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On the X-ray spectra of soft gamma repeaters

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ABSTRACT

Radiation transfer in a scattering medium in a superstrong magnetic field is considered. Because cross-sections depend on frequency, photons with different energies escape layers with different temperatures and therefore the spectrum of the outgoing radiation differs significantly from the equilibrium blackbody or Bose–Einstein spectrum. It is shown that the emergent spectrum (the photon flux per unit energy band) is flat at low energies. Applications of the result to soft gamma repeaters (SGRs) are discussed. Even though the spectrum is strongly distorted when the radiation propagates through the magnetosphere, a flat segment may be observed in the outgoing spectrum if the surface magnetic field of the neutron star is not too high, $B < 10^{15}$ G.

Key words: magnetic fields – radiative transfer – scattering – stars: neutron.

1 INTRODUCTION

Soft gamma repeaters are now widely accepted to be magnetars, i.e. neutron stars with superstrong magnetic fields, $B \sim 10^{15}$ G (for a recent review see Thompson 2002). Bursts are produced by sudden release of the magnetic energy resulting in formation of a hot electron–positron plasma. With the exception of a possible initial short spike, this plasma is trapped in the stellar magnetosphere forming a bubble, which cools by radiation for 10–100 s (Thompson & Duncan 1995, hereafter TD). The radiation flux from this bubble exceeds the Eddington flux orders of magnitude because a very strong magnetic field suppresses the radiation cross-sections (Paczyński 1992; Ulmer 1994; Miller 1995). The optical depth of the bubble is very large therefore the radiation is thermalized inside the bubble or at least a Bose–Einstein spectrum is formed because Comptonization dominates the radiation processes (TD). Because the scattering cross-sections in a superstrong magnetic field depend on frequency (see, e.g., Mészáros 1992), the spectrum of the emerged radiation should differ significantly from the Planck or the Bose–Einstein spectrum (Lyubarskii 1987, hereafter L; Ulmer 1994). Namely because the cross section decreases with the decreasing photon energy, low-energy photons scatter less effectively than those around the ‘thermal peak’ and come from deep in the atmosphere where the temperature is higher. Therefore the radiation flux at low energies exceeds that of the blackbody with the bolometric temperature.

The aim of this research is to study radiation transfer at the parameters appropriate for bubbles formed during bursts in SGRs. Radiation from the surface of magnetars was considered recently by Bulik & Miller (1997), Ho & Lai (2001), Özel (2001), Zane et al. (2001), Lai & Ho (2002). In these conditions, absorption dominates scattering whereas in the bubbles the opposite is true and therefore a special consideration is necessary. Comptonization in a strong magnetic field was considered before in the X-ray pulsar regime (Nagel 1981; Mészáros & Nagel 1985; Lyubarskii 1988a,b); in case the gyrofrequency exceeds the radiation frequencies the scattering physics is essentially the same. In these papers radiation transfer at the constant temperature was considered, which is justified if the Thomson optical depth is not too large. The optical depth of the bubbles formed in SGRs is extremely large (TD); the temperature gradient becomes a crucial factor in this case (L). It will be shown that the photon spectrum (the photon flux per unit energy) of the radiation escaping the bubble photosphere is flat at the energies $\epsilon < T_b$, where T_b is the bolometric temperature in energy units. The high-energy part of this spectrum is distorted by photon splitting when the radiation propagates through the magnetosphere. The low-energy part may be distorted by ion cyclotron absorption. Nevertheless a flat step may be observed in the outgoing spectrum at $E \leqslant 15$–20 keV if the magnetic field at the surface of the star is not too high, $B \leqslant 10^{15}$ G.

The article is organized as follows. In Section 2 the radiation cross-sections and the radiation transfer equation are written out. In Section 3 qualitative estimates are presented that explain why a flat spectrum is formed if radiation cross-sections grow with the photon energy as ϵ^2. It is proved in Section 4 that in the case of interest one can consider radiation transfer in LTE approximation. The results of numerical solution to the radiation transfer equation are presented in Section 5; distortions of the spectrum at the ray path through the magnetosphere are discussed qualitatively in this section also. In the Appendix the vacuum resonance effects are discussed.

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2 BASIC EQUATIONS

A bubble formed during the SGR outburst may be considered as a hot ($T \geq 10\,\mathrm{keV}$), optically thick, rarefied medium embedded in a very strong magnetic field (TD). Within the bubble, the Compton scattering dominates the opacity and the excitation energy of the first Landau level greatly exceeds the photon energy,

$$\epsilon \ll mc^2\left(\sqrt{1+\frac{2B}{B_\mathrm{q}}}-1\right). \tag{1}$$

Here ϵ is the photon energy, B the magnetic field, $B_\mathrm{q} = m^2c^3/(e\hbar) = 4.4\times10^{13}\,\mathrm{G}$ the 'quantum' magnetic field, m the electron mass. In this case the radiation cross-sections depend strongly on polarization of the normal modes. The cross-section of a photon with electric vector perpendicular to $\boldsymbol{B}$ (the extraordinary mode, below E-mode) is strongly suppressed as compared to the cross-section of a photon in the orthogonal polarization state (the ordinary mode, below O-mode). When the dielectric tensor is dominated by vacuum polarization corrections, the scattering cross-sections may be written as (see, e.g., Mészáros 1992)

$$\mathrm{d}\sigma_{\mathrm{O}\to\mathrm{O}} = \frac{3}{4}\sigma_\mathrm{T}\sin^2\theta\sin^2\theta'\mathrm{d}(\cos\theta'), \tag{2}$$

$$\mathrm{d}\sigma_{\mathrm{O}\to\mathrm{E}} = \frac{3}{8}\sigma_\mathrm{T}\left(\frac{\epsilon B_\mathrm{q}}{mc^2B}\right)^2\cos^2\theta\,\mathrm{d}(\cos\theta'), \tag{3}$$

$$\mathrm{d}\sigma_{\mathrm{E}\to\mathrm{O}} = \frac{3}{8}\sigma_\mathrm{T}\left(\frac{\epsilon B_\mathrm{q}}{mc^2B}\right)^2\cos^2\theta'\mathrm{d}(\cos\theta'), \tag{4}$$

$$\mathrm{d}\sigma_{\mathrm{E}\to\mathrm{E}} = \frac{3}{8}\sigma_\mathrm{T}\left(\frac{\epsilon B_\mathrm{q}}{mc^2B}\right)^2\mathrm{d}(\cos\theta'), \tag{5}$$

where σ_T is the Thomson cross-section, θ and θ' are the angles between the direction of propagation of the photon and the magnetic field before and after the scattering, correspondingly.

One can see that at the condition (1) the scattering rates of photons in two polarization modes differ drastically. The same is correct for any absorption-emission process. The radiation transfer may be qualitatively described as follows (L; TD). Diffusion of photons occurs primarily in the E-mode whereas O-photons are locked. Redistribution of photons in frequency occurs predominantly in the O-mode because a photon undergoes much more scatterings in the O-mode than in the E-mode. Because the mode exchange rate is comparable to the scattering rate in the E-mode and because the average change in the photon energy in one scattering is small, $\Delta\epsilon/\epsilon \sim T/mc^2 \ll 1$, the energy of an E-photon changes insignificantly before it converts into an O-photon (or escape if it is close to the surface). Therefore one can neglect the energy redistribution as well as photon production processes in the E-mode and consider the mode exchange term as a source term in the radiation transfer equation for the E-mode photons.

Let us assume for the sake of simplicity that the E-mode photosphere is geometrically thin; then one can consider the radiation transfer in one dimensional approximation and one can also neglect the photon splitting as well as inhomogeneity of the magnetic field in the photosphere. Now the transfer equation for the E-mode photons may be simply written with the use of the cross-sections (3–5) as a balance equation (see also L)

$$\cos\theta\frac{\mathrm{d}n_\mathrm{E}(\theta)}{\mathrm{d}x} = \sigma_\mathrm{T}N\left(\frac{\epsilon B_\mathrm{q}}{mc^2B}\right)^2\left[-n_\mathrm{E}(\theta)+\frac{3}{8}\int n_\mathrm{E}(\theta')\,\mathrm{d}(\cos\theta') + \frac{3}{8}\int\cos^2\theta' n_\mathrm{O}(\theta')\,\mathrm{d}(\cos\theta')\right]. \tag{6}$$

Here N is the total density of electrons and positrons, n_E and n_O are the phase photon densities in the corresponding modes.

3 PRELIMINARY ESTIMATES

Prior to solving equations, let us consider the problem qualitatively. A photon undergoes $\kappa \sim \sigma_{\mathrm{O}\to\mathrm{O}}/\sigma_{\mathrm{O}\to\mathrm{E}} \sim (B/B_\mathrm{q})^2 \times (mc^2/T)^2$ scatterings in the O-mode before it is converted into the E-mode. The corresponding Comptonization parameter is large,

$$y = \kappa\frac{T}{mc^2} = \left(\frac{B}{B_\mathrm{q}}\right)^2\frac{mc^2}{T} \gg 1 \tag{7}$$

which ensures that the distribution function of the O-photons has a Bose–Einstein form. The chemical potential depends on efficiency of the photon production processes. It will be shown in the next section that at the conditions of interest the double-Compton scattering alone supplies enough photons to keep the chemical potential rather small. Therefore in the radiation transfer equation (6), one can substitute n_O by the Planck distribution, $n_B(\epsilon, T) = [\exp(\epsilon/T)-1]^{-1}$, thus reducing the problem to the LTE case.

The radiation transfer may be conveniently described in terms of the Rosseland mean free path (e.g., Mihalas 1978) of extraordinary photons

$$N\lambda_\mathrm{R} = \frac{\int(\sigma_{\mathrm{E}\to\mathrm{O}}+\sigma_{\mathrm{E}\to\mathrm{E}})^{-1}\epsilon^4\exp(\epsilon/T)[\exp(\epsilon/T)-1]^{-2}\,\mathrm{d}\epsilon}{\int\epsilon^4\exp(\epsilon/T)[\exp(\epsilon/T)-1]^{-2}\,\mathrm{d}\epsilon} = \frac{5}{4\pi^2\sigma_\mathrm{T}}\left(\frac{mc^2B}{TB_\mathrm{q}}\right)^2. \tag{8}$$

The temperature distribution within the bubble may be approximately found from the diffusion equation

$$\frac{c\lambda_\mathrm{R}}{3}\frac{dU_\mathrm{B}}{\mathrm{d}x} = \mathcal{F}, \tag{9}$$

where $\mathcal{F}$ is the radiation flux, $U_\mathrm{B} = (2\sigma/c)T^4$ the blackbody energy density in one polarization mode, $\sigma = \pi^2/(60\hbar^3c^2)$ the Stefan–Boltzmann constant. Let us introduce the bolometric temperature, T_b, such that the total radiation flux is presented as $\mathcal{F} = \frac{1}{2}\sigma T_\mathrm{b}^4$. It is appropriate to measure the column density in units of the Rosseland optical depth at the bolometric temperature. Then the solution to equations (8,9) may be written as

$$T = T_\mathrm{b}\sqrt{1+\frac{3}{4}\tau} \tag{10}$$

$$\tau = \frac{4\pi^2}{5}\sigma_\mathrm{T}\left(\frac{T_\mathrm{b}B_\mathrm{q}}{mc^2B}\right)^2\int N\,\mathrm{d}x. \tag{11}$$

The emergent spectrum is determined by the radiation spectrum at the depth corresponding to the mean free path of an E-photon,

$$\frac{5}{4\pi^2}\left(\frac{\epsilon}{T_\mathrm{b}}\right)^2\tau \approx 1. \tag{12}$$

In each layer, the radiation spectrum is close to Planckian however

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photons with different energies come from different layers with different temperatures. Therefore the emergent spectrum may be approximately presented as a blackbody spectrum with the temperature that depends on the photon energy according to equations (10,12). In terms of the photon flux, $F \equiv \epsilon^2 n$, the emergent spectrum may be expressed as

$$F = 0.47\epsilon^2 \left\{ \exp\left[\frac{\epsilon^2}{T_{\mathrm{b}}\sqrt{\epsilon^2 + (3\pi^2/5)T_{\mathrm{b}}^2}} \right] - 1 \right\}^{-1}. \tag{13}$$

The factor 0.47 is introduced to preserve the total energy flux. One can see that at low energies, $\epsilon < T_{\mathrm{b}}$, the emergent spectrum is flat instead of rising as truly thermal distributions do. The reason is that due to energy dependence of the cross-sections, the observer 'sees' at lower energies deeper layers with larger temperatures. The above considerations violates at low enough energy where the bubble becomes transparent for the E-photons. This energy may be estimated substituting into equation (12) the total optical depth of the bubble. One can easily see that for the typical bubble parameters (TD) this energy is too low, $\ll$1 keV, therefore below we will neglect the effects of the finite optical depth. The temperature distribution and the emergent spectrum will be calculated numerically in Section 5.

4 VALIDITY OF THE LTE APPROXIMATION

A balance between the photon production processes, Compton redistribution of the photons in energy and mode exchange scattering determines the radiation spectrum of the O-mode. Even at the E-mode photosphere (the layer with the optical depth for the E-photons about unity) the optical depth for the O-photons is extremely large, $\tau_{\mathrm{O}} \sim (mc^2B/\epsilon B_{\mathrm{q}})^2 \gg 1$, therefore their angular distribution is isotropic. Because the spatial diffusion from this depth takes $\sim\tau_{\mathrm{O}}^2$ scatterings whereas the photon is converted into the E-mode after only $\sim\sigma_{\mathrm{O}\to\mathrm{O}}/\sigma_{\mathrm{O}\to\mathrm{E}} \sim (mc^2B/\epsilon B_{\mathrm{q}})^2 \sim \tau_{\mathrm{O}}$ scatterings, one can neglect the spatial diffusion of the O-photons. The balance equation may be written as (L)

$$Q + \frac{2}{15}\sigma_{\mathrm{T}}\frac{T}{mc^2}\frac{1}{\epsilon^2}\frac{\partial}{\partial\epsilon}\epsilon^4\left[\frac{\partial n_{\mathrm{O}}}{\partial\epsilon} + \frac{n_{\mathrm{O}} + n_{\mathrm{O}}^2}{T}\right] + \frac{1}{4}\sigma_{\mathrm{T}}\left(\frac{\epsilon B_{\mathrm{q}}}{mc^2B}\right)^2\left[\frac{3}{2}\cos^2\theta\int n_{\mathrm{E}}(\theta')\,\mathrm{d}(\cos\theta') - n_{\mathrm{O}}\right] = 0. \tag{14}$$

Here Q is the photon production rate, the second term describes Comptonization in the Fokker–Planck approximation and the third one the mode exchange scattering. The factor 2/15 before the Comptonization operator appears because of angular dependence of the scattering cross-section (Basko & Sunyaev 1975; Lyubarskii 1988a).

Within the bubble, the main photon production mechanisms are photon splitting and double-Compton scattering (TD).[1] Let us show that the double-Compton scattering alone provides enough photons to keep the distribution function close to the Planck one. Photons are produced in the double-Compton scattering process predominantly at low energies, $\epsilon \ll T$, then they gain energy by Comptonization. The photon production rate in a non-magnetized plasma is written as (Lightman 1981)

$$Q = \frac{4\alpha}{3\pi}\frac{\sigma_{\mathrm{T}}}{m^2c^4}\frac{\exp(\epsilon/T) - 1}{\epsilon^3}[n_B(\epsilon, T) - n_{\mathrm{O}}(\epsilon)]I, \tag{15}$$

where, $\alpha = e^2/(\hbar c) = 1/137$,

$$I = \int \epsilon^4[1 + n_{\mathrm{O}}(\epsilon)]n_{\mathrm{O}}(\epsilon)\,\mathrm{d}\epsilon.$$

Because this process has not been considered yet at the condition (1), let us use equation (15) taking into account that at the condition (1) the cross-sections for the O-mode radiation differ from the non-magnetized cross-sections only by some angular factor and therefore the averaged in angles magnetized cross-sections are not much less than the non-magnetized ones.

Comparing different terms in equation (14), one can see that the Comptonization term dominates the balance equation at $\epsilon \sim T$ and therefore the spectrum in this range should be close to the Bose–Einstein one, $n_{\mathrm{O}} = [\exp(\epsilon + \mu/T) - 1]^{-1}$.

Taking into account that a region $\epsilon \sim 4T$ contributes mainly into the integral I, one can write

$$I \approx 24T^5\,\mathrm{e}^{-\mu/T}.$$

The double-Compton terms dominates equation (14) at low enough energies, $\epsilon \ll \epsilon_0$, where

$$\epsilon_0 = \sqrt{\frac{240\alpha}{\pi}\frac{T}{mc^2}}T. \tag{16}$$

So at these energies the spectrum should be close to Planckian, at $\epsilon \gg \epsilon_0$ Comptonization forms a Bose–Einstein spectrum.

One can find the chemical potential μ considering the photon balance in the source (Weaver & Chapline 1974; L). For this purpose let us multiple equations (6) and (14) by ϵ^2 and integrate them over all energies. The Comptonization term vanishes identically because it conserves the total number of photons. Eliminating the mode exchange term, one can write

$$\frac{\mathrm{d}}{\mathrm{d}x}\int \cos\theta\epsilon^2 n_{\mathrm{E}}\,\mathrm{d}\epsilon\,\mathrm{d}(\cos\theta) = 2N\int_0^\infty \epsilon^2 Q\,\mathrm{d}\epsilon. \tag{17}$$

The left-hand side of this equation describes the spatial transfer; deep inside the E-mode photosphere, $\tau \gg 1$, it may be reduced to the diffusion term (L; TD). The source term (the right-hand side) diverges logarithmically at $\epsilon \to 0$ if one substitutes the Bose–Einstein distribution. Taking into account that at $\epsilon \ll \epsilon_0$ the spectrum goes to the Planck distribution and $Q \to 0$ there, one can simply truncate the integral at $\epsilon = \epsilon_0$ and write the source term as

$$\int_0^\infty \epsilon^2 Q d\epsilon = \frac{32\alpha}{\pi}\sigma_{\mathrm{T}}\left(\frac{T}{mc^2}\right)^2 T^3 e^{-\mu/T}\ln\frac{\mu + \epsilon_0}{\epsilon_0}. \tag{18}$$

Because the mode exchange rate is comparable to the scattering rate of the E-photons, one can take $n_{\mathrm{E}} \sim n_{\mathrm{O}}$ (the more careful consideration (L) shows that $n_{\mathrm{E}} \approx n_{\mathrm{O}}$ at $\tau \gtrsim 1$). At the E-mode photosphere, $\tau \sim 1$, one can estimate the left-hand side of equation (17) as T_0^3/x. Substituting the estimate (18) into the right-hand side, one gets, with the aid of equation (11),

$$\ln\frac{\mu + \epsilon_0}{\epsilon_0} \sim \frac{\pi^3}{80\alpha}\left(\frac{B_{\mathrm{q}}}{B}\right)^2 = 0.5\left(\frac{10B_{\mathrm{q}}}{B}\right)^2.$$

[1] At the condition (1) the double-Compton scattering is a non-resonant process resembling the double-Compton scattering in a non-magnetized plasma. The double-Compton scattering in a strongly magnetized plasma studied by Bussard, Alexander & Meszaros (1986), Kirk, Nagel & Storey (1986) and Melrose & Kirk (1986) is a resonance process that may be considered as cyclotron absorption followed by double-photon decay of the excited Landau level.

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One can see that even at the E-mode photosphere the chemical potential is small, $\mu < \epsilon_0 \ll T$.

5 THE EMERGENT SPECTRUM

With the above result, the radiation transfer problem is greatly simplified. Substituting $n_{\mathrm{O}} = n_B(\epsilon, T)$ into equation (6) and making use of the definition (11), one gets

$$\cos\theta \frac{\mathrm{d}n_{\mathrm{E}}(\epsilon,\theta)}{\mathrm{d}\tau} = \frac{5}{4\pi^2}\left(\frac{\epsilon}{T_0}\right)^2 \left[- n_{\mathrm{E}}(\epsilon,\theta) + \frac{3}{8}\int n_{\mathrm{E}}(\epsilon,\theta')\,\mathrm{d}(\cos\theta') + \frac{1}{4} n_B(\epsilon,T)\right]. \quad (19)$$

This is the standard equation for the radiation transfer in a scattering medium with true absorption. The mode exchange process plays the role of true absorption because the spectrum of the O-mode is close to Planckian. This equation should be solved under the condition of the constant radiation flux. The iteration procedure by Unsöld and Lucy, as it was described by Mihalas (1978), was employed (see also Ulmer 1994). The temperature distribution is shown in Fig. 1, and the emergent spectrum in Fig. 2. This solution is consistent with the results by Ulmer (1994) if one takes into account that he presented the spectrum in terms of radiation intensity, $I \propto \epsilon F$. The approximate solution (10,13) fits the numerical solution well enough. One can see that the spectrum is flat at $\epsilon < T_{\mathrm{b}}$.

In the above considerations of the radiation transfer, the vacuum resonance effect (Pavlov & Shibanov 1979; Ventura, Nagel &

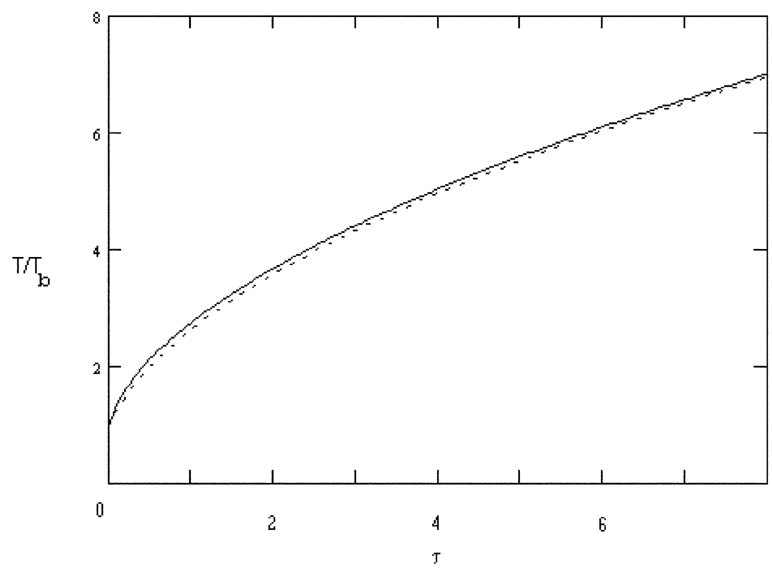

Figure 1. The temperature distribution within the photosphere. Solid line: numerical calculations; dashed: the approximate formula (10).

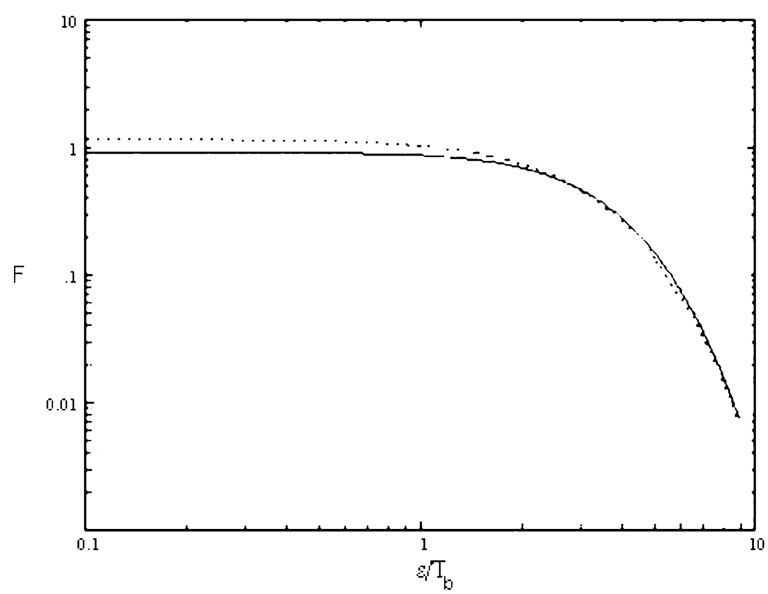

Figure 2. The spectrum of radiation escaping the bubble photosphere. Solid line: numerical calculations; dashed: the approximate formula (13).

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Mészáros 1979) was neglected. It is shown in the Appendix that at the conditions of interest the resonance energy is small, ~ 1 keV; in this band the interstellar absorption should anyway affect the observed spectrum. This also justifies our use of the cross-sections (2–5), which are valid when the plasma contribution into the dielectric tensor is negligible. The vacuum resonance may be important for the emission from the surface of the magnetar (Bulik & Miller 1997; Lai & Ho 2002; Özel 2001) where the plasma density is large.

Note that a flat spectrum is formed at low energies not only in the LTE case but also in the opposite case when the photon production rate is low and the chemical potential is correspondingly large (L). The spectral shape at low energies is determined only by the energy dependence of the scattering cross-section.

The radiation emerged from the bubble propagates through the magnetosphere where various processes may affect the radiation spectrum. The emergent radiation is linearly polarized perpendicularly to the local magnetic field (E-mode radiation). Owing to photon splitting (TD; Baring 1995), E-photons with the energy $\epsilon \gtrsim 30-40$ keV are converted into O-photons with the energy $\sim\epsilon/2$. On account of large scattering cross-sections, the O-mode photons are easily captured by the wind formed by a plasma ablated from the surface (TD); their spectrum is determined by processes in the wind. However at $\epsilon \lesssim 10$ keV the photon splitting does not affect the outgoing spectrum.

In the low-energy band, absorption at the ion cyclotron energy (TD; Thompson, Lyutikov & Kulkarni 2002; Zane et al. 2001),

$$\epsilon_g = 6.3 \frac{Z}{A} \frac{B}{10^{15}\,\mathrm{G}}\ \mathrm{keV},$$

where $Z/A = 1$ for hydrogen and 1/2 for other ions, should be taken into account. Because the magnetic field varies significantly along the ray path through the magnetosphere, a wide absorption trough should be formed in the spectrum. If the surface magnetic field of the star (which is typically not equal to the dipole component of the field) is of about $2 \cdot 10^{15}$ G or larger, the outgoing spectrum will be totally distorted at $\epsilon \lesssim 10$ keV. The low-energy turnover observed in the spectrum of SGR 1806-20 (Fenimore, Laros & Ulmer 1994) may be attributed to the cyclotron absorption. However in sources with lower magnetic fields there should be a range of energies within which the contribution of the photospheric radiation into the total spectrum is significant and therefore the total spectrum should be relatively flat in this range. A sketch of the outgoing spectrum is presented in Fig. 3. The optical depth of the magnetosphere with respect to photon splitting was chosen to be $\tau_{\mathrm{E}\to 2\mathrm{O}} = (\epsilon/30\,\mathrm{keV})^5$ because the splitting cross-section $\propto\epsilon^5$ and the magnetosphere becomes opaque for E-photons at $\epsilon > 30-40$ keV (TD). The photospheric spectrum was diluted by the factor $\exp(-\tau_{\mathrm{E}\to 2\mathrm{O}})$. The spectrum of O-photons produced by the splitting and reprocessed in the wind was chosen in the Wien form

$$F_\mathrm{O} = a\epsilon^2 \exp(-\epsilon/\epsilon_0),$$

where parameters a and ϵ_0 were determined from the conditions that the total amount of the O-photons is twice as much as the amount of the split E-photons and their total energies are the same. The total photon spectrum (E + O) is shown by the solid line; a trough at $\epsilon < 4$ keV demonstrates the influence of cyclotron absorption. One can see that if ϵ_g is not too high, the photon spectrum of the outgoing radiation is nearly flat at $\epsilon_g < \epsilon < 20$ keV. Polarization of this radiation is determined by the ratio of the E and O-mode photon fluxes; at low energies, where E-mode radiation from the bubble dominates, degree of polarization should be large.

The outgoing spectrum may be totally distorted at large radii where the magnetic field is about 10^{12} G and radiation cross-sections are large for both polarization modes. At these radii, the plasma forms a 'blanket', which may absorb and reprocess all radiation from the star (cf. Thompson et al. 2002 who considered reprocession of the persistent emission from the magnetar in a quiet state). Calculation of the observed spectrum with consideration for the reprocession at large radii is out of the scope of this article. One may only anticipate now that the highly super-Eddington radiation flux from the bubble pushes the plasma along the magnetic field lines toward the points, where the

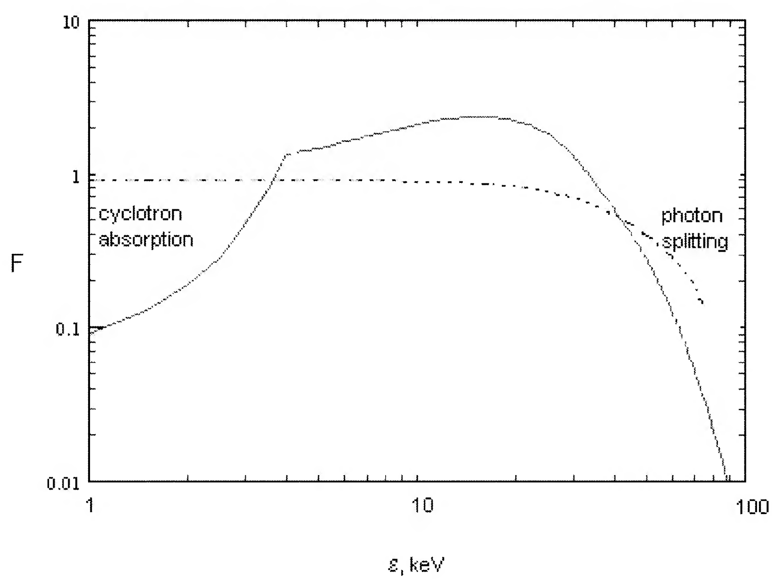

Figure 3. A sketch of SGR spectrum at $T_b = 15$ keV, $\epsilon_g = 4$ keV [$B = 6 \times 10^{14}(A/Z)$ G]. The photospheric spectrum is shown by dashed line. The photon splitting suppress the high-energy tail and forms an excess at $\epsilon \sim 10$ keV. The cyclotron absorption forms a trough at $\epsilon < \epsilon_g$.

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magnetic field is perpendicular to the radius and where the radiation force is compensated by the magnetic pressure, and therefore there should be windows in this 'blanket'. Then the spectrum shown in Fig. 2 may be observed via these windows.

6 CONCLUSIONS

We considered a spectrum of radiation escaping a photosphere of a bubble formed during the SGR outburst. The spectrum is flat at $\epsilon < T_\mathrm{b}$ (Fig. 2) due to the energy dependence of radiation cross-sections in a strong magnetic field. This spectrum is strongly distorted when radiation propagates through the magnetosphere. However if the magnetic field at the surface of the star is not too high, $B \le 10^{15}$ G, a flat spectrum may be observed at energies <20 keV.

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APPENDIX A: ON THE VACUUM RESONANCE

At $B \gg B_\mathrm{q}$ the vacuum resonant energy is (Bulik & Miller 1997)

$$\epsilon_c = \hbar\omega_p\sqrt{\frac{3\pi}{\alpha}\frac{B_\mathrm{q}}{B}} = 4\sqrt{N_{26}\frac{10B_\mathrm{q}}{B}}\,\mathrm{keV}, \tag{A1}$$

where $\omega_p = \sqrt{4\pi e^2 N/m}$ is the plasma density, $N = 10^{26} N_{26}\,\mathrm{cm}^{-3}$. One can see that the vacuum resonance becomes significant at densities $N_{26} \ge 1$. The baryon density in the bubble is small (TD) however the density of electron–positron pairs grows with the depth. Substituting the equilibrium pair density (TD)

$$N = \frac{emB}{\sqrt{2\pi^3}\hbar^2}\sqrt{\frac{T}{mc^2}}\exp\left(-\frac{mc^2}{T}\right)$$

into equation (A1), one can express the resonance energy via the local temperature as

$$\epsilon_c = \sqrt{6}mc^2\left(\frac{2\pi T}{mc^2}\right)^{1/4}\exp\left(-\frac{mc^2}{2T}\right). \tag{A2}$$

Note that the result is independent of the magnetic field. The temperature grows with depth according to equation (10) and therefore ϵ_c grows with depth. An extraordinary photon with the energy ϵ_c escapes from the depth τ if (see equation (12)

$$\frac{5}{4\pi^2}\left(\frac{\epsilon_c}{T_\mathrm{b}}\right)^2\tau \approx 1. \tag{A3}$$

Eliminating τ and T from equations (10), (A2) and (A3) in favor of ϵ_c, one gets

$$\epsilon_c = \left\{1 + \frac{1}{20}\ln\left[\left(\frac{1\,\mathrm{keV}}{\epsilon_c}\right)^{2.5}\frac{T_\mathrm{b}}{10\,\mathrm{keV}}\right]\right\}\left(\frac{T_\mathrm{b}}{10\,\mathrm{keV}}\right)^2\mathrm{keV}.$$

One can see that at $\epsilon > 1{-}2$ keV the effect of vacuum resonance does not affect the spectrum of the emergent radiation.

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