halpern 1985; Vrtilek *et al.* 1986). For spherical geometry, a wind driven by X-rays alone may therefore be relatively small for a degenerate star with mass $m \simeq 10^{-1}\ M_\odot$ in a typical LMXB.

In this paper and elsewhere (Kluźniak *et al.* 1988*b*, *c*, *d*), we argue for the possibility of the emission of soft ($\sim$1 MeV) γ-rays from the vicinity of neutron stars in LMXBs. We will show below that the presence of a soft γ-ray component of large intensity in the radiation reaching the companion will make it impossible to have any kind of hydrostatic corona above the depth corresponding to $P_{\min}$, essentially because equation (3.24) is automatically satisfied with $R \simeq R_s$ whenever a substantial illumination by γ-rays occurs.

We first examine the case of X-ray illumination. In Figure 1*a* we draw as functions of the temperature the curves for the negative cooling per unit mass and per unit pressure, $\Lambda_{H,X}/P = \alpha(T - T_R)/2m_p T^{3/2}$ and for the X-ray heating per unit mass and per unit pressure, $(1/P)\Lambda_X(1 - T/T_X)$ for four different values of the pressure. At a given value of the pressure, the intersection of the heating line $\Lambda_{H,X}/P$ with the curve associated to the cooling rate gives the static, thermal equilibrium as a function of the temperature T. Figure 1*b* depicts the thermal equilibrium (P, T)-curve. A single value of T is obtained for values of the pressure given by, e.g., P_1 and P_3. A value of the pressure equal to P_2 would correspond to three different values of T if a thermal instability did not occur at the minimum of the (P, T)-curve. At the minimum the pressure is $P_{\min}$ (from Fig. 1*a* we see that for $T_X \gg T_R$, $P_{\min} = P_0$ computed at $T = 3T_R$); beyond this point $dP/dT > 0$, so that the part of the equilibrium curve for $3T_R \leq T \leq T_X$ is unphysical. The dotted line of Figure 1*b* represents the response of an atmosphere under the effect of X-ray illumination. The temperature changes discontinuously

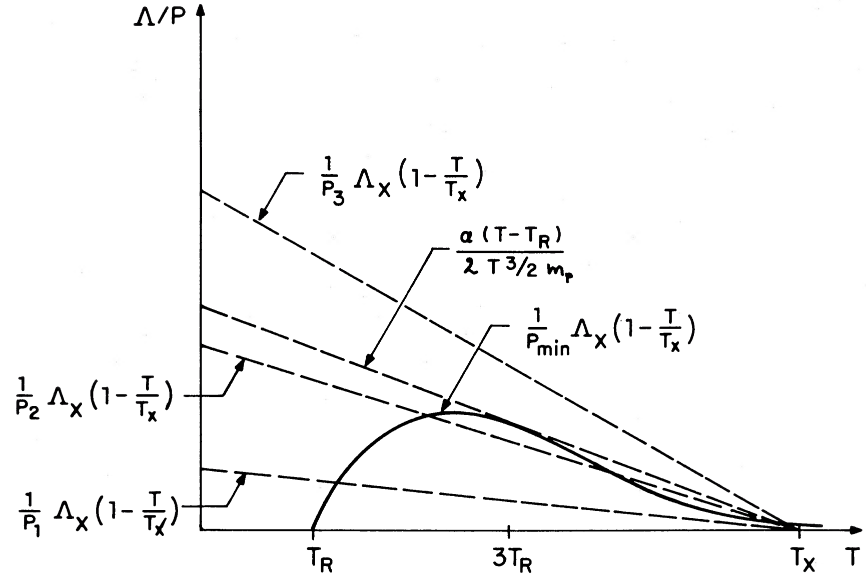

FIG. 1*a*

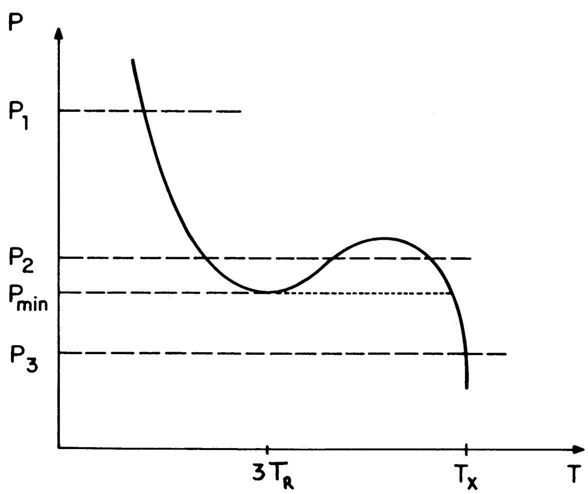

FIG. 1*b*

FIG. 1.—Thermal equilibrium in a stellar atmosphere irradiated by X-rays with a heating per unit mass $\Lambda_H = \Lambda_X(1 - T/T_X)$, where T_X is the temperature for which the absorbing heavy elements become fully ionized. The cooling per unit mass is $\Lambda_C = \alpha P(T - T_R)/2m_p T^{3/2}$, with T_R the radiation temperature of the underlying photosphere. (*a*) Λ_C/P and Λ_H/P as functions of the temperature for several values of the pressure. (*b*) Locus of the points of intersection of (*a*) on the temperature-pressure plane for which $\Lambda_H = \Lambda_C$.

from T_R to $\sim T_X$, and if the final value of the temperature is not large enough, equation (3.15) is not satisfied for a value of the sonic radius close to the photospheric radius. Thus, for a relatively low T_X ($T_X \gtrsim 10^6$ K for a solar composition atmosphere irradiated with a typical LMXB X-ray spectrum), the corona is essentially static.

If a flux of soft γ-rays is part of the incident radiation, the above conclusions need to be changed. Figures 2*a* and 2*b* show the curves analogous to those in Figures 1*a* and 1*b*, upon including a substantial heating rate per unit mass $L_{H,\gamma}$ with the assumption that the equilibrium temperature under γ-ray illumination T_γ satisfies $T_\gamma \rightarrow \infty$. The heavy-line curve in Figure 2*b* corresponds to Figure 2*a*. The dashed lines correspond to different cases that many occur depending on the values of physical quantities governing the evolution of the LMXB. A new minimum, at P_γ, appears in the (P, T)-curve, such that no thermal equilibrium without gas outflow is possible for values of the pressure smaller than $P_* \equiv \min(P_{\min}, P_\gamma)$. When $T_X \gg T_R$, then $P_* = P_{\min}$ as long as

$$\Lambda_\gamma \geq \frac{\sqrt{3}}{2}\left(\frac{T_R}{T_X}\right)^{1/2} \Lambda_X . \tag{3.25}$$

In this case, while the momentum flux is governed by the X-ray heating, a substantial wind will form regardless of the value of T_X because thermal equilibrium at a pressure smaller than $P_{\min}$ is impossible. Thus γ-rays are probably crucial to wind formation because the condition of equation (3.23) may not be easily satisfied for the systems discussed in this paper subject to X-ray illumination alone.[6]

[6] We note that a magnetic field of order $B \gtrsim 10^{-4}$ G present in the outer atmosphere of the companion can make the γ-ray energy deposition local at depths where a substantial wind can be formed.

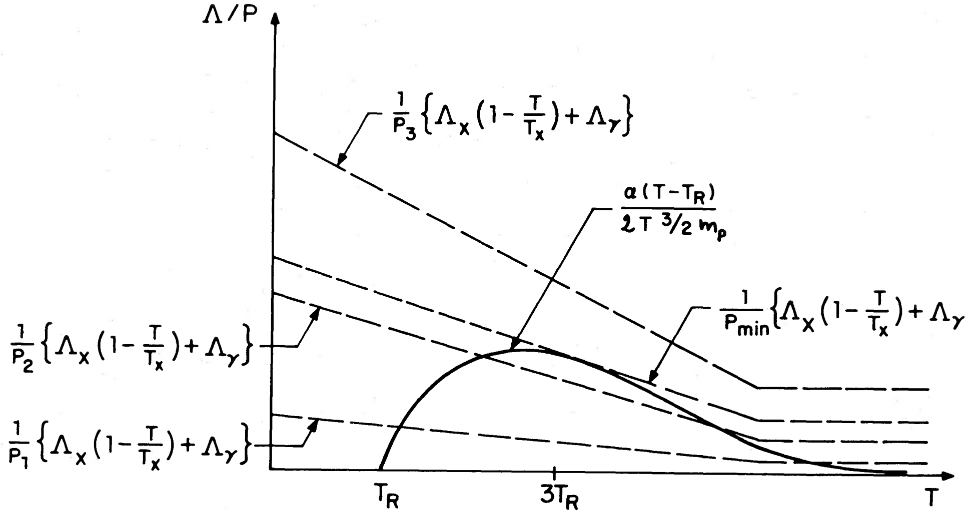

FIG. 2a

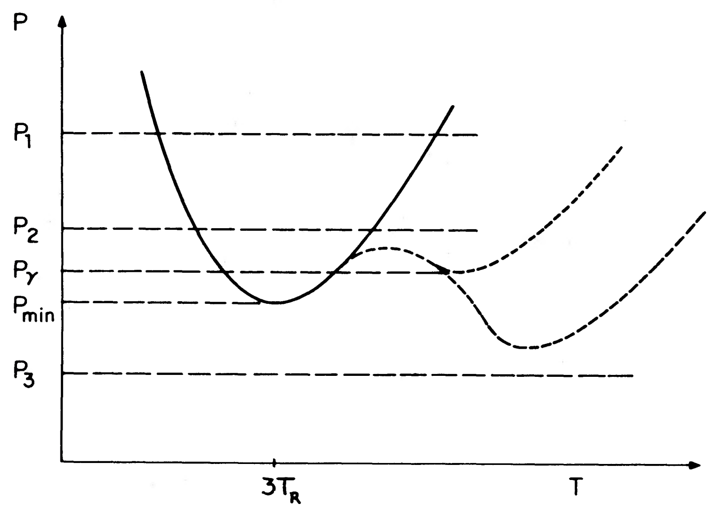

FIG. 2b

FIG. 2.—Same as Fig. 1, but with the inclusion of a substantial γ-ray heating so that $\Lambda_H = \Lambda_\gamma + \Lambda_X(1 - T/T_X)$. The solid curve in (b) corresponds to (a), while the dashed curves follow from the condition $\Lambda_H = \Lambda_C$ for parameters different from those of (a).

Upon substituting in equation (3.7) a dimensionless heating function η which takes into account both X-ray and γ-ray heating, i.e.,

$$\eta = \eta_X(1 - \xi_X x^2) + \eta_\gamma(1 - \xi_\gamma x^2) \tag{3.26}$$

for $\xi_\gamma < \xi_X$, we easily see that whenever

$$\eta_\gamma \geq \frac{1 - \xi_X}{(\xi_X)^{1/2}(\xi_X - \xi_\gamma)}, \tag{3.27}$$

the sonic point occurs at $R_s \simeq R$. In this case only the contribution from the γ-ray heating is relevant, and the wind is sufficiently hot to escape the star with a velocity at the sonic point given by $v_s = (R_s \Lambda_\gamma/3)^{1/3}$. We note that in this case only ξ_γ must satisfy $\xi_\gamma \leq 1$.

In general, one may write an expression for $|\dot{m}_w|$ in terms of $P_{\min}$ of equation (3.3). Basko *et al.* (1977) show, in fact, that under some conditions (see below) one can write the maximum possible $|\dot{m}_w|$, which arises for $R_s \simeq R$, as

$$|\dot{m}_w| v_e \sim 7.2 R^2 P_{\min} . \tag{3.28}$$

Combining equation (3.28) with equation (3.16), we see that we may write, for $\eta_0 \gg 1$,

$$|\dot{m}_w| \sim \Phi R^2 P_{\min} v_e^{-1} , \tag{3.29a}$$

where $\Phi \simeq 7.2$ for the Basko *et al.* process, in which the atmosphere heats up quickly to T_X (which satisfies formula [3.23]) and then flows isothermally in a transonic flow;

$$\Phi \simeq \frac{4\pi}{\eta_0^{1/3}} \frac{\Delta P}{P_{\min}} \lesssim \frac{4\pi}{\eta_0^{1/3}} \tag{3.29b}$$

when $(R_s - R)/R \ll 1$ and $\xi \sim 0$;

$$\Phi \simeq 4\pi\left(\frac{v_e}{2v_X}\right) \frac{\Delta P}{P_{\min}} \lesssim \frac{2\pi v_e}{v_X} \tag{3.29c}$$

when $(R_s - R)/R \ll 1$ and $\eta_0^{-2/3} \ll \xi \leq 1$.

Table 1 shows the values of $|\dot{m}_w|$ calculated from equation (3.29) for the T_R values given in the table and for a value of the ratio Λ/α which corresponds to a relevant pressure $P_{\min} \sim 3 \times 10^{-10}\Lambda$ dynes cm^{-2}. In the computation of the quantities of Table 1 we have assumed that only bremsstrahlung contributes to the cooling ($T_R \sim 10^5$ K) and that the absorption of incident radiation occurs at an absorbing column density equal to 2.5 g cm^{-2}. We leave to § V the discussion of the possible and feasible values of Λ/α for the companion star. For the purposes of the next section, however, we adopt

$$|\dot{m}_w| \simeq 10^{-17} f \hat{L} \text{ g s}^{-1}\,, \tag{3.30}$$

where $\hat{L}$ is the total illumination of the companion in ergs s^{-1} and f is a number of order unity describing the efficiency of the wind excitation process. In the actual computation of the evaporative mass loss (see Table 1) one should multiply the right-hand side of equations (3.28) and (3.29) by the factor $1 + \mathscr{A}$, where $\mathscr{A}$ is the effective albedo appropriate to the radiation reaching the companion (Basko *et al.* 1977).

In order to compare equation (3.30) with the case of Her X-1, one should set $f \simeq 0.1$ for a flat X-ray spectrum between 0.1 and 30 keV and cosmic abundance composition of the companion. This value of f is obtained for Her X-1 because of the higher value of α due to its lower ($\sim 10^4$ K) photospheric temperature T_R.

A final issue relevant to Her X-1 and LMXBs that we shall discuss here is Roche flow geometry. The proximity of the Lagrangian $L1$ point to the stellar surface in that case suggests a plane-parallel geometry approximation. In this geometry, the sonic point

TABLE 1

Quantities Related to Wind-driven Evolution of LMXBs with Degenerate Companions Filling Their Roche Lobe

Parameter	0.1 $M_\odot$ Solar	0.1 $M_\odot$ He	0.03 $M_\odot$ Solar	0.03 $M_\odot$ He	0.01 $M_\odot$ Solar	0.01 $M_\odot$ He
$R/(10^9$ cm)	4	1.6	5.4	2.2	6.4	2.8
$a/(10^{10}$ cm)	2.1	0.83	4.2	1.7	7.2	3.1
$P_{\rm orb}$ (minutes)	23.6	5.8	66.7	17.2	150	42.3
$v_e/(10^7$ cm s^{-1})	8	12.7	3.8	5.9	2	3
$g/(10^5$ cm s^{-2})	8	51	1.4	8	0.3	1.6
$\hat{L}/(10^{35} L_{38}$ ergs s^{-1})	9	9.3	4.1	4.2	2	2
$\Lambda_0/(\Upsilon \times 10^{15}$ ergs s^{-1} g^{-1})	7.2	44	1.8	11	0.61	3.3
η_0/Υ	150	91.6	472	314	1301	912
T_R (10^5 K)	1.1	1.7	0.8	1.2	0.6	0.9
$P_{\min}/(S\Upsilon \times 10^6$ dynes cm^{-2})	2.9	27	0.64	4.8	0.19	1.2
$T_{X,\min}$ (10^6 K)	6.7	39	1.5	8.4	0.42	2.1
$\dot{m}_{\rm I}/(2 \times 10^{18} L_{38} S\Upsilon$ g s^{-1})	2.5	2.7	2.5	2	2	1.6
$\dot{m}_{\rm II}/(2 \times 10^{18} S\Upsilon^{2/3} L_{38}$ g s^{-1})	$0.8L_{38}^{-1/3}$	$0.9L_{38}^{-1/3}$	$0.5L_{38}^{-1/3}$	$0.44L_{38}^{-1/3}$	$0.27L_{38}^{-1/3}$	$0.24L_{38}^{-1/3}$

Note.—In the table the effect of radiation with absorbing column density $\hat{\sigma} = 2.5$ g cm^{-2} and of bremsstrahlung cooling are considered for three different masses and two compositions. Υ is the ratio of 2.5 g cm^{-2} to the effective absorbing column density which is ultimately a function of the incoming spectrum of radiation (see eq. [3.19]). $\Upsilon \sim \frac{1}{2}$ for $\frac{1}{2}$ MeV radiation, and it can be substantially larger than unity if a substantial flux of X-rays illuminates the companion star. S is the ratio of the bremsstrahlung cooling coefficient to the effective one ($S \leq 1$). Bootstrapped evolution will occur whenever $|\dot{m}| \gtrsim 2 \times 10^{18} L_{38}$ g s^{-1} (with L_{38} the neutron star's luminosity in units of 10^{38} ergs s^{-1}), depending upon the efficiency of conversion into energetic radiation near the surface of the neutron star. The symbols are defined as follows: (1) $R \equiv$ degenerate companion radius (from Zapolsky and Salpeter 1969); (2) $a \equiv$ binary separation assuming Roche lobe overflow and a neutron star mass equal to 1.4 $M_\odot$; (3) $P \equiv$ binary orbital period; (4) $v_e \equiv$ escape velocity from an isolated companion star (without Roche potential effects); (5) $g \equiv$ acceleration of gravity at the surface of the isolated companion; (6) $\hat{L} \equiv$ total illumination (in ergs s^{-1}) of the companion, assuming a neutron star luminosity equal to $L_{38} = 1$; (7) $\Lambda_0 \equiv$ heating per unit mass at zero temperature, i.e., illumination per unit area of the companion divided by the effective absorbing column density; (8) $\eta_0 \equiv 8\Lambda_0 R/3v_e^3$ (see eq. [3.13] with $R_s \simeq R$; (9) $T_R \equiv (\hat{L}/2\pi R^2 \sigma_B)^{1/4}$, the surface radiation temperature of the companion, with σ_B the Stefan-Boltzmann constant; (10) $P_{\min} = 3^{3/2}\Lambda_0 T_R^{1/2}/\alpha$ (see eq. [3.3]); (11) $T_{X,\min} \equiv$ lowest allowed value of the temperature for which X-rays can form a substantial excited wind; (12) $\dot{m}_{\rm I}/(2 \times 10^{18} L_{38})$ g s^{-1}, where $\dot{m}_{\rm I} = 7.2(1 + \mathscr{A})R^2 P_{\min} v_e^{-1}$ is the expected wind when $T_X \gtrsim T_{X,\min}$ (see eq. [3.28] multiplied by $1 + \mathscr{A}$, where $\mathscr{A} \simeq 0.4$ is the X-ray albedo; Basko *et al.* 1977); (13) $\dot{m}_{\rm II}/(2 \times 10^{18} L_{38})$ g s^{-1}, where $\dot{m}_{\rm II}$ is the wind expected from eq. (3.29) multiplied by $(1 + \mathscr{A})$, with $\mathscr{A} \simeq 0.2$ the effective albedo for $\frac{1}{2}$ MeV photons (Tavani 1989). Note that bootstrapping occurs when the values in the last two rows are greater than unity. For the last row, that happens for a certain value $L_{38} = L_0 < 1$.

condition, equation (3.4), is modified as follows (London and Flannery 1982):

$$\Lambda_s = -\frac{3}{2} v_s g_s , \tag{3.31}$$

with g_s the gravity acceleration at the sonic point. Equation (3.31) shows that the sonic point always occurs on that side of the Lagrangian $L1$ point which is closer to the neutron star. Let us consider the case in which only γ-rays can effectively illuminate the $L1$ region. It turns out that for very large Λ_γ, radiative cooling (η' of eq. [3.11] cannot be ignored because, when $\xi \simeq 0$, eq. [3.31] means, essentially, $\Lambda_s \simeq 0$). This is a situation where the sonic velocity is very low and the sonic point must be close to the photosphere. Since in this case the overflowing gas is on the accelerating side of the gravitational well, a substantial wind will be formed. Upon rewriting equation (3.31) as

$$\Lambda_\gamma + \frac{\alpha}{m_p} \frac{P}{T^{3/2}} (T_R - T) = \frac{3}{2} v_s |g_s| \tag{3.32}$$

and using $\lambda = \rho_s v_s$, we find, for a completely ionized atmosphere with cosmic abundance,

$$\lambda = \frac{1}{2\alpha} \left(\frac{10}{3}\right)^{1/2} \left(\frac{1}{km_p}\right)^{1/2} \frac{\Lambda_\gamma - (3/2) v_s |g_s|}{1 - 10kT_R/3m_p v_s^2} , \tag{3.33}$$

which yields, when $\Lambda_0/v_s g_s \gg 1$,

$$|\dot{m}_w| \simeq 3 \times 10^{-17} \hat{L} \text{ g s}^{-1} \tag{3.34}$$

for an absorption column of 2.5 g cm^{-2}. Thus, Roche lobe geometry also suggests values of f of order unity (even though the wind velocity is quite low in this case, and the plane-parallel approximation may turn out to be too idealized).

IV. SELF-EXCITED WINDS, ACCRETION, AND LATE EVOLUTION OF LOW-MASS X-RAY BINARIES

In this section we show that, if an equation of the type (3.30) or (3.34) exists, self-excited (bootstrapped) winds can be sustained as long as the mass of a Roche lobe filling companion is larger than some critical mass m_c. If, e.g., $m_c \lesssim 0.04\ M_\odot$, one should observe relatively more sources at high X-ray luminosity than at low luminosities, and a decrease of the LMXB lifetime. The LMXB luminosity statistics, as well as the millisecond pulsar statistics, may be strongly influenced by the bootstrapped evolutionary phase. To see how bootstrapping is effective, we describe the binary evolution of an LMXB driven by a self-excited wind from a low-mass degenerate companion.

To evolve, while the companion star is still filling its Roche lobe, an effective loss of orbital angular momentum is needed. This loss can occur either by some external torque or by some wind-carried loss. If J is the orbital angular momentum of the system, then, to lowest order in m/M,

$$h = 1 + \frac{\dot{a}m}{2\dot{m}a} , \tag{4.1}$$

where we define h as

$$h = \frac{\dot{J}/J}{\dot{m}/m} . \tag{4.2}$$

The quantity h may have contributions from matter flow as well as from torques. For example, when gravitational radiation (GR) is the main mechanism for angular momentum loss, then (Faulkner 1971)

$$h_{\rm GR} = 1.7 \times 10^{-6} \frac{m^2 M^2}{|\dot{m}|} \left(\frac{10^{10}\ \text{cm}}{a}\right)^4 , \tag{4.3}$$

where m, M are expressed in solar units and $\dot{m}$ is in units of $M_\odot$ yr^{-1}. Upon substituting $|\dot{m}| \sim \dot{M}$ from equation (1.1), we find $h_{\rm GR} \propto (m/M)^{-2}$. When, alternatively, magnetic braking (MB) of a companion wind of velocity v is the main mechanism for angular momentum loss, then (Rappaport, Verbunt, and Joss 1983)

$$h_{\rm MB} = 1.2 \times 10^{-7} \frac{m^4 M^{-1/2}}{\dot{m}} \left(\frac{v}{1.3 \times 10^7\ \text{cm s}^{-1}}\right)^{3.7} \left(\frac{a}{10^{10}\ \text{cm}}\right)^{1/2} . \tag{4.4}$$

For $|\dot{m}| \simeq 10^{-8}\ M_\odot$ yr^{-1}, both $h_{\rm GR}$ and $h_{\rm MB}$ become much less than unity when m drops below 0.1 $M_\odot$. Then, in order to have high accretion rates, another source for h is needed.

If matter leaves the system through the Lagrangian $L1$ point in such a way that a fraction β of $|\dot{m}|$ remains in the system, then, if the escaping gas does not acquire any additional specific angular momentum from the binary system, $h_M = 1 - \beta$, and equation (4.1) yields

$$\frac{\dot{a}}{a} = -2\beta \frac{\dot{m}}{m} + \hat{G} , \tag{4.5}$$

with $\hat{G} = (2\dot{m}/m)(h_{GR} + h_{MB}$ + any other non–$L1$ angular momentum loss). Therefore,

$$-\frac{\dot{\Psi}}{\Psi} = \left(\frac{1}{3} + n - 2\beta\right)\frac{\dot{m}}{m} + \hat{G} , \tag{4.6}$$

where Ψ is the ratio of the stellar radius R and the Roche lobe radius R_L, $\Psi = R/R_L$, and where we have assumed for the radius R of the companion star the relation

$$R = Km^{-n} \tag{4.7}$$

with a constant K. When the total angular momentum is conserved, $\hat{G} = 0$, and for constant β and a degenerate star ($n = \frac{1}{3}$), we have

$$a \propto m^{-2\beta} \tag{4.8}$$

and

$$\frac{R_L}{R} \equiv \Psi^{-1} \propto m^{2(1/3-\beta)} . \tag{4.9}$$

Constant Roche lobe overflow requires $\beta = \frac{1}{3}$; when $\beta > \frac{1}{3}$, the companion will contract inside its Roche lobe (Kluźniak *et al.* 1988*a*). The latter case may have been realized during the evolution[7] of the binary system PSR 1957+20 for which $\Psi < 1$ at present (Fruchter, Stinebring, and Taylor 1988). If $\beta < \frac{1}{3}$, a very strong mass loss may occur, yielding the tidal disruption of the companion when it reaches $\Psi = 1$ from below (Ruderman and Shaham 1983). We note that the value $\beta = \frac{1}{3}$ is inconsistent with binary evolution driven by angular momemtum loss caused by gravitational radiation. When $\hat{G}$ is nonzero [for example, for the gravitational torques we have $\hat{G} = 4 \times 10^{-2}(M/M_\odot)^2(m/M_\odot)(a/10^{10}\ \mathrm{cm})^{-4}\ \mathrm{yr}^{-1} \propto m^{11/3}$ for Roche lobe overflow], steady Roche lobe overflow implies from equation (4.6) (Rappaport *et al.* 1987)

$$\frac{\dot{m}}{m} = \frac{\hat{G}}{2(\beta - 1/3)} , \tag{4.10}$$

which suggests a possible instability for $\beta \lesssim \frac{1}{3}$ if gravitational radiation were the only cause of angular momentum loss. From this example, one can realize how self-excited winds may drastically change the evolution of binaries which otherwise would have been driven by gravitational radiation.

When matter is accreted onto the neutron star at a rate $\dot{M} = \beta|\dot{m}|$, the accretion-powered luminosity is

$$L = \xi\beta|\dot{m}|c^2 , \tag{4.11}$$

where ξ is the efficiency for energy conversion into energetic radiation near the surface of the neutron star. This luminosity may be attenuated when it reaches the companion; let the attenuation coefficient be χ, such that

$$\hat{L} = \frac{1}{4}\,\xi\chi\beta\Psi^2 c^2|\dot{m}|\,\frac{R_L^2}{a^2} \equiv \frac{1}{4}\,L\Psi^2\,\frac{R_L^2}{a^2}\,\chi . \tag{4.12}$$

The effective illumination of the companion $\hat{L} = \hat{L}(\dot{m}, m)$ is a function of the solid angle occupied by the companion, as well as of scattering and absorptive phenomena which can occur in the outer disk. The geometrical factor $\hat{Q}^2 \equiv \frac{1}{4}\Psi^2 R_L^2 a^{-2}$ depends on the mass of a degenerate companion m as follows:

$$\frac{1}{4}\Psi^2\frac{R_L^2}{a^2} = \begin{cases} \dfrac{1}{9}\left(\dfrac{m}{3M}\right)^{2/3} & \text{for } \Psi = 1 , \quad (4.13\text{a}) \\[2ex] \dfrac{1}{9}\left(\dfrac{m_0}{3M}\right)^{2/3}\left(\dfrac{m}{m_0}\right)^{2(2\beta-1/3)} = \dfrac{K^2}{4a_0 m_0^{2\beta}}\, m^{2(2\beta-1/3)} & \text{for } \Psi < 1 . \quad (4.13\text{b}) \end{cases}$$

In equation (4.13b) it is assumed that the companion radius satisfies equation (4.7) and a_0 and m_0 are some representative values for the evolution at the present value of β. Only when $\beta > \frac{1}{3}$ does equation (4.13) represent a fully consistent description. (When $\beta < \frac{1}{3}$, the value $\Psi = 1$ could be reached, which requires special care; see eq. [4.10].) Upon substituting equation (4.13) in equation (3.30), we find $|\dot{m}_w|$. The relationship between $|\dot{m}_w|$ and $|\dot{m}|$ depends on the effective mass loss from the binary system. The escaping wind satisfies

$$|\dot{m}_{esc}| \leq |\dot{m}_w| . \tag{4.14}$$

The right-hand side of equation (4.14) is only an upper limit to the $\dot{m}$ which escapes the binary $|\dot{m}_{esc}|$ because of tidal recapture of the evaporative wind material. When the sum of the final wind velocity v_w added to the orbital velocity exceeds the effective escape velocity from the binary, i.e.,

$$\frac{v_w^2}{2} > \frac{GM}{2a}(3 - 2\sqrt{2}) , \tag{4.15}$$

[7] The binary system PSR 1957+20 provides the first observational evidence for an evaporative wind caused by radiation from a millisecond pulsar.

some of the gas outflowing the companion can escape. Assuming the validity of equation (4.15), we will set $|\dot{m}_{esc}| = |\dot{m}_w| \simeq (1 - \beta)|\dot{m}|$. Under Roche lobe overflow conditions, it follows that $|\dot{m}_w| \simeq \frac{2}{3}|\dot{m}|$, and the wind mass-loss rate has to be twice the mass loss which contributes to accretion. When equation (4.15) is not satisfied, the relationship between $|\dot{m}_w|$ and $|\dot{m}|$ is less clear, since part of $|\dot{m}_w|$ may be returned to the companion, and only the remainder will be accreted. Thus we can write (see eq. [3.34]; cgs units are used)

$$10^{-17} f \hat{L} \equiv |\dot{m}_w| = \zeta |\dot{m}| , \tag{4.16}$$

where $\zeta \simeq 1 - \beta$ when equation (4.15) holds, and $\zeta \geq 1$ otherwise.

Figure 3 shows the behavior of the quantity $L\chi$ as a function of $|\dot{m}|$ obtained from equation (4.12) and (4.16). Equation (4.16) is plotted, for constant f, as a straight line through the origin, whose slope depends on m. When $f \propto \hat{L}^{-1/3}$ for large values of $\hat{L}$, e.g., for a large γ-ray illumination, this curve becomes steeper. Equation (4.12) saturates, for $|\dot{m}| \to 0$, to a finite value L_d, the spin-down luminosity of the neutron star (RST). For large $|\dot{m}|$, $L\chi$ may saturate close to the Eddington luminosity, $1.3 \times 10^{38}(M/M_\odot)$ ergs s^{-1}. Three lines derived from equation (4.16) are plotted along with the stable points of the bootstrapping process. Cases (*a*) and (*b*) yield a luminosity of the order of the Eddington luminosity times the attenuation χ of the incident radiation flux. For case (*c*) and for lines with larger slope, the luminosity and $|\dot{m}|$ drop by several orders of magnitude. The range of luminosities between L_1 and L_2 in Figure 3 is not realized.

Case (*c*) determines the critical mass m_c of the companion at which bootstrapping cannot be sustained anymore. Since

$$\hat{Q}^2 = 5 \times 10^{-4} f^{-1} (\xi/0.2)^{-1} \chi^{-1} \beta^{-1} \zeta , \tag{4.17}$$

from equation (4.13a) we find, for Roche lobe overflow, and for a constant f of equation (3.30),

$$m_c \simeq (2.2 \times 10^{-2}\ M_\odot) f^{-3/2} \chi^{-3/2} \left(\frac{\xi}{0.1} \frac{1/3}{\beta}\right)^{-3/2} \left(\frac{M}{1.4\ M_\odot}\right) . \tag{4.18}$$

A critical mass may occur even when the curves representing equation (4.16) steepen above some value of $\hat{L}$ because $f \propto \hat{L}^{-\alpha}$ with $\alpha > 0$. In this case, the value of m_c will, of course, be higher and the value of L_1 lower. Thus, a larger range of luminosities, down to L_1, will be allowed by the bootstrapping mechanism. Note, however, that f may itself depend on m and therefore the f in equation (4.18) will depend on m_c. For example, if equation (3.28) holds, $f_c \propto m_c^{-2/3}$ for $2 \times 10^{-2}\ M_\odot \leq m_c \leq 10^{-1}\ M_\odot$ and equation (4.18) would be meaningless. In this case bootstrapping will always occur for a luminosity L_0 satisfying $L_2 \leq L_0 \leq L_1$ if the photospheric temperature of the companion star is sufficiently large. For a substantial illumination for which $\eta_0 \gg 1$, if v_e in equation (3.28) is replaced by $(R_s \Lambda_0/3)^{1/3} \propto m^{1/3}$, then $f_c \propto m_c^{-1}$, and equation (4.18) should be modified accordingly. At present, not enough is known to determine f to high accuracy, but a value for the critical mass of the order of a few hundredths of a solar mass seems plausible.

If, during the bootstrapped phase with Roche lobe overflow, the mass of the secondary m becomes smaller than m_c, the radiation-driven wind is no longer important and gravitational radiation may (again) be the main source of angular momentum loss from the binary system. As m reaches m_c from above, accretion may become unstable (perhaps with a fluctuating L_X), and then its rate may drop to the value appropriate to a gravitational radiation–driven LMXB. This case is considered schematically for a hydrogen white dwarf secondary in Figure 4 and for a helium one in Figure 5 for $m_c = 10^{-2}\ M_\odot$. In the former case there is a rapid

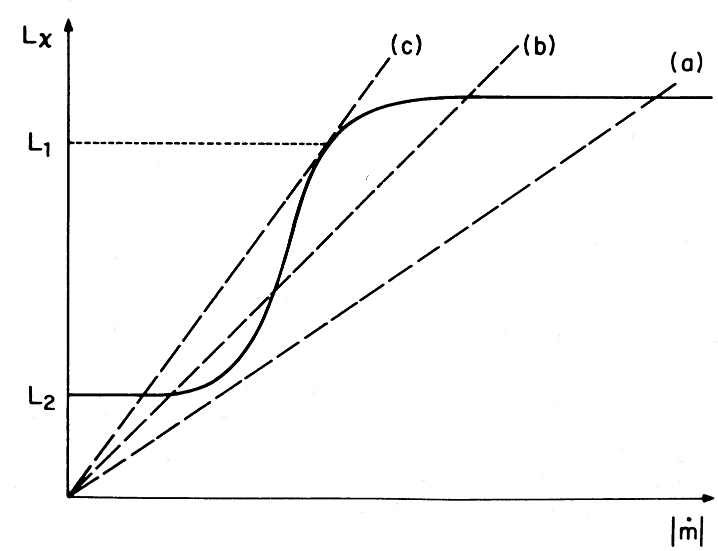

FIG. 3.—Bootstrapping points for various companion masses, $m_{(a)} > m_{(b)} > m_{(c)} \equiv m_c$. The quantity $L\chi$ is the attenuated luminosity reaching the companion. The winding curve very schematically depicts eq. (4.13), while the straight lines depict eq. (4.17). Under realistic conditions, curves (*a*), (*b*), and (*c*) may steepen upward for larger values of $L\chi$ or may be very close to each other because of insensitivity upon the companion mass.

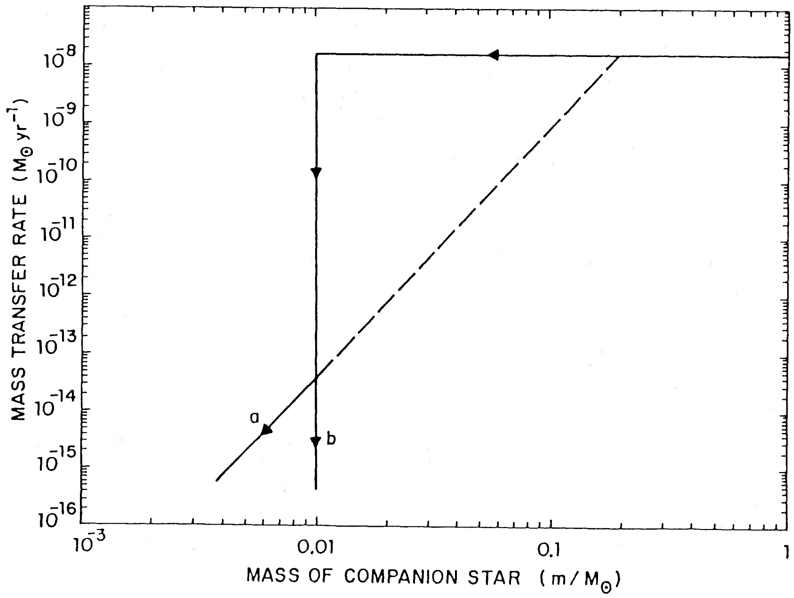

FIG. 4.—Schematic behavior of the MeV γ-ray–induced mass transfer rate in solar masses per year versus the mass of a H white dwarf companion. The dotted line represents the mass transfer rate due to gravitational radiation only. The turnoff transition is shown for a critical mass $m = 10^{-2}\ M_\odot$ (cf. eq. [4.18]). In the bifurcation the vertical evolution (*b*) is for a neutron star spun up to a near-millisecond period; the diagonal path (*a*) is that of a more slowly spinning neutron star. The possibilities are discussed in § VI.

drop in accretion rate and X-ray luminosity by a factor of order 10^5, and in the latter by about 10^4. After this relatively rapid transition, the evolution of the LMXB proceeds by one of two possible alternatives.

a) If the new Alfvén radius still lies inside the light-cylinder radius of the spun-up neutron star, i.e., the neutron star's spin P satisfies the inequality

$$P > (2 \times 10^{-4}\ \mathrm{s})\left(\frac{B_s}{5 \times 10^8}\right)^{4/7} \dot{M}_{18}^{-2/7}\ , \tag{4.19}$$

the accretion continues at the rates shown in Figures 4 and 5 with corresponding L_X well below 10^{34} ergs s^{-1} (see eq. [1.1]). These LMXBs would have an X-ray luminosity indistinguishable from those of cataclysmic variables, but might have a spectrum harder than the typical UV or optical ones (cf., e.g., Córdova and Mason 1983).

b) If the spun-up neutron star period is sufficiently short that inequality (4.19) fails, accretion is completely quenched. For $B_s \simeq 5 \times 10^8$ G and $\dot{M}_{18} \simeq 10^{-4}$, this occurs for $P \lesssim 3 \times 10^{-3}$ s. If the neutron star spin-up has reached its steady state, the

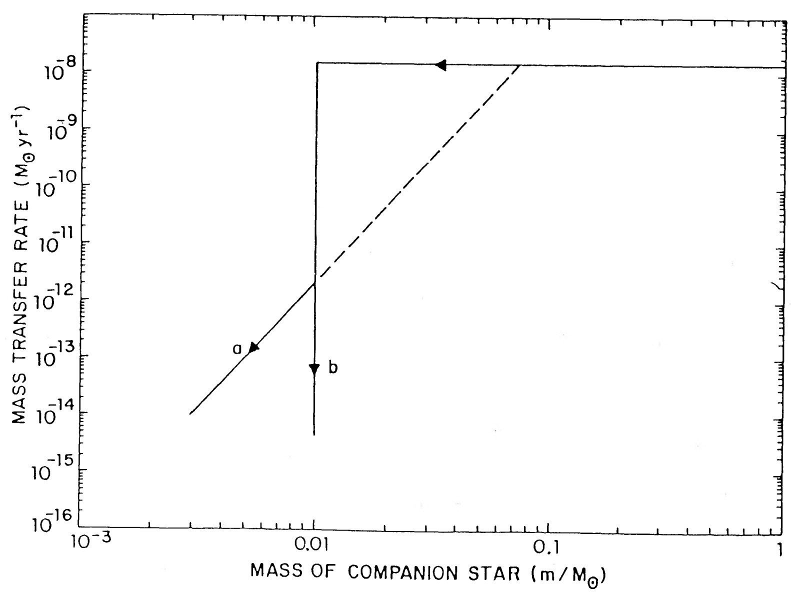

FIG. 5.—Same as Fig. 4, for a He white dwarf companion star

accretion rate which sets this turnoff satisfies the inequality

$$\dot{M}_{14} \lesssim 40 P_{\rm ms}^{-7/2} . \tag{4.20}$$

For rapidly spinning "millisecond" pulsars the corotating magnetosphere radius (R_A) is limited not by a surrounding accretion disk but by its light cylinder. Such neutron stars will therefore radiate as if they were in a vacuum. Unless their magnetic fields are axially symmetric and exactly aligned, they will radiate an electromagnetic wave whose radiation pressure may exceed that which can be balanced by an accretion disk.[8] For a surface dipole field B_s the pressure $p_{\rm md}$ needed to contain the magnetic dipole radiation when $r \gg cP/2\pi$ is

$$p_{\rm md} \sim \frac{B_s R^3}{4\pi r^2} \left(\frac{2\pi}{cP}\right)^4 , \tag{4.21}$$

while the accreting matter pressure p_m falls with distance as

$$p_m \sim r^{-5/2} . \tag{4.22}$$

Thus, if $p_{\rm md}$ is greater than p_m at $r = cP/2\pi$, it remains greater at all larger radii. Therefore the neutron star's dipole radiation pressure may completely turn off any accretion, and any remaining L_γ (or L_X) is from the millisecond pulsar itself. The effect of such radiation on the secondary, especially its ability to evaporate it, is considered in RST.

When bootstrapping is controlled by γ-rays produced near R_A, m_c will also depend on the stellar surface magnetic field, B_s. Furthermore, some of the soft X-ray radiation may, actually, also form in the same region. In this case, ξ of equation (4.11) will depend on $|\dot{m}|$ ($\xi \propto R_A^{-1} \propto |\dot{m}|^{2/7}$ for a matter-dominated inner accretion disk) and the critical mass will depend on the magnetic field (e.g., $m_c \propto B_s^{6/7}$ for $f = $ const.). This dependence indicates the possibility of four different regimes in the evolution of LMXBs with light secondaries, as indicated schematically in Figure 6.

1. B_s is initially large enough that m_c exceeds m. For a degenerate dwarf secondary with $m > 10^{-1}\, M_\odot$, gravitational radiation alone can keep $\dot{M} \simeq 10^{18}$ g s^{-1} and L_X at the Eddington limit.
2. As m diminishes to $\sim$(3–5) $\times 10^{-2}\, M_\odot$, $|\dot{m}|$ and L_X fall by about 2 orders of magnitude. This would last for about $m/\dot{m} \simeq 10^7$ yr. (This is also the time scale sometimes estimated for magnetic field decay of a neutron star.)
3. If B_s drops toward 3×10^8 G, m_c also diminishes, so that $m > m_c$ and equation (4.16) will, at some moment, be satisfied. At that time, the accretion rate and L_X will rise to L_1 of Figure 3.
4. After an additional interval of $m/\dot{M} \simeq 3 \times 10^6$ yr, m will drop below the new m_c and accretion will be enormously diminished or cease completely as discussed in alternatives *a* and *b* above.

We speculate that some LMXBs which are observed with $L_X \sim 10^{36}$–10^{37} ergs s^{-1} are in phase 2 of the above evolution scenario, which is not a very long-lived stage and so may not involve many LMXBs with very low mass secondaries. Large observed fluctuations of $|\dot{m}|$ in several LMXBs may be related to variations in the effective optical depth for absorption of the X-rays, γ-rays, and $e^\pm$ wind on their way from the inner edge of the accretion disk to the companion star. An increase in such absorption could suppress radiation-sustained evaporative winds from the secondary star. The accretion disk feeding would then be driven by

[8] Even after accretion turnoff, the pulsar radiation does not exert any stress on an infinitesimally thin flat disk whose spin axis is exactly parallel to that of the pulsar.

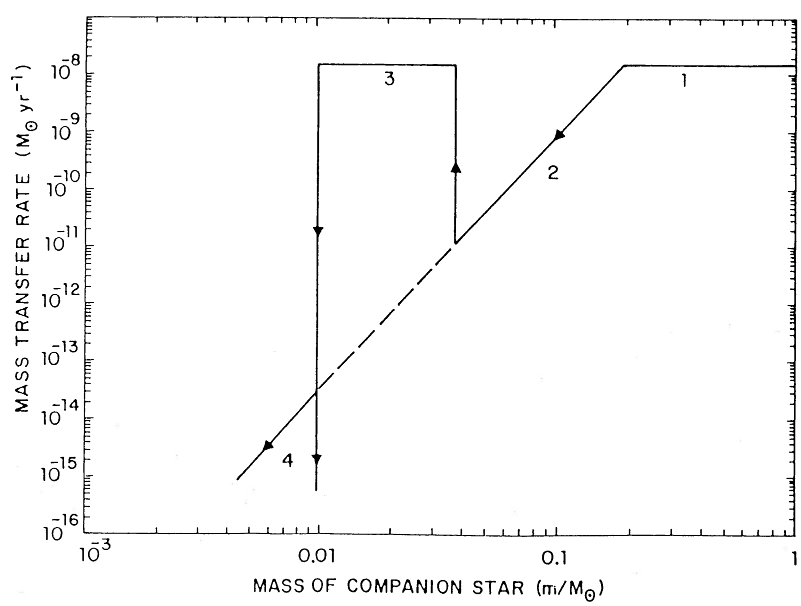

FIG. 6.—The four possible stages in the evolution of LMXBs discussed in § IV

gravitational radiation until the optical depth drops and allows the X-rays and/or the γ-ray/$e^{\pm}$ wind from the primary to be effective again in binding up a large $|\dot{m}|$ from the secondary. (We note that a relatively long fluctuation time scale may be produced by any bootstrap mechanism which couples an intervening absorption of radiation with $|\dot{m}|$.)

V. EFFECTS ON WINDS OF STELLAR COMPOSITION AND OF RADIATION INTENSITY AND SPECTRUM

For neutron star luminosities close to the Eddington limit and for binaries containing degenerate companions filling their Roche lobes, typical photospheric temperatures T_R are close to 10^5 K. This high-temperature outer atmosphere of the companion is then characterized by a cooling rate smaller than the cooling rate due to H and He recombination and heavy-element line cooling. The value for the relevant pressure $P_{\rm min}$ (eq. [3.3]) is therefore large compared with the Her X-1 value (London, McCray, and Auer 1981). However, as we see from Table 1, only under optimal conditions of hydrodynamical acceleration will the external radiation penetrating to the depth $\sim m_p/\sigma_{\rm T}$ (with $\sigma_{\rm T}$ the Thomson cross section) produce values of mass loss suitable to a bootstrapping mechanism for secondary masses as low as 0.03 or even 0.01 $M_\odot$. How could bootstrapped self-excited winds be maintained for small secondary masses?

For luminosities close to the Eddington luminosity, the values of Table 1 imply, by equation (3.5), a region for which $\eta_0 \gg 1$, so that from equation (3.15) we would only be able to limit the wind speed v_s (thus raising $\dot{M} \propto P_s R_s^2/v_s$) by having external radiation with an adequate value of T_X ($\xi \simeq 1$). Under these conditions we also expect equation (3.28) to hold. Let us consider X-rays first.

The most effective absorbers are C, N, and particularly O. There is some evidence for the presence of Ne and Mg in cataclysmic variables (see, e.g., Starrfield, Sparks, and Truran 1986) which also makes possible the existence of O-Ne–rich companions (Iben and Tutukov 1984). An increased abundance of such metals in the companion, from enrichment through primary nova outbursts or supernova events or because of suitable main-sequence evolution, may yield a value of $P_{\rm min}$ larger by a factor of several compared with the cosmic abundance case, as long as the effective particle mass μ does not become too large. (Otherwise the right-hand side of eq. [3.24] should be multiplied by μ/m_p and a higher T_X have to be achieved to form a wind.)

While Λ, hence $P_{\rm min}$, could rise with increased metal abundance (depending on the exact ionization balance in the outer atmosphere of the companion and on the incident spectrum of radiation), the very fact that these X-rays *are* strongly absorbed may cause two problems:

1. The X-rays may be substantially attenuated before reaching the companion because of absorption by cold matter in the disk or elsewhere. The persistent lack of eclipses in LMXB observations support this view. Clearly, the presence of a very large scattering corona does not seem possible in observed LMXBs, because short-time variability of the X-ray emission up to 100 Hz is present in most of them (Hasinger 1988). However, the possible occurrence of a short-lived phase in which a corona with temperature equal to 10^6 K is present in the LMXB as a common envelope around the two stars[9] could solve the problem of X-ray absorption. The soft X-rays would be scattered in the corona surrounding the binary, and they would excite a very substantial ($\geq$ a few times 10^{-8} $M_\odot$ yr^{-1}) wind. With a size of 10^{11} cm and average number density $n \sim 10^{14}$ cm^{-3}, such a corona would degrade any γ-ray from the primary down to $\sim$20 keV and could reprocess $\gtrsim$10% of the neutron-stellar radiation down to $\sim$0.1 keV photons. These soft X-rays may be very efficient for wind excitation from the secondary. The effective capture area of the companion would increase by a factor of $\sim$4, and so would the value of $|\dot{m}|$. These LMXBs would not be quasi-periodic oscillation (QPO) sources, for the photon diffusion time across the corona would be of order 10^2 s. However, the lifetime of the system in that phase would be quite short: $|\dot{m}_w|$ values above 10^{-8} $M_\odot$ yr^{-1} may give a lifetime as short as 10^6 yr for a sufficiently small mass of the companion. Thus this phase is practically unobservable, but after its completion a neutron star binary with a very light companion, of mass m_c, would be left. The total mass contained in such a corona would be of order 3×10^{-10} $M_\odot$ and may build up by the same kind of bootstrapping processes as those exciting the secondary's wind. If the primary's accretion-driven luminosity L is somewhat larger than the Eddington luminosity $L_{\rm E}$, any large-scale corona could not have a density much above $\sim 10^{12}$ cm^{-3}; otherwise, much more than 10^{18} g s^{-1} would be ejected from the system. Only when L begins to drop below $L_{\rm E}$ can that density be exceeded. For an appropriate value of T_X the wind will escape from the star with velocity v_X just below the binary escape velocity and may sustain a corona. The $e^{\pm}$ wind formed at the magnetospheric boundary (see § VI) could provide a "starter" to this particular bootstrapping process.

2. For the values of the net heating per unit mass Λ_0 given in Table 1, the densities at which radiative cooling balances heating are so large ($\sim 10^{16}$ cm^{-3}) that it is possible that few X-rays could reach the depth where the pressure is close to $P_{\rm min}$. Thus pressures less than $P_{\rm min}$ may govern $|\dot{m}_w|$, decreasing the latter to below the values of Table 1. However, it is not inconceivable that the scale height for those densities in the presence of a strong illumination will be sufficiently small to render the coronal region thin after all. That scale height will be of order R/η, so that the optical depth $\tau \sim \rho R/\eta$ will be independent of Λ. Furthermore, illumination of the companion's atmosphere by a relatively large flux of γ-rays may favor the penetration of X-rays to larger depths. This may occur because γ-rays can bore holes in the disk and atmosphere of the companion (see § VII) and because the ionization of heavy elements throughout the corona may be substantial owing to the heating produced by reprocessed radiation. Thus this problem could be overcome as well.

We turn next to consider possible illumination by soft γ-rays (see § VI), which have an energy deposition cross section which does not depend on the state (or composition) of the companion's atmosphere. Because they penetrate deeper into the atmosphere, they may not be as effective as X-rays. We emphasized in § III the role of these γ-rays in transforming a not-hot-enough corona heated by X-rays into a wind. For the remainder of this section we will consider illumination by γ-rays alone. Even though they will be able to reach the companion better than the X-rays, γ-rays (for the same reason) give both potentially less heating per particle and higher wind velocities ($\xi \simeq 0$ effectively). Both of these effects tend to reduce $|\dot{m}_w|$ under the conditions present in the relevant LMXBs. The

[9] Perhaps because the wind itself has, in this case, a velocity smaller than the escape velocity from the system.

soft γ-rays which may be emitted from the vicinity of the neutron star may, however, reach the companion even when the softer X-rays do not. For Compton heating, $\sim$MeV γ-rays are more effective in depositing their energy than softer X-rays. We also note that the mean free path in the atmosphere for electrons scattered by the γ-rays is of the same order of magnitude as that for the γ-rays themselves, and therefore γ-ray heating tends to be nonlocal. However, even a small magnetic field ($\geq 10^{-4}$ G) present in the outer atmosphere of the companion could effectively confine the struck electrons and yield a local deposition of energy.

In the extreme case of nonlocal heating, a quasi-isothermal accelerating region may, again, be produced, allowing the use of equation (3.31) for γ-rays as well. The constant temperature of that region will be fixed naturally by comparing the rates at which heat is conducted away from outflowing gas to the rate at which γ-rays deposit their energy. This will introduce an effective v_s which may be small enough to be interesting. When localizing magnetic fields are present, however, the velocity of the sonic point will be of order $\frac{1}{2}v_e \eta_0^{1/3}$ and $|\dot{m}_w|$ will drop, so that m_c of equation (4.18) may become larger than $4 \times 10^{-2}\ M_\odot$.

VI. LOW-ENERGY γ-RAYS AND $e^\pm$ WINDS FROM LOW-MASS X-RAY BINARIES

The emission of γ-rays from a weak dipole moment neutron star qualitatively changes the expected γ-ray emission from that discussed in § II. The Alfvén radius R_A, near where the γ-ray accelerator is assumed to be, is related to the neutron star dipole moment $\tilde{\mu}$ and accretion rate $\dot{M} = LR_{ns}/GM$ by

$$R_A \simeq 3 \times 10^8 \tilde{\mu}_{30}^{4/7} L_{37}^{-2/7}\ \text{cm} \sim 10^6 \left(\frac{B_s}{5 \times 10^8}\right)^{4/7} \dot{M}_{18}^{-2/7}\ \text{cm} \tag{6.1}$$

as long as $R_A > R$ (Alpar *et al.* 1982). In equation (6.1), L_{37} is the total (X-ray) luminosity powered by accretion down to the neutron star surface (in units of 10^{37} ergs s^{-1}) and $\tilde{\mu}_{30}$ is the neutron star's magnetic moment (in 10^{30} G cm^3).

For an LMXB with $B_s \simeq 10^9$ G and $L_{37} \simeq 10$, $R_A \simeq 2R_{ns}$. This would be expected to cause two dramatic changes in the γ-ray flux from an LMXB compared with that from an X-ray pulsar with much larger $\tilde{\mu}$. First, much more power can be mobilized at the accretion disk magnetosphere interface which powers the γ-ray energy source,

$$L_\gamma(\text{LMXB}) \sim L_X/2 \sim 5 \times 10^{37}\ \text{ergs s}^{-1}\ , \tag{6.2}$$

which is a γ-ray flux much larger than that from Her X-1. Even more power could be produced there if the millisecond period neutron star happens to be spinning down at that time (Priedhorsky 1986; Tavani 1989; RST). Second, the character of the γ-ray flux must be very different. At the magnetosphere accretion disk interface of Her X-1, $B \lesssim 10^6$ G and 10^{12} eV γ-rays can escape before magnetic conversion into $e^\pm$ pairs. But, in general,

$$B(R_A) \sim \tilde{\mu} R_A^{-3} \sim \tilde{\mu}^{-5/7}\ , \tag{6.3}$$

so that the smaller dipole in an LMXB results in a larger magnetic field at the Alfvén radius. With $B_s \simeq 10^9$ G and $R_A \simeq 2R$, most of the VHE γ-rays more energetic than 10 GeV would be materialized into $e^\pm$ pairs by the magnetic field before escaping. The ultimate photon energies into which VHE γ-rays are degraded in order to escape the region where $r \sim R_A \simeq 2R$ may depend upon unknown details of how the energy ultimately emitted as γ-rays is initially injected. However, rather general arguments suggest that much of the finally emitted photon intensity may be degraded into the MeV range (Kluźniak *et al.* 1988*b*). For γ-rays above 10 GeV the local B is opaque. Between 10 GeV and 10 MeV the optical depth τ_γ for pair production on accretion-powered X-rays of energy E_X is

$$\tau_\gamma = \frac{L_X \sigma_{\gamma X}}{4\pi R_A c E_X} \simeq \frac{L_X (3/8) \sigma_T (mc^2)^2}{4\pi R_A c E_X^2 E_\gamma} \log\left[\frac{4E_\gamma E_X}{(mc^2)^2}\right], \tag{6.4}$$

as long as $E_\gamma E_X \gg (mc^2)^2$, with $\sigma_{\gamma X}$ the pair-production cross section of γ-rays on X-rays and σ_T the Thomson cross section. For $E_X \simeq 10$ keV, $E_\gamma = 10^2$ MeV, $L = 10^{38}$ ergs s^{-1} and $R_A = 3 \times 10^6$ cm, $\tau_\gamma \simeq 10^3$. This is sufficiently large that γ-rays can escape only if their energy falls below that needed to make pairs on accretion X-rays: $E_\gamma < (mc^2)^2/E_X$. For an assumed maximum $E_X \simeq 25$ keV this window opens only when E_γ falls below about 10 MeV.

The precise spectral range into which the energy of VHF γ-rays will finally be emitted depends upon a series of processes and presently unknown initial conditions. Thus, for example, a 10^{12} eV γ-ray would be expected to be emitted initially almost parallel to the local magnetic field B (possibly by inverse Compton scattering) by an electron[10] moving along B. The TeV photons in the 10^8 G field materialize as soon as their pitch angle with the magnetic field becomes as large as $\theta \sim 10^{-2}$, so this is the expected initial pitch angle of the pairs. However, since

$$B^2\theta^2/8\pi \ll L_X/4\pi R_A^2 c \tag{6.5}$$

for such values of θ, these pairs would lose their energy to inverse Compton scattering on X-ray photons before they emit synchrotron radiation. With the large values of τ_γ in equation (6.4) the new $e^\pm$ pairs will be formed at essentially the same location and pitch angle. The cascade will continue in a similar fashion until τ_γ becomes large enough for the γ-Comptonized X-rays to travel a large fraction of R_A before making the pair. At that stage, when $E_\gamma \simeq 10$ MeV, θ becomes of order unity and synchrotron losses exceed the Compton ones. The synchrotron photons are emitted at frequencies around and below 1 MeV in a range of magnetic field values. A more detailed discussion of the resulting power-law spectra for energies much below 1 MeV is given elsewhere (Kluźniak *et al.* 1988*b*, *c*). The expected high-energy power-law spectra may possibly have been observed in the low-intensity optically thin state of some LMXBs (see, e.g., XB 1916−05 [White and Swank 1982], 4U 1705−44 [Langmeier *et al.* 1987]) and in the Galactic center region (Riegler *et al.* 1985).

[10] Or by a π^0 meson produced by an accelerated proton colliding with disk matter.

In the absence of a convincingly detailed model for VHE γ-ray production, it is difficult to offer a quantitative estimate for the MeV intensity from a LMXB. Additional uncertainties in the local magnetic field where $e^{\pm}$ pairs emit synchrotron radiation and the fate of lower energy $e^{\pm}$ trapped on closed magnetic field lines (which may convert much of their energy into $\frac{1}{2}$ MeV annihilation γ-rays) make it even harder to obtain a reliable estimate of the MeV γ-ray flux. For our numerical examples we somewhat arbitrarily assume that the L_{γ} of equation (6.2) is emitted at or below around 1 MeV, typical of that sometimes observed to be emitted from the Galactic center region, which has, indeed, been attributed to LMXB sources (Kluźniak *et al.* 1988*b*, *c*). Not all LMXBs are expected to exhibit these spectra, because only a small fraction of them have appropriately low magnetic fields (and secondary masses). Section VII deals with several questions raised by the existence of a corona optically thick to X-rays which probably exists in some LMXBs.

We further note that the optical depth for the scattering of X-rays from a hot polar cap on the spinning neutron star surface by these $e^{\pm}$ pairs could be as large as

$$\tau_X \simeq \left(\frac{L_{\gamma}}{1\ \text{MeV}}\right)\frac{\sigma_{\text{T}}}{4\pi R_{\text{A}}\, c} \simeq 50 L_{38}\ . \tag{6.6}$$

It could then take a time $t \sim \tau_X(R_{\text{A}}/c) > 3 \times 10^{-3}$ s for an X-ray to diffuse out from the stellar surface through the surrounding $e^{\pm}$ flux, perhaps enough to smooth out in time and especially in angle a signal whose time scale is of the order of milliseconds (as well as to give some energy degradation of the most energetic X-rays) but not to affect significantly the longer pulse time scale of quasi-periodic objects (Brainerd and Lamb 1987; Bussard *et al.* 1988). With so many scatterings (τ_X^2 per X-ray photon) the $e^{\pm}$ cloud would cool down to the X-ray photon temperature when Eddington X-ray fluxes are prevalent. At sub-Eddington luminosities, however, the cloud may remain quite hot. A hard synchrotron spectrum may not escape when L_{38} is of order unity, emerging intact only when L_{38} drops to several percent.

We note that there is considerable inferential evidence (White and Holt 1982; White, Stella, and Parmar 1988; Stollman *et al.* 1987; Mason 1986) for an X-ray scattering corona, of optical depth $\tau_X \simeq 3$–10, around the neutron star in several LMXBs. The canonical view attributes the corona to X-ray–sustained evaporation of disk matter (White and Holt 1982). In the above discussion, at least part of the corona cloud be the $e^{\pm}$ could whose source is near R_{A}. We consider next the question of whether these (various) coronae may prevent γ-ray beams originating relatively near the primary from reaching the secondary.

VII. WILL THE HIGH-ENERGY RADIATION REACH THE COMPANION?

There are three classes of observations which have been interpreted as providing evidence for the presence of coronae in LMXBs. Perhaps the most compelling one is the absence of sharp, total eclipses. If one argues that LMXBs are observed preferentially from directions well outside the binary plane, one must postulate the presence of obscuring material in that plane. The diffuse eclipses observed in some LMXBs could then be thought of as viewed through regions where the anisotropic obscuring material is only partially obscuring. There seems to be no evidence, however, for reprocessing of X-rays on such large scales in LMXBs. A less extreme model (White and Holt 1982) is based on the notion that the eclipses are not sharp and complete because the observed X-ray source is more extended than the binary companion. This would be true if a Compton-thick scattering corona of dimension exceeding that of the companion scatters the X-rays from a source which it encloses by itself. Such a corona may be formed by X-ray heating of the accretion disk or by dissipation of acoustical or magnetic energies there. Another possible origin may be companion evaporation by intercepted hard X- or γ-rays. Model calculations (White and Holt 1982) suggest that the path to the companion star would be optically clear to scattering once the primary luminosity is down to several times 10^{37} ergs s^{-1}, and that only at higher luminosities should one expect an extended X-ray source.

The second class of observations has been interpreted as providing an indication of the temperature of the corona. Various attempts to fit the X-ray spectra seem to improve once one introduces Comptonization in a relatively hot electron corona.

The third class of observations is that of the hard and the soft lags in the temporal and power spectra of QPO sources (van der Klis 1988). If these are interpreted as due to scatterings in a corona whose optical depth is indicated by the fitted energy spectra (see, e.g., Stollman *et al.* 1987), coronal dimensions can be inferred which are closer to the magnetospheric radius than to the binary separation, mostly because hard time lags observed in the horizontal-branch state of GX 5-1 and Cyg X-2 are of the order of several milliseconds (van der Klis 1988). However, Cyg X-2 shows hard lags of order ~70 ms in the normal branch state (Mitsuda 1986).

We note finally that some X-ray systems show direct evidence of hard X-rays reaching the companion (such as the Her X-1/HZ Her system; Bahcall and Bahcall 1972; Forman, Jones, and Liller 1972; Parmar *et al.* 1984) or reaching outside the system (as in Cyg X-3 and GX 5-1; Levine *et al.* 1984). All this by itself would seem to imply that Compton-thick coronae with temperatures of order a few kilo electron volts are not standard to all X-ray binaries, and at any rate it is not clear yet how to put all this evidence for coronae into a coherent picture.

We turn now to reasons why a beam of MeV γ-rays produced near the primary may be expected to be able to reach the secondary companion even though Eddington limit LMXBs may have coronae which are optically thick to X-rays. The nonthermal emission mechanisms for MeV γ-rays usually involve ultrarelativistic particles constrained to move along magnetic field lines. This generally results in strongly nonisotopic radiation. Possible support for a beam structure for the 10^{12} eV γ-rays from Her X-1 may be the long interval between observations of them interpreted as a time variation in beam sweepout directions.

a) Hole Boring

A pair of intense beamed cones of γ-rays which rotate rapidly through a corona can push most of the coronal matter out of the hollowed angular wedge they sweep through (cf. Fig. 7). For simplicity we assume that each cone has an opening half-angle θ_0 about

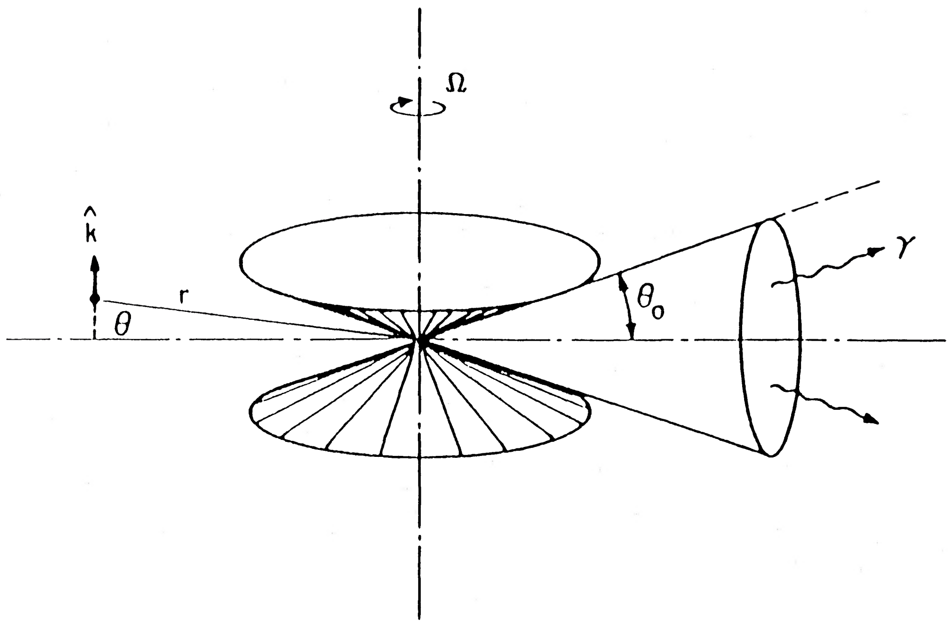

FIG. 7.—Schematic representation of the rotating γ-ray cone discussed in § VII and the hollowed wedge it sweeps out

a cone axis which sweeps the equatorial plane. Then the averaged γ-ray flux intensity in the beam cone at distance r from its apex is

$$i = \frac{L_\gamma}{4\pi r^2 \sin\theta_0} . \tag{7.1}$$

When $L_\gamma > (\sin\theta_0)L_{\rm Edd}$, the coronal matter within the region swept by unattenuated γ-ray beams would be pushed by a radiation force which exceeds the neutron star gravitational pull. This would be the case even if the special symmetry of Figure 7 is not appropriate. Scattered electrons (accompanied by protons and nuclei) would be accelerated out of the hollowed wedge, through its upper and lower surfaces, until the sum of the neutron star gravity (and centrifugal force) which pulls it back in,

$$F_{\rm grav} \simeq m_{\rm H} \frac{GM}{r^2} (\sin\theta) , \tag{7.2}$$

is less than the push of scattered radiation on an electron in the opposite direction

$$F_R = \frac{L_\gamma \sigma_{\rm T}}{4\pi r c \lambda_\gamma} \exp\left(-\frac{r}{\lambda_\gamma}\right) \frac{\sin\theta}{\sin\theta_0} , \tag{7.3}$$

where λ_γ is the (assumed constant) mean free path for large energy loss of MeV γ-rays by scattering on residual matter in the γ-ray emission cone. That lost energy is reemitted as synchrotron photons if a large enough magnetic field is present, or by inverse Compton boosting of accretion-powered X-rays which traverse the hollowed wedge. For $L_X \simeq 10^{38}$ ergs s^{-1} and $E_X \simeq 1$ keV, λ_e, the mean free path for electron–X-ray scattering, satisfies

$$\frac{\lambda_e}{r} \simeq \frac{4\pi r^2 E_X c}{L_X \sigma_{\rm T} r} \simeq 10^{-4} r_7 , \tag{7.4}$$

with $r_7 = r/(10^7$ cm). Thus for $\theta_0 > 10^{-2}$ there will be a conversion of lost γ-ray energy into much lower energy photons which push with a Thomson cross section on electrons in the cone. This is the origin of $\sigma_{\rm T}$ in equation (7.3). Where $|F_R|$ locally exceeds $|F_{\rm grav}|$, matter will be evacuated from the wedge until λ_γ increases enough for an equality to be reached. Then for a steady state

$$\frac{L_\gamma}{\sin\theta_0}\left(\frac{r}{\lambda_\gamma}\right) \exp\left(-\frac{r}{\lambda_\gamma}\right) = L_{\rm Edd} . \tag{7.5}$$

A solution exists as long as

$$L_\gamma > 2.718(\sin\theta_0)L_{\rm Edd} . \tag{7.6}$$

The stable one has $\lambda_\gamma > r$, so that the residual electron density in the wedge gives small total attenuation of the γ-ray beam throughout the corona's radial extent. The wedge will, however, remain open between the repeated beam passages only if the sound speed v_s of the coronal matter above and below it is not large enough to permit the transfer of matter deeply into the wedge during that interval:

$$\theta_0 r > \frac{2\pi}{\Omega} v_s . \tag{7.7}$$

For a canonical corona with temperature ~ 5 keV, $v_s \simeq 10^8$ cm s^{-1}. For $\theta_0 r > 10^6$ cm, equation (7) is satisfied as long as the period P for the return of a beam through γ-ray–irradiated coronal regions satisfies

$$P < 10^{-2} \text{ s} . \tag{7.8}$$

Note that once a hole has been bored, X-rays may leak through it as well. There are other attributes of coronae which may make γ-ray beam hole boring easier than indicated in equations (7.5)–(7.7). If a corona has an optical depth for X-ray scattering of order 5–10, then it is already almost transparent to γ-rays which have very considerably less than a Thomson cross section for scattering. If the γ-ray energy is much greater than 1 MeV, it will scatter mainly forward, and may still penetrate, although with degraded energy. In any event, if energy is deposited in it, the heated beam channel will expand and thus reduce the optical depth even if the radiation which results is insufficient to give enough pressure all by itself to push out beam channel coronal matter. On the other hand, if the corona is mainly an $e^{\pm}$ cloud or wind, gravity plays no important role in containing the corona, and the push of equation (7.3) is opposed only by the relatively small pressure of surrounding $e^{\pm}$ plasma. For a canonical coronal plasma at $T \simeq 5$ keV extending to a distance r away from the γ-ray source, this pressure can be overwhelmed by that from L_γ, $L_\gamma/4\pi rc\lambda_\gamma$, unless $L_\gamma < 10^{34} r_7$ ergs s^{-1}.

Finally, the $e^{\pm}$ component may diffuse through the corona and still reach the secondary, where it can deposit its rest energy into annihilation γ-rays. The expected $e^{\pm}$ wind has a total flux

$$\dot{N}_{\pm} \simeq \frac{L_X R}{2 m_e c^2 R_A} . \tag{7.9}$$

Each e^+ reaching the secondary's atmosphere will give a pair of γ-rays (each of energy $\sim m_e c^2 \simeq 511$ keV) within a surface thickness less than $\hat{\sigma}_T$. The range of the annihilation γ-rays is $\hat{\sigma}_T$, so that a power

$$\hat{L}_{\pm} = 2 m_e c^2 \dot{N}_{\pm} = \hat{L}_X \frac{R}{R_A} \simeq \hat{L}_\gamma \tag{7.10}$$

would be initially given to the local electrons to a depth $\hat{\sigma}_T$.

b) *Subsidence of Coronae*

Insofar as the primary's canonical corona may ultimately be sustained by a γ-ray–powered wind and accretion from the secondary, a drop in the incident γ-rays because of any residual coronal absorption would cause a subsidence of the corona. This could then result in a sub-Eddington steady state such that enough γ-rays reach the secondary to sustain the accretion. There is also a time lag between Roche lobe overflow and the emission of γ-rays when that overflow reaches the neutron star magnetosphere, which may be of order 10^3–10^6 s. Then, instead of a steady state, we may, on these time scales, have varying X-ray luminosity together with anticorrelated coronal surges and strong γ-ray fluxes reaching the secondary.

c) *Observations of Hard Spectra from LMXBs*

We note that some LMXBs, bright persistent sources and bursters, do indeed have power-law spectra not degraded by the passage through a 5 keV corona (White and Mason 1985; White, Stella, and Parmar 1988). For example, one of the two presently known compact binaries with degenerate companions, XB 1916−053, shows an energy spectral index around 0.5–0.7 for X-ray energies between 1 and 20 keV (White and Swank 1982; Smale *et al.* 1988). For several sources there is substantial evidence for high-energy emission at energies larger than 20 keV (Levine *et al.* 1984). It is argued elsewhere (Kluźniak *et al.* 1988*b*, *c*) that both the MeV γ-rays, sometimes observed from the direction of the Galactic center region, and the hard power-law X-ray spectrum around and above 10^2 keV have their origin in several near-Eddington LMXBs.

VIII. SUMMARY

Radiation from weakly magnetized neutron stars in LMXBs can be so efficient in driving substantial winds from secondaries with mass smaller than 10^{-1} $M_\odot$ that the secondaries may continue to feed the accretion disk with a mass loss near the neutron star's Eddington limit. This may occur even when other mechanisms for removing angular momentum from the binaries, such as gravitational radiation, are too weak to sustain large mass transfers. These binaries will then pass through this mass range in 10^7 yr or less and be relatively rare.

This bootstrapped accretion process, in which accretion powers X-rays and an $e^{\pm}/\gamma$-ray wind which effectively sustain the accretion, may turn off suddenly when the evaporating secondary mass falls below a few hundredths of a solar mass. When this occurs, accretion and the X-ray luminosity which results from it drop by several orders of magnitude and may be quenched completely. Although the arguments for MeV γ-ray/$e^{\pm}$ winds powered by accretion onto low magnetic field neutron stars in LMXBs are not yet compelling, that such winds should exist is a rather suggestive inference from reported TeV γ-ray observations of X-ray pulsars in binaries. The LMXB version of these γ-rays could play a crucial role in exciting the companion wind, either as the sole radiative component or as the component that transforms X-ray–heated coronae into a wind.

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