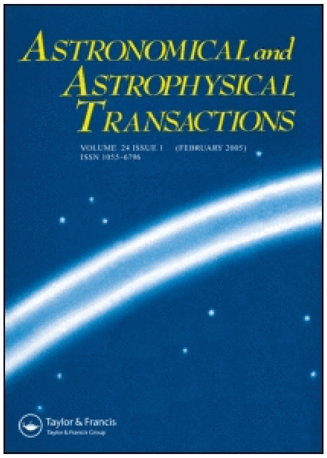

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E. Ya. Vilkoviskij [a]
[a] V. G. Fessenkov Astrophysical Institute of Kazakh Academy of Sciences, Observatory AFI, Kazakhstan

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RADIATION-ACCELERATED IONS IN THE HOT STAR WINDS AND THE NUCLEAR REACTIONS IN A NON-EQUILIBRIUM STATE

E. Ya. VILKOVISKIJ

V. G. Fessenkov Astrophysical Institute of Kazakh Academy of Sciences, 480068 Almaty, Observatory AFI, Kazakhstan

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The problem of the hot star stellar winds heated and accelerated by resonant-scattering ions is briefly discussed. We show that the "kinetic heating" mechanism is rather common in the winds of the hot stars and can explain the warm winds and the X-ray luminosities of the stars. Moreover, in some cases detectable γ-rays fluxes are expected.

1 INTRODUCTION

It is commonly accepted in the radiation-driven wind theory (Lucy and Solomon, 1970; Castor *et al.*, 1975; and others) that the velocities of resonant-scattering ions V_i are small relative to the proton thermal velocity $V_i < \bar{V}_{Tp} = (2kT/m_p)^{1/2}$ (T is the temperature and m_p is the proton mass), and a heating effect of the drift is small relative to other heating and cooling processes (radiative, in particular).

The conditions under which these assumptions are inapplicable were analyzed by Vilkoviskij (1981) (P1) and Vilkoviskij and Tambovtseva (1988) (P2), where it was shown that ions can reach much higher velocities ($V_i \gg \overline{V}_{Tp}$) in the outer regions of hot star stellar winds due to the dominance of the radiation pressure force over the proton friction force. We have shown that the heating of the outer parts of the winds by this process ("kinetic heating") can present a physical foundation for the warm wind model (Lamers and Snow, 1978).

Recently, Springmann and Pauldrach (1992) (SP) came back to the question and obtained the conclusions qualitatively similar to those in our papers, yet they made some mistakes, so we consider the problem in more detail.

Particle kinetics and runaway effects are investigated in Section 2 for both optically thin and thick cases; heat balance and the X- ray problem are considered in Section 3, whereas in Section 4 we discuss the nuclear reaction problem. The main conclusions are summarized in Section 5.

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2 PARTICLE KINETICS IN THE RADIATION-DRIVEN WINDS

2.1 The Thin Optical Depth Approximation

When the optical depth at the main (resonant) transition λ_i of the ion i (charge $z_i e$, mass m_i) is small ($\tau_i \ll 1$), the distribution function $DF(v_i)$ for the ion velocities v_i is determined by interactions with other plasma particles and with the radiation field of the stellar photosphere $I(r, \theta, \nu_i)$. The latter process can be presented as an action of the average radiative force

$$FR_i = \pi e^2/(m_e c^2) F(r, \nu_i) f_i, \tag{1}$$

where $F(r, \nu_i)$ is the radiation flux density at the resonant transition frequency and f_i is the oscillator strength. The friction force acting on the moving ion ($x_{ik} = (v_i - \overline{v}_k)/v_{Tp}$) as a result of the collisions with the particles $k(z_k, m_k)$ is

$$FP_i = 4\pi e^4 z_i^2 N_p (kT)_p^{-1} \ln \Lambda \sum_k C_k z_k^2 G(x_{ik}), \tag{2}$$

where $C_k = N_k/N_p$ is the partial density of the k-particles relative to that of protons, $G(x)$ is the Chandrasekhar function (for Maxwell's distribution functions of v_k), $x_{ik} \cong \gamma_{ik} v_i / v_{Tp}$, $v_{Tp} = (2kT_p/m_p)^{1/2}$ and $\gamma_{ik} = (T_k/T_p \cdot m_i m_k/(m_i + m_k))^{1/2}$ (in the wind comoving frame, $\overline{v}_k = 0$ for protons, He ions and electrons). The simplest approximation for $G(x)$ is $G(x) \cong x/(1 + x^3)$.

The function $FP(x_{ip})$ has two maxima: the first maximum is due to protons and helium and the second one is due to electrons. If $FR_i < FP_i^{\max}$, we have two stable (in the case $dFP_i/dv_i > 0$) and two unstable ($dFP_i/dv_i < 0$) ranges of the ion velocity v_i. If $FR_i > FP_i^{\max}$, we have an unstable ("runaway") situation with continuously accelerated ions. The criterium for the transition from the state (1) with $V_i < \bar{V}_{pT}$ and the distribution function $DF1(v_i)$ (close to the Maxwellian one) to the state (2) $V_i \gg \bar{V}_{pT}$ (with a non-Maxwellian $DF2(v_i)$) is

$$FR_i \geq FP_i(x_{ip} = 1). \tag{3}$$

Using the continuity condition for the wind matter, $dM/dt = 4\pi\rho v r^2$ ($\rho = N_p \sum C_k m_k$), we can transform (3) to the form

$$V_7 \geq 0.7 \dot{M}_6 z_i^2 / (R_1^2 T_4 \varphi(\lambda_i T_*)), \tag{4}$$

where V_7 is the wind velocity in the 10^7 cm/s units, $M_6 = (dM/dt)/(10^{-6} M_\odot/y)$, $R_1 = R_*/(10 R_\odot)$, R_* is the stellar radius. $\varphi(\lambda_i T_*) = f_i \lambda_3^{-3}/(\exp(14.4/(\lambda_3 T_4) - 1)$, $\lambda_3 = \lambda_i/(10^3 A)$; and $T_4 = T_{ef}/(10^4 K)$ (we accept the Plank photosphere spectra for the estimation).

If the ionization state of an accelerated ion is conserved in the "runaway" state, the ion would reach the terminal velocity

$$v_i(\infty) = R_* 2FR_i(r = R_*)/(R_0 m_i)^{1/2}, \tag{5}$$

where $R_0 > R_*$ is the radius at which the runaway begins.

For ζ Pup star and the CIV ion, $v_i(\infty) \sim 4 \times 10^9$ cm/s if $R_0 = 100R_*$. The ion velocity and the acceleration time t_a depend on R as

$$v_i = v_i(\infty)(1 - R_0/R)^{1/2}; \quad t_a \cong (R/v_i(\infty))(1 - R_0/R)^{1/2}. \tag{6}$$

In reality, however, the ionization state of ion is not conserved, it changes due to ionization and recombination processes, thereby changing $FR_i(\nu_i, f_i)$ and FP_i (changing z). For instance, the Li-like ions CIV, NV, OVI have strong resonant lines in the region of the $F(\nu_i)$ maximum, but their neighbouring $z+1$ ionization states have not, so the first ones can "run away" to the large velocities but they return back to the lower velocities in the $z+1$ state. As a result the behaviour of an atom (ion) in the velocity space is rather complicated and, therefore, there is a problem of finding the distribution function $DF2$, which is certainly non-Maxwellian.

The situation is more complicated due to a non-local character of the process: an ion is accelerated over the distance range from R_0 to R_1 but brakes (in the wind comoving system) at $R > R_1$ in the region with other plasma parameters. So below we shall operate with the "most probable velocity" instead of $DF2(v_i)$, which we shall determine in Section 3.

2.2 The Case of Large Optical Depth

At the Sobolev approximation (Sobolev, 1947), the optical depth in the line λ_i is

$$\tau_i \cong (\pi e^2/m_e c) N_i f_i \lambda_i/(dV/dR), \tag{7}$$

where dV/dR is the local wind velocity gradient.

The corresponding absorption probability of the $h\nu_i$ quantum is

$$\beta_i \cong (1 - \exp(-\tau_i))/\tau_i. \tag{8}$$

This does not mean at all that every ion in the local volume can absorb with the β_i probability. Rather β_i is the fraction of the ions which have the radial velocity projections exceeding the Doppler width of the absorption line. The most important mistake of SP is that they take the radiation force acting on a single ion as a mean (average) force from the acceleration $g_L^{\rm rad}$ of the optically thick stellar wind matter: $FR_i = g_i^{\rm rad} m_i \cong g_L^{\rm rad}\rho/N_i$, where is ρ the plasma density and N_i is the ion density (see SP for the definition of Γ_L, their Section 3), and as a result they hardly underestimate the force FR_i acting on the ions displaced out of the absorption line shadow. (Note that in P1 this approach was used for the estimation of a lower limit of the FR_i only, and we mean a thin-depth limit of the $g^{\rm rad}$ for the calculations).

In the real optically thick case the condition (3) is valid for a part of the ions located (in the velocity space) out of the absorption line "shadow", but it must be supplemented with the condition for the accelerated ion to be not hidden in the shadow by collisions with other particles:

$$FR_i/(m_i \nu_{ip}) \geq (2kT_p/m_p)^{1/2}, \tag{9}$$

where $1/\nu_{ip}$ is the time for the ion velocity direction change.

So when the "runaway conditions" (3) and (9) are fulfiled together in the optically thick case, the DF of the ions consists of two parts: the main part $DF1$ is close to the Maxwellian distribution, and the non-Maxwellian one $DF2$ is for runaway particles. To determine exactly the fraction $q2$ of the particles in the latter state, one should solve the the corresponding kinetic equation; as a lower limit, we accept $q2 \cong 1/\tau_i$ (where $\tau_i \gg 1$), then

$$DF_i = (1 - q2)DF1_i + q2\, DF2_i. \tag{10}$$

3 THE KINETIC HEATING AND HEAT BALANCE

With known $DF(v_i)$, the specific heat power divided by N_p^2 is

$$H = N_i/N_p^2 \int_0^\infty v_i^* DF(v_i)^* FP_i dv_i \tag{11}$$

(note that H is independent of N_p because DF and FP are $\propto N_p$).

In the optically thin case, when the condition (3) is not fulfilled ($v_i/v_{Tp} \ll 1$), the $DF(= DF1)$ is close to the "shifted Maxwellian" function $F_M(v + \overline{v}_i)$ with $\overline{v}_i \sim v_{Tp} FR_i/FP_i(x_{ip} = 1)$, and $H \cong \overline{v}_i N_i FR_i$. It reaches a maximum at $\overline{v}_i \cong v_{Tp}$,

$$H1^{\max} \cong \pi e^4 z_i^2 N_i (N_p k T_p)^{-1} \ln \Lambda^* v_{Tp}, \tag{12}$$

but at this point the condition (3) is realized and DF is transformed to $DF2$ (the non-Maxwellian one, with a high $\overline{v}_i$).

As we have mentioned above, it is difficult to determine the function $DF2$ and we shall use instead the most probable velocity $\overline{v}2_i$ in this (runaway) state. We can argue that $\overline{v}2_i$ is close to the electron thermal velocity v_{Te}:

i) The recombination time for ions is about $t^{\rm rec} \sim 10^{13} T_4^{1/2}/(N_e z_i^2)$ s; the ionization time for the most abundant ions is of the same order, so it is enough to accelerate the ions with strong lines (only those go to the runaway state) to $\geq 10^8$ cm/s in the hot star winds (see (5) and (6) for the estimations).

ii) When v_i exceeds v_{Te}, the probability of the collision ionization increases, so Li-like ions become He-like and then "switch off" the radiative force and drift to lower velocities, as it is shown above.

With the most probable velocity $v_i \sim v_{Te}$, the specific heat divided by N_p^2 is $H \cong \overline{v}_{Te} N_i z_i^2 FP_i(x_{ie} = 1)\, N_e/N_p^2 \cong N_e/N_p\, (m_p/m_e)^{1/2} H1^{\max}$. So, we can estimate

$$H2 \cong (30 \div 60) H1^{\max}. \tag{13}$$

Numerically, $H1^{\max}(N_i) \sim 2 \times 10^{-22} z_i^2 T_{p4}^{-1/2} (n_i/10^{-4})$ erg cm^3/s and $H2(N_i) \sim 10^{-20} z_i^2 T_{e4}^{-1/2} (n_i/10^{-4})$ erg cm^3/s, where $n_i = N_i/N_p$.

In the optically thick case ($\tau_i \gg 1$), $DF1$ is close to the Maxwellian one without any velocity shift (being in the shadow of the absorption) and so $H1(\tau_i) \cong 0$ and $H2(\tau_i) \cong q2H2(N_i)$ with $q2 \geq 1/\tau_i$. In fact, the $q2$ value increases with r because of the frequency of particle transitions from $DF1$ to $DF2$ exceeds the reverse process frequency due to the "nonlocality" effects.

Now we shall briefly discuss the consequences of the "kinetic heat" for the heat balance and the temperature structure of the winds.

In the "cold wind" model, the heating by ionization processes is balanced by radiative losses and the electron temperature slowly decreases with radius mainly due to the dynamic cooling. The situation changes when the "kinetic heat" is switched on. Simple estimations show that for the O- and early B-star winds the conditions (3) and (4) are satisfied mainly for the ions with $z = 1$ and $z = 2$ (having strong lines), in the range of the wind velocities $\sim 10^7 - 10^8$ cm/s. It is essential for the heat balance even if $N_i/N_p \geq 10^{-7}$ and so temperature increases to $T_e \geq 10^5$ K. Then the ions with $z = 3$ and 4 become important, including the most abundant Li-like ions.

Of course, the temperature increase depends on the star and wind parameters ($T_{\rm eff}, R_*, \dot{M}$ and V_∞), but as a rule the heat power exceeds the critical value $H_k \cong 7 \times 10^{-22}$ ergs cm^3/s (the maximum of the radiative cooling) when $V_w \leq V_\infty$ and $R_k \sim 100R_*$ for O and early B stars. Then the "temperature runaway" occurs and temperature rises up to $T_e \sim (10^6 - 10^7)$ K at $R \geq R_k$.

So the physical picture including the kinetic heat is close to the empirical "warm wind" model and, moreover, it predicts an outer hot corona at $R \geq R_k$. The X-ray luminosity of the corona is $L_x = 4\pi \int_{R_k}^{\infty} \varepsilon_X R^2 dR$. With $\dot{M} = 4\pi\rho v r^2$ and $\varepsilon_X \cong 2 \times 10^{-27} N_e^2 T_e^{1/2}$ ergs/cm^{3}s, we have

$$L_X \cong 7 \times 10^{32} T_{e7}^{1/2} \dot{M}_6^2 / (R_1 V_8^2 (r_k/100)) \text{ ergs/s}, \tag{14}$$

where $R_1 = R_*/10R_\odot, r_k = R_k/R_*, T_{e7} = T_e/(10^7 K)$ and V_8 is the wind velocity (in km/s) at R_k. (We do not include, in the estimation of the X-luminosity, the radiation from the $T_e < 10^6 K$ region and the line radiation because of a larger absorption in these cases).

We can predict a wide scatter of the X-luminosities due to the stellar parameter scatter and the X-ray flux variability (characteristic times from some hours to days) due to stellar-wind variabilities.

4 NUCLEAR REACTIONS AND THE NON-EQUILIBRIUM NUCLEAR SYNTHESIS PROBLEM

If the velocities of the accelerated ions (at the runaway state) exceed the nuclear reaction (NR) thresholds (due to the collisions with other particles moving at the wind velocities V_w), nuclear reactions occur and one can detect the resulting γ-rays.

At the distance close to R_k (somewhat less than the hot corona radius), the acceleration of Li-like ions continues and, for the most luminous hot stars, the ions reach velocities close to $V_i \sim 10^9$ cm/s (see (5) for the estimations), high enough for ion-proton and other reactions (like $^{14}N(p,\gamma)^{15}O$, emitting the 1.38 MeV γ-line). Note that the electron thermal velocity $v_{Te} \sim 2 \times 10^9 T_{e7}^{1/2}$ cm/s in the coronal region $R > R_k$ and the ionization time increases due to the density drop.

We can estimate the reaction specific rate as $\varepsilon_\gamma \sim N_i N_p v_i \sigma_\gamma E_\gamma$ and the γ-luminosity, as $L_\gamma \sim 4\pi \int_{Rk}^{\infty} \varepsilon_\gamma R^2 dR$. By analogy with (14), we have

$$L_\gamma \cong 1.6 \times 10^{28} n_{i4} v_{i9} \sigma_{27} E_1 \dot{M}_6^2 / (R_1\ V_8^2 (r_k/100))\ \text{ergs/s}, \tag{15}$$

where $n_{i4} = (N_i/N_p)/10^{-4}$; $v_{i9} = v_i/(10^9$ cm/s); $\sigma_{27} = \sigma_\gamma/(10^{-27}$ cm^2); and $E_1 = E_\gamma$ (MeV).

Note that Lamb *et al.* (1983) detected γ-lines from SS 433 with the luminosity $L_\gamma \sim 10^{37}$ergs/s, and Boyd *et al.* (1984) argued that the γ-flux has a thermonuclear reaction origin. Their Model B is close to the picture described above, and so the radiation acceleration mechanism can, in principle, explain the SS 433 case too (the radiation acceleration is in accordance with the fact that the SS 433 jet terminal velocity is determined by the hydrogen line-looking effect). In both cases we suppose the non-equilibrium nuclear synthesis (NES) (i.e. involving non-Maxwellian *DF* particles).

It should be mentioned that earlier we had discussed the NES problem before the stellar wind problems were explored (Vilkoviskij, 1970), and we proposed to use the ion acceleration in a plasma with a large electric current (Gurevich, 1961). However, it became clear that the currents required are too large, out of the TOKAMAK stability conditions.

The resonant-scattering ion acceleration is another possibility for the NES. Of course, the realization of this idea in nuclear fusion projects involves some specific technical problems and will be discussed in a separate work; here we stress that the NES problem has not only astrophysical applications.

The detection of γ-rays from SS-433 and (probably) from hot stars can show us that in solving the laboratory nuclear fusion problem we should use not only the stellar interior example (the equilibrium state), but stellar envelopes (non-equilibrium situation) as well.

5 CONCLUSIONS

i) The inclusion of the kinetic heat into the theory of stationary radiation-driven winds of hot stars predicts warm winds with outer hot coronas for the stars earlier than B3–B4 spectral classes.

ii) Nuclear reactions of the radiation-accelerated ions can generate γ-rays which can be detected in nearby most luminous hot stars.

iii) Here we consider the ideal steady winds only, but the conclusions can be valid in the outer parts of real shock-structured winds too.

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