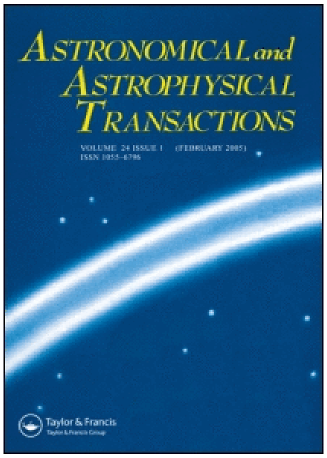

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INTERSTELLAR BUBBLES WITH A VARIABLE ENERGY INPUT RATE

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Influence of the energy input rate change on the interstellar bubble dynamics and X-ray luminosity is discussed. Analytical formulae for spherically-symmetric bubble with a power-law energy input rate are derived. The properties of the nebula Z364 in M33 galaxy and of the nebula N51D in the Large Magellanic Cloud are discussed.

1 INTRODUCTION

A fundamental property of massive stars is the powerful mass outflow (stellar wind) driven by the radiation pressure and possibly by other mechanisms (Gershberg and Shcheglov, 1964; Morton, 1967; Bisnovatyi–Kogan and Nadyozhin, 1972; Willis and Garmany, 1987; Lozinskaya, 1992). It has recently been recognized that compact groups of massive stars concentrated (OB-associations) within limited region of space inject enormous amount of energy into surrounding medium and influence strongly the dynamics, energy balance, and other properties of the interstellar medium in spiral and irregular galaxies. The hot bubble interior buffers the discrete supernova explosions that allows to treat them as a source of continuous energy input (Mac Low and McCray, 1988) as well. Thus the theory of the stellar wind interaction with interstellar medium developed by Pikel'ner (1968), Avedisova (1971), Castor *et al.* (1975), Weaver *et al.* (1977) can be applied to both of phenomena.

Bochkarev and Lozinskaya (1985) have applied this theory to predict X-ray emission from eight ring nebulae around WR and O_f stars. A weak diffuse X-ray emission from the stellar wind bubble around the WR star HD192163 was discovered by Bochkarev (1988). But the X-ray flux in the 0.2-3 keV band was estimated to be an order of magnitude lower that predicted by the model of Weaver *et al.* (1977). Bochkarev and Zhekov (1990) explained this low X-ray luminosity by deviation of the ionization state of the hot gas within the cavity from the equilibrium one.

Chu and Mac Low (1990), Chu *et al.*, (1993) revealed diffuse X-ray emission from seven OB/HII complexes in the LMC. Application of Weaver *et al.* (1977)

theory to some of the observed nebulae (N51D, N44) has shown that calculated X-ray luminosities fall an order of magnitude below that derived from Einstein and ROSAT observations. Chu and Mac Low (1990) considered observed X-ray emitting complexes as early stages of the superbubble evolution. They suspect that the dynamics of these bubbles is controlled by the joint action of the stellar winds from the embedded OB stars and that the X-ray excess is from the first supernova explosion near the edge of the wind-blown cavity. Chugai (1993) has proposed another mechanism to explain powerful X-ray emission from SN1986j. He has assumed that the amplification of the X-ray emission occurs due to inhomogeneity of the supernova ejecta.

In the same time analysis of the observational data has shown (de Jager *et al.*, 1988) that stellar mass-loss rate changes by more than an order of magnitude during the main sequence lifetimes of OB stars and later during the blue, red supergiant and luminous blue variable phases. The computation of the integrated output of mass, momentum, and energy from a population of massive stars (Leitherer *et al.*, 1992) has shown an obvious growth of the collective wind power with time in the continuous starburst model as well.

It is interesting to reexamine main predictions of the Weaver *et al.* (1977) model for the power-law energy input rate and to discuss possible influence of the wind power growth on the observational parameters of the wind-driven bubbles.

2 THE VARIABLE WIND MODEL

Let us consider expansion of the spherically symmetric bubble with the energy input rate

$$L = L_0 t^{\epsilon} \tag{1}$$

in a homogeneous interstellar medium with density ρ_0. After a short adiabatic phase (Avedisova, 1971; Castor *et al.*, 1975) a dense cold shell is formed with motion following from the equation of mass, momentum, and energy conservation. In the thin layer approximation (Imshennik, 1977; Bisnovatyi–Kogan and Blinnikov, 1982; Bisnovatyi–Kogan and Silich, 1995) these equations can be written as:

$$M = \frac{4}{3}\pi\rho_0 R^3, \tag{2}$$

$$\frac{d}{dt}(MU) = 4\pi R^2(P_{\rm in} - P_{\rm ext}), \tag{3}$$

$$\frac{dE_{\rm th}}{dt} = L - 4\pi R^2 P_{\rm in} U, \tag{4}$$

$$\frac{dR}{dt} = U, \tag{5}$$

where M is the mass of the swept-up interstellar gas concentrated in the thin shell after the radiative shock front, R and U are the shell radius and velocity, $E_{\rm th}$ is the thermal energy of the hot rarefied gas within the cavity, $P_{\rm in}$ and $P_{\rm ext}$ are the pressures within the cavity and in the outer interstellar matter. Neglect the outer gas pressure and put $P_{\rm ext} = 0$. Equations (1)–(5) are coupled with the expression for the internal gas pressure

$$P_{\rm in} = (\gamma - 1)\frac{3E_{\rm th}}{4\pi R^3}, \tag{6}$$

where γ is the ratio of specific heats. Substituting M and U from (2) and (5) into (3) and (4), $E_{\rm th}$ from equation (6) into equation (4) and then excluding internal gas pressure from (3) and (4), one can get an equation of the shell motion:

$$\frac{d^2}{dt^2}(R^3\dot{R}) + (3\gamma - 2)\dot{R}R^{-1}\frac{d}{dt}(R^3\dot{R}) = \frac{9(\gamma-1)L_0 t^{\varepsilon}}{4\pi\rho_0 R}. \tag{7}$$

It is easy to see from a simple substitution, that this equation has a power-law solution:

$$R = \left[\frac{125BL_0}{\rho_0}\right]^{1/5} t^{(3+\varepsilon)/5}. \tag{8}$$

where

$$B = \frac{9(\gamma-1)}{4\pi(3+\varepsilon)(7+4\varepsilon)[2+4\varepsilon+(3\gamma-2)(3+\varepsilon)]}. \tag{9}$$

The shell expansion velocity U and pressure within the cavity $P_{\rm in}$ are defined by the expression:

$$U = \frac{3+\varepsilon}{5}\frac{R}{t} = \frac{3+\varepsilon}{5}\left[\frac{125BL_0}{\rho_0}\right]^{\frac{1}{3+\varepsilon}} R^{\frac{\varepsilon-2}{3+\varepsilon}}, \tag{10}$$

$$P_{\rm in} = \frac{(3+\varepsilon)(7+4\varepsilon)\rho_0}{3\cdot 5^{\frac{2\varepsilon}{3+\varepsilon}}}\left[\frac{BL_0}{\rho_0 R^{2-\varepsilon}}\right]^{\frac{2}{3+\varepsilon}}. \tag{11}$$

A large temperature gradient between the cold shell and the hot gas within the cavity causes the shell to evaporate and add mass to the bubble interior at rate (Castor *et al.* 1975):

$$\dot{M}_{\rm ev} = \frac{16\pi}{25}\frac{\mu_{\rm in}}{k}CT_c^{5/2}R, \tag{12}$$

where $\mu_{\rm in}$ is the mean mass per particle within the cavity, $\mu_{\rm in} = \frac{14}{23}m_H$ for the completely ionized gas with 10% of helium by number ($n_{\rm He}/n_H = 0.1$), k is Boltzman's constant, $C = 6.0 \times 10^{-7}$ ergs cm^{-1} s^{-1} K$^{-7/2}$ is the coefficient of the thermal conductivity $k = CT^{5/2}, T_c$ is the temperature near the remnant centre. As shown by Castor *et al.* (1975) and Weaver *et al.* (1977), the evaporated mass reaches the self-similar distribution within the cavity with the particle number density

$$n = n_c(1-x)^{-2/5}, \tag{13}$$

where n_c is the particle number density near the bubble center, and $x = r/R$ is the current dimensionless radius. The pressure within the cavity may be considered uniform due to the high temperature and sound velocity; this defines the temperature distribution as:

$$T = T_c(1-x)^{2/5}. \tag{14}$$

Integration of the equation (13) over the bubble volume gives mass of the gas within the cavity:

$$M_{\text{in}} = \frac{125\pi}{39}\mu_{\text{in}} n_c R^3. \tag{15}$$

Taking into account that $n_c = P_{\text{in}}/kT_c$ where P_{in} is defined by equation (11) and combining derivative of equation (15) with respect to t with equation (12), we get a differential equation for T_c as a function of time.

$$\frac{1}{T_c}\frac{dT_c}{dt} + \frac{624C}{3125}\frac{T_c^{7/2}}{P_{\text{in}}R^2} - \frac{\varepsilon+1}{t} = 0, \tag{16}$$

where P_{in} and R are defined by the formulae (8) and (11). Equation (16) has a power-law solution that can be expressed as a function of the shell radius R:

$$T_c = 5^{\frac{2}{7}\left(\frac{15+2\varepsilon}{3+\varepsilon}\right)}\left[\frac{(41+27\varepsilon)(3+\varepsilon)(7+4\varepsilon)\rho_0}{13104C}\right]^{\frac{2}{7}}\left[\frac{BL_0}{\rho_0}\right]^{\frac{2}{7}\frac{3}{3+\varepsilon}} R^{\frac{2}{7}\frac{4\varepsilon-3}{3+\varepsilon}}. \tag{17}$$

The central particle number density then follows from the equation of the state and expression (11):

$$n_c = \frac{1}{3k}\left[\frac{13104C}{41+27\varepsilon}\right]^{2/7}\left[(3+\varepsilon)(7+4\varepsilon)\rho_0\right]^{5/7}\left[\frac{BL_0}{5^{\frac{15+9\varepsilon}{4}}\rho_0}\right]^{\frac{8}{7(3+\varepsilon)}} R^{\frac{2}{7}\frac{3\varepsilon-11}{3+\varepsilon}}. \tag{18}$$

Thus we have a full description of the internal bubble structure and bubble dynamics that makes it possible to calculate the evolution of the bubble X-ray luminosity:

$$L_x = 4\pi\int_0^R n_i^2(r)\Lambda_x[T(r)]r^2\,dr, \tag{19}$$

where ion number density $n_i = \frac{11}{23}n$ and gas temperature distributions $T(r)$ inside the remnant are given by the formulae (13)–(14), (17)–(18). The X-ray emissivity falls down dramatically for temperatures below the critical value $T_{\text{cr}} \approx 5\times10^5$ K. So the outer regions of the bubble interior do not add to the total bubble X-ray luminosity. Chu and Mac Low (1990) have estimated that over the temperature range 2.5×10^6–10^8 K the X-ray emissivity Λ_x of the equilibrium plasma within the Einstein energy band of 0.2—3.5 keV is almost independent of the gas temperature. It can be approximated within a 25% accuracy by the constant value

$$\Lambda \approx 9\times10^{-24}\xi \text{ergs cm}^3\ \text{s}^{-1}, \tag{20}$$

where $\xi = Z/Z_\odot$ is metallicity. Integration of (19) then gives:

$$L_x = 4\pi n_{\text{ci}}^2 R^3\Lambda_x I(\tau) = \frac{484\pi}{4761\cdot 5^{\frac{4(15+9\varepsilon)}{7(3+\varepsilon)}}}\left[\frac{13104}{41+27\varepsilon}\right]^{4/7}$$
$$\times\ [(3+\varepsilon)(7+4\varepsilon)]^{10/7}B^{\frac{16}{7(3+\varepsilon)}}\left(\frac{R}{R_0}\right)^{\frac{19+33\varepsilon}{7(3+\varepsilon)}} R_0^{-3}I(\tau)\Lambda_x, \tag{21}$$

where

$$R_0 = \left[\frac{k^{7(3+\varepsilon)}}{C^2 \rho_0^{7+5\varepsilon} L_0^8}\right]^{1/(41+27\varepsilon)}, \tag{22}$$

$$I(\tau) = \frac{125}{33} - 5\tau^{1/2} + \frac{5}{3}\tau^3 - \frac{5}{11}\tau^{11/2}, \tag{23}$$

$\tau = \frac{T_{cx}}{T_c}$, $n_{\rm ci}$ is the ion number density near the remnant centre.

3 DISCUSSION

The main observational parameters of the stellar wind bubbles are their radii. We have wrote therefore other observational quantities and bubbles parameters as functions of the independent variable R. The role of the stellar wind evolution in the bubble dynamics follows from the formulae (8)–(10). Bubble expansion velocity falls down quickly for the constant wind mechanical luminosity $\varepsilon = 0$ and becomes independent of radius for the square growth ($\varepsilon = 2$) of the input wind power (see also Ostriker and McKee, 1988). This leads, as has been emphasized by Oey and Massey (1994), to a substantial underestimation of the bubble lifetime if the constant wind power is assumed. The evolution of the central bubble temperature T_c is quite different in these two cases as well. The value of T_c reduces with bubble expansion if $\varepsilon < 0.75$ and grows with time in the opposite case. This leads to a more powerful mass evaporation rate for greater ε than follows from the equation (12). Therefore one can expect the mass of the hot bubble interior $M_{\rm ev}$ and bubble X-ray luminosity to be greater on the late stage of the bubble evolution for the increased energy input rate. Figure 1 illustrates these effects for different values of ε. Here R_1 is a fiducial radius. The coincidence of the expansion velocities, central temperatures or evaporated mass for different values of ε at $R = R_1$ is reached by the variation of the parameter L_0. For a given L_0, the bubble parameters depend on the power-law index ε and bubble radius.

Recently Oey and Massey (1994) have studied two remarkable giant (with radii $R = 66$ pc) H_α nebulae in galaxy M33. These correspond to Z362 and Z364 nebulae in the Couttes *et al.* (1987) H_α survey. They have identified these nebulae with stellar wind bubbles generated by single, isolated O stars. Using the evolutionary tracks of Maeder (1990), the spectral types and bolometric luminosities of the stars, they have roughly evaluated the star mass to be in the range 25–40$M_\odot$ for the Z362 nebula and 40–60$M_\odot$ for Z364. The best fit for the Z364 nebula corresponds to an O9 star with stellar wind power $L = 2.38 \times 10^{36}$ ergs s^{-1} at $R = 66$ pc and power-law exponent $\varepsilon \approx 1.9$. The parameter L_0 then can be found from equations (1) and (8) as follows:

$$L_0 = \left(\frac{5^3 B}{\rho_0 R^5}\right)^{\varepsilon/3} L^{(3+\varepsilon)/3}(R). \tag{24}$$

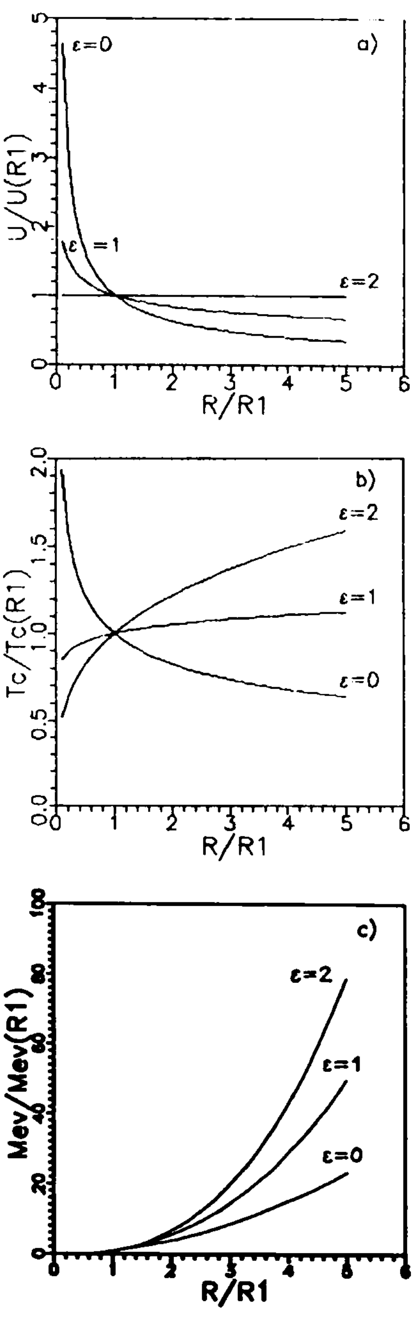

Figure 1 The wind bubble parameters as functions of the bubble radius for different wind mechanical luminosity. a, The remnant expansion velocity; b, the bubble central temperature; c, the mass evaporated from the cold shell.

It makes it possible to appreciate the X-ray evolution of the nebula. Taking as input parameters $R = 66$ pc, $L(R) = 2.38 \times 10^{36}$ ergs s^{-1}, ambient gas particle number density $n_0 = 0.3$ cm^{-3}, metallicity $\xi = Z/Z_\odot = 1/4, \gamma = 5/3$, the X-ray cut-off temperature $T_{\rm cr} = 5 \times 10^5$ K, and the X-ray emissivity $\Lambda_x = 9 \times 10^{-24}\xi$ ergs cm^3 s^{-1}, one can find parameters B and R_0 from equations (9) and (22), L_0 from the formula (24), then find bubble central temperature from the formula (17), and at last calculate the bubble X-ray luminosity (21). The results of calculations for different values of ε are shown in Figure 2. From the figure we find no evidence for observational difference in the bubble dynamics or X-ray luminosity at the current time. The expansion velocities of the shell were quite different early. The apparent difference in the X-ray luminosities may be expected for a greater bubble radius.

For a constant wind mechanical luminosity, the observed radius and expansion velocity along with external gas density and metallicity define a unique X-ray luminosity. But for a variable wind power we have an additional wind parameter, ε. The question that arises in this connection is, how strongly variability of the wind power can influence the bubble X-ray luminosity?

To answer this question, let us assume the bubble radius R and expansion velocity U to be known. Then equation (8)–(10) define the parameter L_0 as a function of ε:

$$L_0 = \left(\frac{5}{3+\varepsilon}\right)^{3+\varepsilon} \frac{\rho_0}{125B} U^{3+\varepsilon} R^{2-\varepsilon}. \tag{25}$$

The central temperature T_c and the X-ray luminosity L_x can be calculated as earlier. The action of the power-law index ε on the X-ray luminosity then follows from the ratio:

$$\frac{L_x(\varepsilon)}{L_x(0)} = \left(1 + \frac{4}{7}\varepsilon\right)^{10/7} \left(\frac{3}{3+\varepsilon}\right)^{6/7} \left(\frac{41}{41+27\varepsilon}\right)^{4/7} \frac{I(\tau_\varepsilon)}{I(\tau_0)}. \tag{26}$$

It is wonderful that power indices in this relation do not depend on ε. This predicts a weak dependence of the X-ray luminosity on the wind variation if the bubble radius and expansion velocity are fixed. This conclusion is illustrated in Figure 3, where variation of the X-ray luminosity with ε is shown for the N51D nebula parameters: $R = 53$ pc, $U = 30$ km s^{-1}, $n_0 = 0.4$ cm^{-3}, $\xi = 1/3$. The dashed lines show the same, but for different expansion velocities: $U = 10$ km s^{-1} (curve *2*) and $U = 100$ km s^{-1} (curve *3*). Thus growth of the stellar wind power cannot explain the excess of the diffuse X-ray emission from this nebula. As follows from Figure 3, the ε-dependence of the X-ray luminosity becomes stronger for the low expansion velocities. This is a consequence of the more sharp change of $I(\tau)$ due to the central temperature and volume of plasma that can emit in the X-ray band.

4 SUMMARY

1. Analytical formulae for the dynamics and X-ray luminosities of interstellar bubbles with a power-law energy input rate are found.

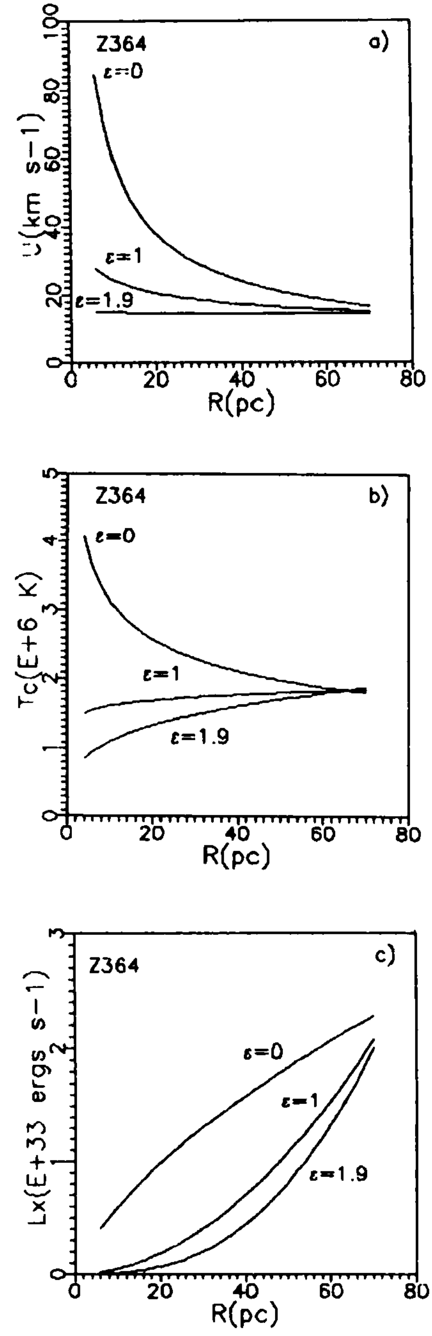

Figure 2 Evolution of the Z364 nebula for different energy input rates. a, The bubble expansion velocity; b, the bubble central temperature; c, the bubble X-ray luminosity in the Einstein energy band 0.2–3.5 keV.

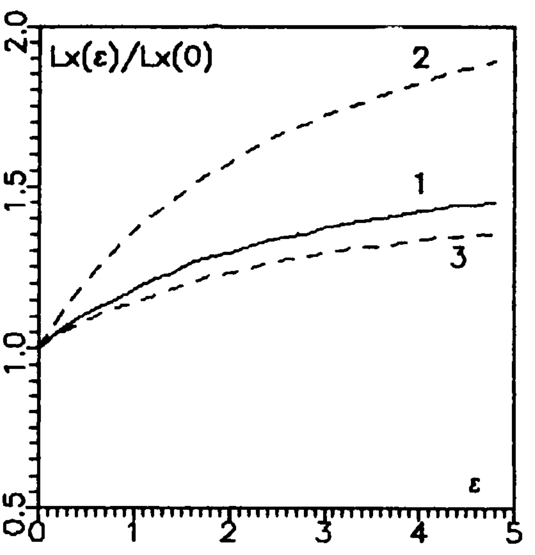

Figure 3 The X-ray luminosity as a function of the energy input rate parameter ε for the bubble-like N51D nebula. Curve *1* is for the observed expansion velocity $U = 30$ km s^{-1}, curves *2* and *3* are the same, but for $U = 10$ km s^{-1} and $U = 100$ km s^{-1}.

2. The variation of the energy input rate influences strongly the bubble expansion. However, if the bubble radius and expansion velocity are known, only a weak dependence of the bubble X-ray luminosity on the power index ε may be achieved. This dependence becomes apparent for small (near the sound speed) velocities of the bubble expansion only.

3. X-ray luminosity of the Z364 nebula in the galaxy M33 is estimated to be $\approx 2 \times 10^{33}$ ergs s^{-1} in the Einstein energy band 0.2–3.5 keV.

4. Growth of the stellar wind power with time cannot explain the excess of the diffuse X-ray emission from the N51D nebula that was discovered in the Einstein energy band by Chu and Mac Low (1990).

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