Non-pulsed gamma radiation from binary system with a pulsar

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ABSTRACT
Consider binary system with a millisecond pulsar ejecting relativistic particles and an optical star emitting soft photons with the energy $\omega \approx 1 - 10$ eV. These low-energy photons are scattered by the relativistic electrons and positrons of the pulsar’s wind. The scattered photons forms a wide spectrum from hard X-ray band up to gamma band $\varepsilon \approx 1 - 1000$GeV. When the pulsar wind is isotropic the luminosity of gamma radiation $L_{\gamma} = L_{\gamma}(\psi)$ depends heavily on the angle $\psi$ between the directions to the optical star and to the observer from the pulsar. During the orbital motion this angle varies periodically giving rise to the periodical change of the observed intensity of the gamma radiation and its spectrum. We calculated the spectral shape of the scattered hard photons. Under the assumption that the energy losses of the relativistic particles are small we receive analytical formulas. We apply our results to the binary system PSR B1259-63 and show that if the wind from the Be star is accounted for then it is possible to reproduce the observed spectrum.

Key words: binaries:general - pulsars:general - pulsars:individual:PSR B1259-63 - X-rays:stars

1 INTRODUCTION
The principal accessible store of energy in a pulsar is its rotational energy, which is liberated at a rate $L_p \approx I \frac{\partial \omega}{\partial t} \approx 10^{32} \div 10^{36}$erg/s (see, for example, a table of pulsars parameters ($L_p, p, p$) given in the book of Beskin et al.,1993 ), where $I \approx 10^{45}$g cm$^{-2}$ moment of inertia, $p$ is a period, $p$ is a deceleration of a pulsar. The bulk of the pulsar energy $L_p$ is transferred to the pulsar wind which consists of electrons, positrons and probably heavy ions, $L_{\omega} = gL_p$, $g \leq 1$. The Lorentz factor of the relativistic particles in the wind may vary in range $\gamma \approx 10 - 10^9$ (e.g. Manchester&Taylor,1977). The intrinsic gamma ray luminosity of pulsed emission from the short periodic pulsars is of the order of $L_{\gamma}/L_p \approx 0.01$ (see Arons, 1991). Below we discuss the different (induced) mechanism of non pulsed gamma radiation generation in binary with a pulsar and an optical star. This mechanism could result in considerably higher ratio of $L_{\gamma}/L_p$. The preliminary results of this work were published in the paper of Chernyakova&Illarionov, 1997.

Consider the case of binary with a pulsar ejecting relativistic particles and an optical star emitting soft photons in optic and UV band with the energy $\omega \approx 1 - 10$ eV. These low-energy photons are scattered by the pulsar wind relativistic electrons and positrons. The energy of the photon after the inverse Compton scattering is very high - $\varepsilon_{\max} \approx \omega \gamma^2$ in the Thomson limit ($\omega \gamma \ll mc^2$) and $\varepsilon_{\max} \approx mc^2\gamma$ in the opposite case ($m$ is a mass of an electron, $c$ is a light velocity). The scattered photons form a wide spectrum from hard X-ray band up to gamma band $\varepsilon \approx 1 - 1000$GeV. The relativistic particle scatters the photon preferably along the direction of the particle velocity. As a result while soft photons are directed radially from the optical star, the scattered hard photons move radially from the pulsar. Here and below we assume that the pulsar wind particles are radially directed from the pulsar inside the effective scattering volume.

In the case of the presence of an obstacle for the pulsar wind and mainly when the matter flow from the optical star is rather intensive the radial flow of the relativistic wind is destroyed. Then the system of shock waves resulting from the collision of two winds appears between the pulsar and the optical star and the trajectories of the electrons and positrons beyond the shock change. In the work of Tavani and Brookshaw (1991) the case of weak (in comparison with the pulsar wind) matter outflow is discussed. The hydrodynamics of the collision between the relativistic and the nonrelativistic winds is closely analogous to the hydrodynamics of the two nonrelativistic winds collision, which was intensively discussed by different authors applied to the WR+OB binaries since the work of Prilutskii&Usov (1976).

The total luminosity of scattered hard photons $L_{\gamma}$ is

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equal to the total particle energy losses \( L_{\text{loss}} \) through the inverse Compton scattering in the course of the particle motion from the pulsar to the infinity. To estimate \( L_\gamma \) we note that the rate of the energy losses of a relativistic particle moving in a radiation field with energy density of soft photons \( w_{\text{soft}} \approx L_\gamma / 4\pi c^2 \) is about \( mc^2 d\gamma / dt \approx -w_{\text{soft}} \sigma_T c^2 \gamma^2 \) in the Thomson limit. Here \( L_\gamma \) is a luminosity of the optical star, \( \sigma_T \) is a binary stars separation and \( \sigma_T = 8\pi (c^2 / m c^2)^2 = 6.65 \times 10^{-25} \text{cm}^2 \) is the Thomson cross-section. Hence the decrease of the Lorentz factor of a particle is \( \Delta \gamma \sim \gamma \left[ 1 - \left(1 + \gamma / \gamma_*\right)^{-1} \right] \), \( \gamma_* = 4\pi am c^3 / \gamma_T L_* \).

So the rate of all particles total energy loss is

\[
L_\gamma = L_{\text{loss}} = L_w \Delta \gamma / \gamma \sim L_w \left(1 - \frac{1}{1 + \gamma / \gamma_*} \right).
\]

Thus in the case \( \gamma \ll \gamma_* \), the gamma-ray luminosity from binary is proportional to the luminosity of the optical star \( L_* \) and to the luminosity of the wind \( L_w \). \( L_\gamma = K L_w \), where the transformation parameter \( K = \gamma / \gamma_* \). In the close binary system with an optical star with high luminosity \( \gamma \gg \gamma_* \).

In this case practically all energy of the wind transfers to the energy of scattered photons \( L_{\gamma, \text{max}} \approx L_w L_* \).

In each unit of the volume in the region where the gamma radiation is generated the source function of gamma radiation is highly anisotropic. When the pulsar wind is isotropic the gamma luminosity \( L_\gamma \) ( \( L_\gamma = \int 2\pi L_\gamma (\psi) d\cos \psi \) ) has an azimuthal symmetry around the line connected binary companions and depends heavily on the angle \( \psi \) between the directions to the optical star and to the observer from the pulsar, see Figure 1. In our case of free radial relativistic wind the binary system emits the maximum energy of gamma radiation in the direction of the star \( (\psi = 0) \) and the minimum energy of radiation in the opposite direction \( (\psi = \pi) \). The spectrum of the radiation and the maximal radiated energy also depend on \( \psi \).

During orbital motion \( \psi \) varies periodically giving rise to the periodical change of the intensity of the gamma radiation coming from the binary system to the observer.

In section 2 we calculate the spectral shape \( L_\gamma (\varepsilon, \psi) \) of the scattered hard photons in the case of arbitrary value of the parameter \( \Delta \gamma \), going beyond the Thomson limit. Under the assumption \( K \ll 1 \) we receive analytical formula for the \( L_\gamma (\varepsilon, \psi) \).

In section 3 we apply our results to the binary system PSR B1259-63 and find that under the assumption of the power law relativistic particles spectrum, \( \frac{dN}{d\varepsilon} = 0.4L_w \left[ mc^2 \left( \gamma_{\text{min}}^{0.4} - \gamma_{\text{max}}^{0.4} \right) \right] \gamma^{-2.4} \) in the range \( 10 < \gamma < 500 \), our model describes the observe photon spectrum rather good but the intensity is less than the observed one by a factor about 30. This discrepancy is due to the presence of the mass outflow from the Be star which disturbed the free flow of the pulsar wind. The centrally located shock appears between the pulsar and the star due to the interaction between the two winds. The big differences between the values of the velocities of the particles from the different sides of the tangential discontinuity will lead to the growth of the in-stabilities and the two winds will be macroscopically mixed between the shocks. Then the heavy non relativistic wind slows down the volumes filled by the relativistic electrons and positrons and they acquire essentially non relativistic hydrodynamic drift velocity \( v_d \) along the shock while the energy of electrons and positrons does not changes significantly. With the decrease of the hydrodynamic velocity of the relativistic plasma the time which it spends near the optical star increases in \( c/v_d \) times. The effective transformation parameter \( K_{\text{eff}} \sim \frac{\gamma}{\gamma_*} K \) thus can be large enough to overcome the discrepancy between the simple theory and observations.

## 2 Generation of Gamma Radiation

Let us consider an interaction between the relativistic pulsar wind and the soft radiation from the companion and find the spectral and angular dependence of the outgoing hard radiation \( L_\gamma (\varepsilon, \psi) \). We assume that the pulsar wind and the soft emission from the companion are isotropic and that there is no mass outflow from the optical star. We treat both the pulsar and the optical star as a point sources. In this case the trajectories of the particles and soft non scattered photons are directed radially from the pulsar and from the optical star correspondingly.

Let I denote the location of the small element of volume \( dV \) at a distance \( r \) to the pulsar and at an angle \( \theta_2 \) to the line of sight \( P\bar{O} \) (see Figure 1). In this element of volume optical photons with the energy \( \omega \) scatters by the relativistic particles (electrons or positrons). The angle between the directions of the photon and the particle movements before the interaction is denoted by \( \theta_1 \). The photon scattering angle is designated by \( \theta \). Let \( l \) signifies the distance from this volume to the companion \( S \).

We are interested in the case of a distant observer and thus a vector along the direction to the observer from any point near the system may be considered as parallel to such a vector from the pulsar. Only the photons scattered in the direction of the observer within the small solid angle \( \Omega = S \cos \chi / D^2 \) will reach the observer. Here \( S \) is the area of the observer’s surface, \( D \) is a distance from the pulsar to the observer and \( \chi \) is an angle between the direction to the observer from the pulsar and from the volume. As \( \chi \) is of the order of \( a / D \ll 1 \), we may neglect the variance of \( \Omega \) from point to point. As we have already mention relativistic particle scatters the photons preferably along the direction of the particle velocity and thus \( \theta_2 \ll 1 \). From this fact it follows that angles \( \theta \) and \( \theta_1 \) depend only on \( r \) and with an accuracy of the order \( \theta_2 \sim \sqrt{\omega / \varepsilon} \) are equal.

### 2.1 Anisotropy of the energy radiated from the binary system in the Thomson limit

Lets calculate the \( \psi \)-dependence of the luminosity of the hard radiation emitted from the binary system \( L_\gamma (\psi) \) in the Thomson limit (\( \omega \gamma / mc^2 \ll 1 \)). The luminosity of scattered hard photons \( L_\gamma (\psi) \) is equal to the particle energy losses in the course of the particle motion from the pulsar to the infinity. So at first lets calculate the total energy losses \( L_{\text{loss}} \) of the electron moving along the radial trajectory at an angle \( \psi \) to the line \( PS \) in the radiation field of the optical star. The rate of the relativistic (\( \gamma \gg 1 \)) electron’s energy losses is given by

\[
\frac{1}{c} \frac{d\gamma}{dt} = \frac{d\gamma}{dr} = -\frac{L_* \sigma_T \gamma^2 \left(1 - \cos \theta_1 \right)^2}{4\pi mc^3 \left( l^2 \right)}.
\]
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where

\[ l = \sqrt{(r - a \cos \psi)^2 + a^2 \sin^2 \psi}, \]
\[ \cos \theta_1 = (r - a \cos \psi) / l. \]

Upon integrating (1) over radius from \( r = 0 \) to \( r \) we obtain the dependence of the Lorentz factor of the electron on \( r \):

\[ \gamma = \gamma_0 \left[ 1 + K \Phi(\psi) + \frac{K}{\sin \psi} \left( \frac{2}{\sqrt{1 + x^2}} - 3 \frac{\arccot x}{2} - \frac{x}{2(1 + x^2)} \right) \right]^{-1}, \tag{2} \]

where

\[ \Phi(\psi) = \frac{3(\pi - \psi)}{2 \sin \psi} - 2 - \frac{\cos \psi}{2} \]

is a beaming function which as it will follows from (5) represents the angular dependence of the energy losses of the electrons.

Energy radiated from the binary system is determined as

\[ x = \frac{r/a - \cos \psi}{\sin \psi} \]

and \( \gamma_0 \) is the initial Lorentz factor of the electron. Letting \( r \to \infty \) we find from (2) the total change of the Lorentz factor \( \Delta \gamma(\psi) = \gamma_0 - \gamma(x = \infty, \psi) = \gamma_0 \left[ 1 - (1 + K \Phi(\psi))^{-1} \right] \).

As the bulk of the energy losses of the relativistic electron transfers to the scattered photons moving in the direction of the electron movement the angular dependence of the energy radiated from the binary system \( L_\gamma(\psi) \) agrees with the angular dependence of the energy losses of the electrons. The total energy losses of the isotropic flow of all relativistic particles in the direction of the observer located at an angle \( \psi \) to the line \( PS \) in a small solid angle \( \Omega \) are proportional to the amount of particles moving in this direction \( \sigma_c L \Omega / 4 \pi = L_\gamma L \Omega / 4 \pi \gamma_0 mc^2 \) in the case of monoenergetic pulsar wind, and to the total energy losses of one particle \( mc^2 \Delta \gamma(\psi) \)

\[ L_\gamma(\psi) \Omega \approx L_{\text{loss}}(\psi) \Omega = \frac{L_w \Delta \gamma(\psi)}{4 \pi \gamma_0} \Omega = \frac{L_w}{4 \pi} \left( 1 - \frac{1}{1 + K \Phi(\psi)} \right) \Omega. \tag{5} \]

This is the energy transferred to the gamma band. Figure 2 shows the dependence \( \Phi(\psi) \). The beaming function \( \Phi(\psi) \) is rapidly decreasing function. At \( \psi \ll 1 \Phi(\psi) \approx \pi / \psi \) and at \( \pi - \psi \ll 1 \Phi(\psi) \approx (\pi - \psi)^4 / 120 \). Equation (5) is valid for the arbitrary value of \( K \) but it has no use for \( \psi \) less then \( R_\ast / a \) and bigger then \( R_\ast / a \) (\( R_\ast \) - is the radius of the optical star) as the optical star is not a point source. That’s why we don’t worry about the fact that \( \Phi(\psi) \to \infty \) as \( \psi \to 0 \).

If \( \psi \) is not too small then it follows from (5) that under the condition \( K \ll 1 \) the energy decrease is small and

\[ L_\gamma(\psi) = \frac{L_w \gamma_0 \Phi(\psi)}{4 \pi \gamma_0}. \tag{6} \]

Upon integrating (6) over \( \psi \) from \( \psi = 0 \) to \( \psi = \pi \) we obtain the total energy transferred to the gamma band in the case \( K \ll 1 \)

\[ L_\gamma = \left( \frac{3}{8} \pi^2 - 2 \right) \frac{\gamma_0 L_w}{\gamma_\ast}. \tag{7} \]

2.2 The spectral and the angular dependence of the outgoing radiation

Now lets calculate the spectral and the angular dependence of the outgoing radiation \( L_\gamma(\omega, \varepsilon, \gamma, \psi) \) in the case of arbitrary value of the parameter \( \omega \gamma / mc^2 \), going beyond the Thomson limit. The luminosity of the scattered quanta moving at an angle \( \psi \) to the line \( PS \) within a unit solid angle in a unit of time in the energy range from \( \varepsilon \) to \( \varepsilon + d\varepsilon \) may be written in the following form:

\[ L_\gamma(\omega, \varepsilon, \gamma, \psi) d\varepsilon = \int \varepsilon n_\pm n_\text{c} (1 - \beta \cos \theta_1) \sigma_K d\varepsilon, \tag{8} \]

where \( \beta = v/c \), \( n_\pm \) is the density of the particles, \( n_\text{c} = L_\ast / (4 \pi c^2 \omega) \) is the density of the optical photons, \( \sigma_K(\theta_1, \theta_2, \omega, \varepsilon) \) - is the Klein-Nishina cross-section (Jauch & Rohrlich 1976)

\[ \sigma_K = \frac{3 \sigma_T}{16 \pi \gamma^2} \omega^2 \left( 1 - \beta \cos \theta_1 \right)^{-2} \times \]

\[ \left[ \frac{1 - \cos \theta}{\gamma^2 (1 - \beta \cos \theta_1 (1 - \beta \cos \theta_2))} - 1 \right]^2 + 1 \]

\[ + \left( \frac{\varepsilon}{mc^2 \gamma} \right)^2 \left( 1 - \cos \theta_2 \right)^2 \left( 1 - \beta \cos \theta_1 \right) \].

The energy of the scattered photon is (Berestetskii et al. 1971)

\[ \varepsilon = \omega \frac{1 - \beta \cos \theta_1}{1 - \beta \cos \theta_2 + \frac{\omega}{mc^2 \gamma} (1 - \cos \theta_1)}. \tag{10} \]

From (10) it follows that

\[ \cos \theta_2 = \beta^{-1} \left[ 1 - \frac{\omega}{\varepsilon} (1 - \beta \cos \theta_1) + \frac{\omega}{mc^2 \gamma} (1 - \cos \theta_1) \right]. \tag{11} \]

We are interesting in case \( \gamma \gg 1 \) and in all formulas that followed we use the expansion \( \beta = 1 - \frac{1}{8 \pi \gamma} \). In the spherical coordinates \( d\varepsilon = 2 \pi r^2 dr d\cos \theta_2 \) is \( 2 \pi \gamma^2 (1 - \cos \theta_1) d\varepsilon \). When \( \varepsilon \gg \omega \), then \( \theta_2 \ll 1 \) and thus

\[ \cos \theta_1 \approx \cos \theta = \frac{r/a - \cos \psi}{\sqrt{(r/a - \cos \psi)^2 + \sin^2 \psi}}. \tag{12} \]

Now it is possible to rewrite (9) as an integral over \( r \)

\[ L_\gamma(\omega, \varepsilon, \gamma, \psi) = \frac{N_\pm n_\pm \omega}{8 \pi a^\gamma \varepsilon} \int_0^{r_{\max}} \left( 1 - \cos \theta_1 \right)^2 \sigma_K \frac{dr}{(r/a - \cos \psi)^2 + \sin^2 \psi}. \tag{13} \]

where \( N_\pm \) is the number of photons going from the companion per second and \( N_\ast \pm \) is the number of particles leaving the pulsar per second, \( n_\pm = N_\omega / 4 \pi c^2 \), \( n_\pm = N_{\ast \pm} / 4 \pi c^2 \), \( r_{\max} \) is determined from the condition \( \cos \theta_2 = 1 \). If Lorentz factor of a particle is constant then from (14), (15) follows:
to obtain:

\[ l = \frac{4\omega\gamma^2}{\varepsilon}, \quad \lambda = 1 - \frac{\varepsilon}{mc^2\gamma}. \]

In the case \( K \ll 1 \) in Thomson limit the Lorentz factor of the particles is practically constant in the effective scattering volume. In the case \( \omega\gamma/mc^2 \gg 1 \) this statement is also correct. In this case if the mean free path of a particle is much bigger then the binary separation: \( l_{free} = \sigma_K\gamma_{m}\langle l \rangle /4\pi\omega^2c \gg a \) (\( \sigma_K_{m} \) is a full Klein-Nishina cross-section) then the influence of the field of soft photons on the spectrum and flux of the relativistic particles may be neglected. In the highly relativistic case \( \omega\gamma/mc^2 \gg 1 \) \( \sigma_K_{m} \sim \sigma_Y mc^2/\omega\gamma \) and we note that from \( K \ll 1 \) it follows that \( l_{free} \gg a \).

After a substitution of the expression \( L_{\gamma} \) in (13) and the following integration under the assumption that \( \gamma \) is constant we obtain:

\[ L_{\gamma} = \frac{N_e N_\omega}{\gamma^2} F(\omega, \varepsilon, \gamma, \psi), \quad (15) \]

where the constant \( \kappa = \frac{3\omega mc^2}{2\sin c} \) and in the range

\[ \omega \ll \varepsilon \leq \varepsilon_{max} = 2\gamma^2\omega - \cos \psi + \frac{2\omega mc^2}{\gamma}(1 + \cos \psi) \]

the function \( F(\omega, \varepsilon, \gamma, \psi) \) is determined as follows:

\[ F = \frac{1}{\sin \psi} \frac{1}{\sin \psi} \left( \lambda + \frac{1}{\lambda} \right) \left( 2 \arctan z - \psi \right) \]

\[ \left( \frac{3\mu^2}{2\lambda^2} \right) (2z^3 - 3z^2 \tan \frac{\psi}{2} + \tan \frac{3\psi}{2}) \right), \quad (17) \]

and is equal to zero when \( \varepsilon > \varepsilon_{max} \). When \( \omega \ll \varepsilon \ll \varepsilon_{max} \) the equation (17) simplifies:

\[ F = \frac{2}{\sin \psi} \frac{\pi - \psi}{\sin \psi}. \]

In the Thomson limit \( \varepsilon_{max} = 2\gamma^2(1 + \cos \psi) \) and hence

\[ 1 - \lambda = \frac{2\omega mc^2}{\gamma}(1 + \cos \psi) \ll 1. \]

Under these conditions equation (18) becomes

\[ F = \frac{1}{\sin \psi} \frac{1}{\sin \psi} \left( 4 \arctan z - 2\psi \right. \]

\[ \left. - \frac{8}{3\mu^2} \left( 2z^3 - 3z^2 \tan \frac{\psi}{2} + \tan \frac{3\psi}{2} \right) \right). \quad (18) \]

If we integrate equation (13) with \( F = F_T \) over the energy of the scattered photons \( \varepsilon \) we receive the \( \psi \)-dependence of the total outgoing energy in accordance with the equation (13). The relationship between functions \( \Phi(\psi) \) and \( F_T \) is:

\[ \Phi(\psi) = \frac{3}{8} \int_0^{\varepsilon_{max}} \frac{\varepsilon}{\omega^2} F_T d\varepsilon \]

Under the real conditions the spectrum of the optical photons is not monochromatic. If the companion’s radiation is similar with a radiation from the black body with temperature \( T \), then the number of soft photons in the energy range from \( \omega \) to \( \omega + d\omega \) is:

\[ dN_\omega = \frac{15L_\gamma}{\pi^2 T^4} \frac{\omega^2 d\omega}{\omega^2 - T^2}. \]

Replacing \( N_\omega \) in formula (14) with \( dN_\omega \) and integrating over \( \omega \) we receive:

\[ L_\gamma(\omega, \varepsilon, \gamma, \psi) = \int F(\omega, \varepsilon, \gamma, \psi) \frac{\omega d\omega}{(\omega^2 - T^2)} (15). \]

After formulas (14) - (20) were received under the assumptions that Lorentz factor is constant. If we consider in the Thomson limit the dependence of the energy \( \varepsilon_{max} \) on \( \gamma \) (according to (13) then from formula (13) the luminosity of the scattered radiation is:

\[ L_{\gamma}(\omega, \varepsilon, \gamma_0, \psi) = \tau N_e N_\omega \frac{\varepsilon}{\omega_{00}} F_\gamma(\omega, \varepsilon, \gamma_0, \psi), \quad (21) \]

where

\[ F_\gamma(\omega, \varepsilon, \gamma_0, \psi) = \frac{r_{max}}{\gamma^2} \int \frac{\gamma^{-1} - 1}{(r/a - \cos \psi)^2 + \sin^2 \psi} \frac{dr}{a}. \quad (22) \]

The dependences of \( r_{max} \) and \( \cos \theta_2 \) on \( r \) are determined by the formulas (13), (21).

Figure 7 represents spectra of the scattered radiation obtained in the Thomson limit for different values of the parameter \( K = \gamma_0/\gamma_s \) with \( \gamma_0 = 10^6 \). Spectra are calculated taking into account the dependence of \( \gamma \) on \( r \). Curve \( K = 0 \) corresponds to the case of constant Lorentz factor, which is described by formula (3). It can be seen from this figure that the bigger is the value of parameter \( K \) at a given angle, the softer is the spectrum of the hard radiation. This figure also illustrates that with the growth of parameter \( K \) the total energy of the outgoing radiation increases tending to \( L_{\gamma} \) as \( K \rightarrow \infty \).
3 Discussion

We considered the scattering of the soft photons emitted by the optical star by the relativistic electrons and positrons from the pulsar wind. We showed that the intensity of the hard radiation came about from such a scattering is proportional to the luminosities of the pulsar and the optical star. But it is worth to mention that the luminosity of optical star itself depends on the luminosity of the pulsar as it absorbs and reradiate part of the energy of the pulsar wind. If we denote by $L_{\text{st}}$ the intrinsic (un-illuminated) luminosity of the star and by $L_{\text{ind}} = f L_{p} R_{p}^{2}/4a^{2}$ the energy of the pulsar reradiated by the optical star ($f$ is the fraction of the energy flux from the pulsar incident on the surface of the optical star that is converted into soft emission) then the total luminosity of the star is $L_{\text{tot}} = L_{\text{st}} + L_{\text{ind}}$. PSR 1957-20 is an example of the system where the illumination of the optical star by a pulsar affects the luminosity of the optical star. A detailed study of the optical light curve of PSR 1957+20 done by Callanan et al. (1995) shows that the secondary is close to fill its Roche lobe with 7% - 20% of the incident flux converted to optical emission. Thus the described mechanism of the generation of the hard radiation is important not only in case of a highly luminous optical star in binary but also in case of a close binary with bright pulsar and an optical star with a big radius.

3.1 Comparison with the observations.

The only known binary system which contains radio pulsar and emits non pulsed X-ray radiation is PSR B1259-63 system. Since its discovery, the PSR B1259-63 system has been observed several times at X-ray energies. ROSAT observed the PSR B1259-63 system near to apastron in 1991-1992 (Cominsky et al. 1994, Greiner et al. 1995). In January 1994 the X-ray emission from the system PSR B1259-63 during the periastron passage was observed by telescopes ASCA and OSSE (Grove et al. 1995, Kaspi et al. 1995). ASCA also observed the post-periastron emission from this system on February 28, 1994.

The PSR B1259-63 system contains the radio pulsar with the spin period $P = 47.76$ ms, rotating around the massive Be star SS 2883. According to Manchester et al., 1995 the orbital eccentricity is $e = 0.87$, the projected semimajor axis is $a_{\text{maj}} \sin i = 3.9 \times 10^{13}$ cm and the longitude of periastron is $\omega = 139^\circ$. Spin-down luminosity of the pulsar is $L_{p} \simeq 9 \times 10^{35}$ erg/s. The mass function $f(M_p) = \frac{(M_p \sin i)^3}{(M_p + M_{\text{st}})} = 1.53 M_{\odot}$, the luminosity of the Be star is $L_{\text{st}} = 2.2 \times 10^{38}$ erg/s (Johnston et al. 1992, 1994). The distance between companions at periastron is $a = a_{\text{maj}}(1 - e) \sim 10^{13}$ cm.

The effective transformation parameter $K_{\text{eff}} \sim \frac{C_{\gamma}}{L}$ thus can be large enough to overcome the discrepancy between the simple theory and observations.

Under the assumption of the power law energy distribution of the relativistic electrons and positrons after the shock $\frac{dN_{\gamma}}{d\gamma} = A \gamma^{-s}$ and their isotropic velocity distribution we can estimate the intensity of the X-ray radiation from the system. Such a distribution can be either the result of the particle acceleration on the shock or in the case of the intrinsic power law distribution of the electrons and positrons in the pulsar wind and thin collisionless shock, passing which the relativistic particles do not change their energy distribution.

Lets consider an element of volume $dV$ filled with the relativistic plasma locating in the shock on a distance $R$ from the Be star. In the case of the relativistic electrons and positrons isotropic distribution the number of particles produces the photons with energy $\varepsilon$ moving in the direction of the observer in a unit solid angle is $\frac{dN_{\gamma}}{d\gamma} d\varepsilon d\Omega d\gamma = \frac{0.5 \varepsilon^{2}}{c} \frac{d\varepsilon}{d\gamma} d\Omega = \frac{0.5 \varepsilon^{2}}{c} \frac{d\varepsilon}{d\gamma} \cos \theta \cos \psi d\Omega$ (see section 2.2) Then according to the number of photons coming to the observer in the unit of time per unit square per MeV is

$$F_{\text{uv}}(\varepsilon, \psi) = \int_{\gamma_{\text{min}}}^{\gamma_{\text{max}}} F_{bb}(\varepsilon, \gamma, \psi) \gamma^{-s-2} d\gamma.$$ (24)
\[ dN = \frac{3\sigma_T AM_{p,c}}{32\pi \omega D^2} \times \int_{\sqrt{2}\omega(1-\cos \theta)}^{\infty} \gamma^{-s-2} \left[ 1 + \left( 1 - \frac{\varepsilon}{\omega \gamma^2 (1 - \cos \theta)} \right)^2 \right] d\gamma dV (25) \]

where \( \theta \) is a photon scattering angle. For the simplicity we take the photon density along the shock constant and equal to the one at a distance \( R = 10^{13} \text{cm} \). The factor \( AV \) can be estimated from the energy conservation law by equating the energy entering the shock per second \( L_B \Omega / 4\pi \) and the energy leaving the shock per second \( AV mc^2 \gamma_{\text{min}}^{2-s} / a(s-2) \), where \( \Omega \) is the solid angle under which the shock wave is seen from the pulsar. Thus we have \( AV \sim L_B \Omega (s-2) / 4\pi mc^2 \gamma_{\text{min}}^{2-s} \).

Then integrating (25) over the volume of the shock we have

\[ N \sim \frac{2\pi}{D^2 R^2} \frac{L_B \Omega}{mc v_d} \frac{(s-2)(2\omega)^{(s-3)/2}}{\gamma_{\text{min}}^{2-s}} \frac{\pi}{(1 + s)} e^{-(1+s)/2}. \]  

For \( \gamma_{\text{min}} = 10, s = 2.4, D = 2\text{kpc}, v_d = 10^8 \text{cm/s} \) we have

\[ N \sim 5 \times 10^{-3} \left( \frac{\varepsilon}{100 \text{keV}} \right)^{-1.7} \text{ph/s/cm}^2 / \text{MeV} \]  

while the observable value of the radiation at a 100 keV is \( 2.8 \times 10^{-3} \text{ph/s/cm}^2 / \text{MeV} \). Thus if the small drift velocity is taken into account it is possible to explain the observed spectral shape and intensity of the X-ray radiation. It is also worth to mention that for the particles with big Lorentz factor \( \gamma > \gamma_{\text{min}} v_d / c \) the assumption of constant Lorentz factor will be not valid due to the inverse Compton losses and the break in the photon spectrum will appear. The break in the photon spectrum at \( \varepsilon = \omega \gamma_{\text{min}}^2 v_d^2 / c^2 \) will appear. The index of the photon spectrum after the break will be bigger then the original one at 1/2. It can be also seen from (25) that the focus effect will take place - the intensity of the radiation is proportional to the \( (1 - \cos \theta)^{(1+s)/2} \) and thus the major part of the radiation will be emitted toward the direction of the optical star.

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Figure Captions

1) The geometry under consideration. $P$ is a pulsar, $S$ is an optical star, $I$ is a point of interaction of the electrons and photons, $O$ is an observer. Note that these points are not in one plane in general case.

2) The dependence of the beaming function $\Phi$ on an angle $\psi$.

3) The dependence of $F_{bb}$ on $\varepsilon$ for different values of $\psi$ for $\gamma = 10^4$ and $\gamma = 10^6$.

4) The $\psi$-dependence of the energy gone with the scattered photons in a unit of time in a unit solid angle. The energy is normalized to $C = \frac{15\pi^2}{4} N_s \gamma_L^2 r^2$.

5) The dependence of $F_{bb}$ on $\varepsilon$ for different positions of the companions in the circular orbit for $\gamma = 10^4$ and $\gamma = 10^6$.

6) The orbital angle $\varphi$ dependence of the total energy gone with the scattered quanta in the unit of time in the unit solid angle for the different values of inclination angles a) $\gamma = 10^4$, b) $\gamma = 10^6$.

7) Spectra of the scattered radiation obtained for different values of the parameter $K = \frac{\gamma_0}{\gamma_*}$. Spectra are calculated taking into account the dependence of $\gamma$ on $r$.

8) We apply our model to the system PSR B1259-63. Theoretical result (solid line), OSSE spectrum of emission from 1994 January 3-23 (black dots) and schematic extrapolations to the power-law spectra of ASCA observations from 1993 December 28 (dashed line), 1994 January 10 (dotted line), and 1994 January 26 (dashed-dotted line) are shown. ASCA extrapolations obtained from the analysis of Kaspi et al. 1995, OSSE results are taken from the paper of Grove et al. 1995.
Figure 3
Figure 4
Figure 5
Figure 6
\[ \psi = \pi/4, \Phi(\pi/4) = 2.65 \]

\[ \psi = \pi/2, \Phi(\pi/4) = 0.36 \]

**Figure 7**
Figure 8

\[ N_{e^-}(\gamma) \sim \gamma^{-2.4} \]

\( \gamma_{\text{min}} = 10 \)

\( \gamma_{\text{max}} = 500 \)