LOW-MASS NORMAL-MATTER ATMOSPHERES OF STRANGE STARS AND THEIR RADIATION

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ABSTRACT

The quark surface of a strange star has a very low emissivity for X-ray photons. I find that a small amount of normal matter at the quark surface with temperature in the range $10^7 \text{ K} \leq T_s \ll mc^2/k \approx 6 \times 10^{10} \text{ K}$ is enough to produce X-rays with high luminosity $L_n \approx (10^{32} - 10^{35})[\Delta M/(10^{-20} M_\odot)]^2 \text{ ergs s}^{-1}$. For a total atmospheric mass $\Delta M \sim 10^{-9} \text{ M}_\odot$, this luminosity may be as high as the Eddington limit. The mean energy of X-ray photons that are radiated from such a low-mass atmosphere of a strange star is $\sim 10^2[T_s/(10^8 \text{ K})]^{3/5} \approx 30 - 300$ times greater than the mean energy of X-ray photons that are radiated from the surface of both a neutron star and a strange star with a massive normal-matter envelope, $\Delta M \sim 10^{-5} \text{ M}_\odot$, for a fixed temperature at the stellar core. This raises the possibility that some black hole candidates with hard X-ray spectra are, in fact, strange stars with a low-mass atmosphere. The X-ray emission from single strange stars is estimated.

Subject headings: radiation mechanisms: thermal — stars: atmospheres — X-rays: stars

1. INTRODUCTION

Strange stars have been proposed by Witten (1984) as a new class of astronomical compact objects. If Witten’s idea is true, at least some part of the compact objects known to astronomers as pulsars, powerful accreting X-ray sources, X-ray and $\gamma$-ray bursters, etc., might be strange stars, not neutron stars, as is usually assumed (Alcock, Farhi, & Olinto 1986; Glendenning 1990; Caldwell & Friedman 1991; Madsen & Olesen 1991; Weber et al. 1996). Strange quark matter with a density of $\sim 4 \times 10^{14} \text{ g cm}^{-3}$ can exist, by hypothesis, up to the surface of strange stars. The quark surface is a very poor radiator at energies $\epsilon < 20 \text{ MeV}$ (Alcock et al. 1986), but the presence of “normal” matter (ions and electrons) at the quark surface may restore the ability of the surface to radiate soft photons (this is like painting with black paint on a silver surface). There is an upper limit to the amount of normal matter at the quark surface, $\Delta M \approx 10^{-5} \text{ M}_\odot$ (see, e.g., Glendenning & Weber 1992), set by the requirement that the density of the inner layer of normal matter at the quark surface cannot be greater than $\rho_s \approx 4.3 \times 10^{12} \text{ g cm}^{-3}$. Such a massive envelope of normal matter with $\Delta M \sim 10^{-5} \text{ M}_\odot$ completely obscures the quark surface.

If a neutron star is not too young (age $t \gtrsim 10^3$ yr), the stellar interior may be divided into two regions: the isothermal core with density $\rho > \rho_s \approx 10^{10} \text{ g cm}^{-3}$ and the outer envelope with $\rho < \rho_s$ (see, e.g., Guzman, Pethick, & Epstein 1983; Nomoto & Tsuruta 1987; Schaff 1990). At the core temperature $T_r \sim 10^8 - 10^9 \text{ K}$, the temperature decreases by a factor of $\sim 30 - 300$ in the envelope, from $T_r$ at the inner boundary of the envelope to $\sim 10^7[T_r/(10^8 \text{ K})]^{3/5} \text{ K}$ at the neutron star surface. Since $\rho_s \sim \rho_s$ for strange stars with massive envelopes, $\Delta M \sim 10^{-5} \text{ M}_\odot$, the temperature variation between the quark surface and the surface of the normal-matter envelope is more or less the same as the core-to-surface temperature variation of neutron stars for a fixed temperature at the stellar center. Besides, if the strange quark matter is superfluid, the cooling behavior of the quark cores of strange stars is more or less similar to the cooling behavior of the isothermal cores of neutron stars (see, e.g., Weber et al. 1996). Therefore, at an age $t \gtrsim 10^7$ yr, photon radiation from the surfaces of compact objects is not a good observational signature of strange stars with massive envelopes of normal matter.

It has been pointed out (e.g., Haensel, Paczyński, & Amsterdam 1991; Burrows & Hayes 1995; Hartmann & Woosley 1995; Cheng & Dai 1996) that the temperature in the interiors of both neutron stars and strange stars at the moment of their formation is very high, $T_s$ on the order of a few times $10^{11} \text{ K}$. The mass of gas that is ejected from the surface of such a hot neutron star during $\sim 10$ s since its formation is about $10^{-3} - 10^{-2} \text{ M}_\odot$ (see, e.g., Woosley & Baron 1992; Levinson & Eichler 1993; Woosley 1993). This value is considerably larger than the upper limit on the mass of normal-matter envelopes of strange stars. The input physics for calculations of gas outflow from the envelopes is similar for both hot neutron stars and hot strange stars. Therefore, it is natural to expect that if any normal matter remains at the surface of a strange star at $t \gg 10$ s, its mass is many orders smaller than the maximum. In the process of gas accretion onto such a strange star, the stellar atmosphere has to pass through the stage with $\Delta M \ll 10^{-5} \text{ M}_\odot$ before reaching the maximum, $\Delta M \sim 10^{-5} \text{ M}_\odot$, that has usually been assumed in studies of the thermal structure and photon radiation of strange stars. Below, photon radiation of a nonmagnetic strange star with a low-mass atmosphere, $\Delta M \ll 10^{-5} \text{ M}_\odot$, is considered.

2. STRUCTURE AND PHOTON RADIATION OF THE ATMOSPHERE

We consider here a star consisting of a core of strange quark matter surrounded by a low-density atmosphere of normal matter with a mass $\Delta M$ that is many orders of magnitude smaller than $10^{-5} \text{ M}_\odot$. The core acts on the atmosphere as a heat reservoir. The thermal structure and photon radiation of the atmosphere can be found by solving the heat transfer problem with $T = T_3$ as a boundary condition at the inner layer of normal matter. We assume that the temperature of the quark surface is $T_3 \approx 10^7 \text{ K}$ and that the hot gas of the atmosphere emits mainly as a result of free-free transitions (Gaz & Salpeter 1983).

A young strange star cools very rapidly as a result of intense...
neutrino emission (see, e.g., Weber et al. 1996). Therefore we restrict our consideration to the case of nonrelativistic temperatures, \( T_s \ll T_0 = mc^2/k \approx 6 \times 10^6 \) K. At such temperatures, the energy loss per unit gas volume by bremsstrahlung radiation is

\[
Q_{\text{br}} = \frac{C_1N_eZ^2T^{1/2}}{1 \text{ erg s}^{-1} \text{ cm}^{-3}}, \tag{1}
\]

where \( N_e = \rho/m_p A \) is the ion density in cm\(^{-3}\), \( N_e = NZ \) is the electron density in cm\(^{-3}\), \( m_p \) is the proton mass, \( A \) is the mass number of ions, \( Z \) is their electrical charge, \( T \) is the temperature in kelvins, \( C_1 \equiv 1.4 \times 10^{-2} g(T) \), and \( g(T) \) is the frequency-averaged Gaunt factor, which is in the range 1.1–1.5. Choosing a value of 1.2 for \( g(T) \) will yield an accuracy to within about 20% (Karzas & Latter 1961).

At \( T \ll T_0 \), the scale height, \( \Delta x \), of the atmosphere is very small compared with the stellar radius \( R \), and a plane-parallel approximation may be used. In this approximation all parameters depend upon only one coordinate, \( x \), which is the distance from the quark surface.

The set of equations that describes the structure of the atmosphere will be the equation of hydrostatic equilibrium and the energy transport equation:

\[
\frac{dP}{dx} = -\frac{\rho GM}{R^2}, \quad \frac{dT}{dx} = -\frac{Q_{\text{br}}}{\rho}. \tag{2}
\]

where \( P = (N_e + N)kT = \rho kT/m_p \mu \) is the gas pressure, \( \mu = A/(1 + Z) \) is the mean molecular weight, \( G \) is the gravitational constant, \( M \) is the stellar mass, and \( F \) is the heat flux due to both thermal conductivity and convection (Schwarzschild 1958). The absorption of radiation is ignored in the energy transport equation because in our case the atmosphere is optically thin for the bulk of the radiation (see below).

### 2.1. Isothermal Atmosphere

The characteristic time of heat conduction in the atmosphere is

\[
\tau_{\text{heat}} \simeq \frac{k(1 + Z)N_e(\Delta x)^2}{kN_e} \text{ s,} \tag{3}
\]

and the characteristic cooling time of the hot atmospheric plasma via bremsstrahlung is

\[
\tau_{\text{cool}} \simeq 1.5 \times 10^{11} \frac{(1 + Z)T^{1/2}}{Z^2N_e} \text{ s,} \tag{4}
\]

where \( \eta \approx 10^{-7} Z^{-1/2} \text{ ergs s}^{-1} \text{ K}^{-1} \) is the coefficient of heat conductivity for a rarefied, totally ionized plasma (Spitzer 1967).

When \( \tau_{\text{heat}} \ll \tau_{\text{cool}} \), the atmosphere is nearly isothermal, \( T \approx T_s \), and the equation of hydrostatic equilibrium (eq. [2]) can be integrated immediately:

\[
\rho = \rho_0 \exp (-x/\Delta x), \tag{5a}
\]

where

\[
\Delta x = \frac{R^2kT_s}{GM\mu m_p}. \tag{5b}
\]

In this case, the photon luminosity of the atmosphere is

\[
L = 4\pi R^2 \int_0^\infty Q_{\text{br}} \, dx \approx \frac{4 \times 10^{10} Z}{A(1 + Z)} \left( \frac{R}{10^6 \text{ cm}} \right)^4 \left( \frac{M}{1 M_\odot} \right) \times \left( \frac{T_s}{10^4 \text{ K}} \right)^{1/2} \left( \frac{\Delta M}{10^{12} \text{ g}} \right)^2 \text{ ergs s}^{-1}, \tag{6}
\]

where \( \Delta M = 4\pi R^2 \rho \Delta x \) is the total mass of the atmosphere.

Using equations (3) and (4), the condition \( \tau_{\text{heat}} \ll \tau_{\text{cool}} \) may be written as a limitation of the atmospheric mass: \( \Delta M \ll \Delta M_\odot \), where

\[
\Delta M_\odot \approx 7 \times 10^{11} \frac{A}{Z} \left( \frac{T_s}{10^4 \text{ K}} \right)^{3/2} \left( \frac{R}{10^6 \text{ cm}} \right)^2 \text{ g}. \tag{7}
\]

### 2.2. Convective Atmosphere

The heat flux due to thermal conductivity, which is responsible for the heat transport in the atmosphere when \( \Delta M < \Delta M_\odot \), is \( F = -\eta dT/dx \). Using this, we can rewrite the equation of hydrostatic equilibrium (eq. [2]) in the form

\[
\frac{d\rho}{dx} = -\frac{\rho}{T} \left( \frac{GM\mu \rho}{R^2} - \frac{F}{\eta} \right). \tag{8}
\]

In steady state, which we assume in this Letter, we have the following boundary condition for \( F \) at \( x = 0 \): \( F_{\text{tot}} = L_1/4\pi R^2 \). From this condition and equation (8), the gradient of the density at the quark surface, \( x = 0 \), is

\[
\frac{d\rho}{dx} \bigg|_{x=0} = -\frac{\rho_0}{T_s R^2} \left( \frac{GM\mu \rho}{k} - \frac{L_1}{4\pi \eta} \right). \tag{9}
\]

If the heat transport in the atmosphere is only due to thermal conductivity and the photon luminosity of the atmosphere is higher than

\[
L_1 = \frac{4\pi GM\mu \rho}{k} \eta \bigg|_{x=r-\tilde{r}} \approx 3 \times 10^{33} \frac{\mu}{Z} \left( \frac{M}{1 M_\odot} \right) \left( \frac{T_s}{10^4 \text{ K}} \right)^{5/2} \text{ ergs s}^{-1}, \tag{10}
\]

then the density of normal matter near the quark surface should increase with the distance from the surface, \( \left. (d\rho/dx) \right|_{x=0} > 0 \). However, such a density behavior is unstable, and convection develops in the atmosphere at \( L > L_1 \). The value of \( L_1 \) coincides with the photon luminosity (eq. [6]) after substituting \( \Delta M = \Delta M_\odot \) from equation (7) into equation (6).

If convection takes place in the stellar atmosphere, it is a good approximation, as a rule, to say that the temperature gradient is equal to the adiabatic one, i.e., \( PT^{\gamma/(\gamma - 1)} = \text{const} \), where \( \gamma \) is the ratio of the specific heats at constant pressure and at constant volume. In this case, from equation (2), the temperature and density distributions are

\[
T = T_s \left[ 1 - \frac{(\gamma - 1)x}{\gamma T_s} \right], \quad \rho = \rho_0 \left( \frac{T_s}{T} \right)^{1/(\gamma - 1)}, \tag{11}
\]

where \( \Delta x \) is determined by equation (5b).

Using equations (1) and (11), one can obtain the photon luminosity of the convective atmosphere, \( L = 4\gamma L_1/(3\gamma + 1) \).
where $L$ is the photon luminosity of the isothermal atmosphere, which is given by equation (6). Since $\gamma = 1$, the value of $L$ is in the range $L \leq L < \frac{3}{2}L$. For a rarefied, totally ionized plasma, we have $\gamma = \frac{3}{2}$ and $L = \frac{3}{2}L$. As noted earlier, the accuracy of equation (1) for the energy loss $Q_{\text{er}}$ is about 20%. Moreover, we did not take into account the general relativistic effects, which are $\sim 20\%$. Therefore the accuracy of our calculations of bremsstrahlung radiation from the atmospheres of strange stars is several times 10%. We can see that the difference between $L$ and $L$ is within the accuracy of our calculations.

The adiabatic approximation may be used to estimate the photon luminosity of the convective atmosphere only if the characteristic time of convection $t_{\text{conv}} = \Delta t/t_{\text{conv}}$ which is the main process of heat transport at $\Delta M > \Delta M_0$, is smaller than the characteristic cooling time $t_{\text{cool}}$ for the atmospheric plasma, where $t_{\text{conv}}$ is the convective velocity, which is limited by the velocity of sound, $v_{\text{conv}} \approx c_s = (\gamma p/\rho)^{1/2}$. Using this and equation (4), the condition $t_{\text{conv}} < t_{\text{cool}}$ may be written in the form $\Delta M < \Delta M_0$, where

$$\Delta M_0 \approx 4 \times 10^{12} A \left( \frac{T_3}{10^8 \text{ K}} \right) \left( \frac{R}{10^8 \text{ cm}} \right)^2 \text{ g.}$$  \hspace{1cm} (12)

At $\Delta M = \Delta M_0$, the photon luminosity is

$$L_2 \approx 6 \times 10^{12} \left( \frac{1 + Z}{Z} \right)^3 \left( \frac{M}{1 M_\odot} \right) \left( \frac{T_3}{10^8 \text{ K}} \right)^{1/2} \text{ ergs s}^{-1}. \hspace{1cm} (13)$$

For $M = 1.4 M_\odot$, $T_3 \approx 2 \times 10^8 \text{ K}$, and $Z = 1$, the value of $L_2$ is $\sim 10^{56}$ ergs s$^{-1}$, which is only 2 orders of magnitude smaller than the Eddington limit, $L_{\text{Edd}} \approx 1.3 \times 10^{53} (A/Z) [M/(1 M_\odot)] \text{ ergs s}^{-1}$.

The mean free-free optical depth of the atmosphere is

$$\tau_0 \approx \alpha_0 \Delta x \sim 10^{-9} \left( \frac{L}{10^{34} \text{ ergs s}^{-1}} \right) \left( \frac{T_3}{10^8 \text{ K}} \right)^{-4},$$  \hspace{1cm} (14)

where $\alpha_0 \approx 10^2 T_3 Q_{\text{er}}$ is the Rosseland mean of the free-free absorption coefficient. At $T_3$ greater than a few times $10^8 \text{ K}$, the atmosphere is optically thin, $\tau_0 \ll 1$, up to $L \sim L_{\text{Edd}}$.

For $\Delta M > \Delta M_0$, both thermal conductivity and convection are not able to account for the cooling of atmospheric matter, and a thermal instability develops in the atmosphere. As a result, the atmosphere cannot be in hydrostatic equilibrium during a time longer than $\tau_{\text{cool}}$, and it is strongly variable on a timescale of a few times $(2R^2 \Delta x/GM)^{1/2} \sim 10^{-5}[T_3/(10^8 \text{ K})]^{1/2} \text{ s}$. Consideration of this variability and estimation of the photon luminosity at $\Delta M > \Delta M_0$ are under way and will be published elsewhere. Here it is worth noting only that at $T_3$ greater than a few times $10^8 \text{ K}$ and $\Delta M > \Delta M_0$, the tendency of the photon luminosity to increase with increase of $\Delta M$ has to be held up to $L = L_{\text{Edd}}$.

3. CONCLUSIONS AND DISCUSSION

The photon luminosity of a strange star with a low-mass atmosphere, $\Delta M < \Delta M_0$, is given by equation (6) with an accuracy of several times 10% irrespective of the atmospheric structure. The photon luminosity may be very high, up to $\sim L_{\text{Edd}}$. It is very important for discovery of strange stars that the mean energy of X-ray photons that are radiated from such stars is substantially higher than the mean energy of X-ray photons that are radiated from the surface of both a neutron star and a strange star with a massive envelope of normal matter, $\Delta M \sim 10^{-5} M_\odot$ for a fixed temperature at the stellar core.

A source of X-rays may be a strange star if it meets the following criteria:

1. The X-ray emission (or at least one of its components) may be fitted by thermal emission of optically thin plasma at $kT$ up to $\sim 10^5$ keV.
2. The X-ray flux is variable at the X-ray luminosity $L_x > L_2 \approx (0.05–5) \times 10^{39} [T/(10^8 \text{ K})]^{1/2} \text{ ergs s}^{-1}$, depending upon the composition of the atmosphere. This variability is either irregular or quasi-periodic.
3. The mass of the compact X-ray source is on the high side for neutron star masses, because conversion of a neutron star to a strange star requires a very high density at the center of the neutron star (see, e.g., Alcock et al. 1986).

A few enigmatic X-ray sources that are considered as black hole candidates (see, e.g., Cherepashchuk 1996) answer these criteria and may be, in fact, strange stars with a low-mass atmosphere. They are 1E 1740.7–2942, GRS 1915–105, GRO J0422+32, GX 339–4, and S4 433. Some other powerful X-ray sources that are black hole candidates as well, for example, Cyg X-1, have both hard X-ray spectra that may be fitted by emission of optically thin plasma and strong X-ray flux variability. The existing lower limits to the masses of these X-ray sources are substantially higher than the Oppenheimer-Volkoff limit for a strange star (Cherepashchuk 1996). However, these objects may be, in principle, triple systems with a strange star (cf. Bahcall et al. 1974).

The thermal energy of a strange star itself is not enough for the star to be a powerful X-ray source, $L_x \sim (0.1–1) L_{\text{Edd}}$, for a long time, $\tau \approx 10^{10}$ yr, and accretion of gas onto the strange star is necessary to account for such a strong, prolonged X-ray emission. The kinetic energy of the accreted gas may be transformed into emission of the strange star atmosphere in the following way: Let the magnetic field at the stellar surface be $\sim (0.3–1) \times 10^{11} \text{ G}$. This field canalizes the gas motion along the field lines to the quark surface. In this case, the kinetic energy of ions at the surface is about 100 MeV nucleon$^{-1}$, which is $\sim 5$ times greater than the Coulomb barrier at the quark surface (Alcock et al. 1986). Accreted particles penetrate through the quark surface (cf. below), and they are dissolved into quark matter. As a result, the quark core is heated at the magnetic poles. The process of heat transport through the core is very fast because of the very high heat conductivity of quark matter, and therefore the quark core is nearly isothermal. Then the energy that is released in the process of gas accretion is radiated from the normal-matter atmosphere more or less isotropically, just as discussed above.

The total atmospheric mass that restores the ability of the quark surface to radiate X-ray photons is extremely small. At first sight, even if the rate of gas accretion onto a hot strange star is very small, the atmospheric mass increases rapidly up to the value at which the photon luminosity is $\sim L_{\text{Edd}}$. But this does not necessarily happen. The point is that ions of accreted gas, in the process of their motion through the atmosphere, collide with the atmospheric ions and draw them into the quark surface. As a result, a steady state of the atmosphere may be achieved before the photon luminosity is $\sim L_{\text{Edd}}$. Let us estimate the photon luminosity at such a steady state. For definiteness, we assume that a single strange star with $M \approx 1.4 M_\odot$ is at rest in a uniform ionized gas of pure hydrogen. We shall assume that the gas is at rest at infinity with a density of


\( \sim 1 \text{ cm}^{-3} \) and temperature of \( \sim 10^7 \) K. This situation will correspond to accretion in a typical interstellar H II region. In this case, the accretion rate is \( M \approx 2 \times 10^{10} \text{ g s}^{-1} \) (see, e.g., Shapiro & Teukolsky 1983). The efficiency for radiation of the gas during its accretion is very low, i.e., the gas motion is nearly adiabatic. The mean kinetic energy of accreted protons at the surface is \( E_{\text{kin}} \approx 100 \text{ MeV} \). The temperature of accreted gas near the stellar surface is \( T_s \approx 10^{11} \) K (\( kT_s \approx 10 \) MeV). The bulk of accreted protons passes through the quark surface. Protons can be reflected from the quark surface back into the atmosphere only if, in the frame of the star, the kinetic energy of their radial motion to the stellar surface is smaller than the Coulomb barrier. Taking this into account and assuming that the accreted protons are Maxwellian, the rate of the atmospheric mass accumulation is \( \dot{M}_a \approx \exp \left( -\frac{E_{\text{kin}}}{kT_s} \right) M \approx \exp \left( -10 \right) M \). The characteristic time of the atmospheric mass decrease due to bombardment by the accreted protons is \( \Delta t_a \approx \left( n_s \sigma v_s \right)^{-1} \approx 4\pi R^2 m_p / M \), where \( n_s \approx \dot{M} / 4\pi R^2 v_s m_p \) is the density of accreted protons at the surface, \( v_s \) is their velocity, and \( \sigma \approx 10^{-28} \text{ cm}^2 \) is the cross section for proton-proton collision at energies of \( \sim 100 \text{ MeV} \). The atmospheric mass at the steady state is \( \Delta M_a \approx \dot{M}_a \Delta t_a \approx 4 \exp \left( -\frac{E_{\text{kin}}}{kT_s} \right) R^2 m_p \sigma^{-1} \approx 10^{12} \text{ g} \), which does not depend upon the accretion rate. The steady state may be reached in \( \sim \Delta t_a \sim 10^4 \) s. From equation (6), for \( T_s \sim 10^7-10^8 \) K and \( \Delta M_a \approx 10^{12} \text{ g} \) the expected X-ray luminosity of a single strange star is \( L_x \sim 10^{34} \text{ ergs s}^{-1} \), that is, of the order of X-ray luminosity of a single neutron star in \( \sim 10^5-10^6 \) yr after its formation. But, in the case of such a strange star, the expected X-ray spectrum is much harder than the X-ray spectrum of the thermal radiation from a single neutron star with the same luminosity.

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