Photosphere Recession and Luminosity of Homologous Explosions Revisited

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Abstract

By assuming the photosphere located at the outmost edge of the ejecta, Arnett et al. (1980, 1982, 1989) presented the light curves of homologous explosions in supernovae analytically and numerically to include recombination effects. Actually as homologous expansion proceeds, the photosphere recedes deeper into the ejecta. In this situation, the photospheric radius increases at early times and decreases later on, which can be described by a simple method proposed by Liu et al. To study how the photosphere recession affects the luminosity evolution, we impose a boundary condition on the photosphere to determine the spatial and time distribution of the temperature of the ejecta, which is clarified to be reasonable. We find that the photosphere recession reduces the luminosity compared with the previous result without the recession, which can be tested with observations of Type-IIP supernovae.

Unification Astronomy Thesaurus concepts: Core-collapse supernovae (304)

1. Introduction

The analytical and semianalytical light-curve models allow crude and quick estimates of explosion energy, initial radius, the mass of ejected matter, and 56Ni to be made from the observed data of supernovae (SNe). All these models assume a stationary photosphere fixed at the outmost edge of the ejecta (Arnett 1980, 1982, hereafter A80 and A82; Arnett & Fu 1989; Nagy et al. 2014b). Because of the continuous expansion of the ejecta, the photosphere radius always increases. But the photospheric radius $R_{\text{ph}}(t) = [L_{\text{bol}}(t)/4\pi\sigma T_{\text{eff}}^4(t)]^{1/2}$ of some supernovae with well observed bolometric luminosity $L_{\text{bol}}$ and effective temperature $T_{\text{eff}}$ show an early rising and late falling behavior. To explain this evolution, Liu et al. (2018) allows the photosphere to recede within the ejecta as it expands and derives an ordinary differential equation to describe this recession. Although indeed showing a general rising/falling behavior, their results cannot fit the data quite well. One possibility is the complex relation between the photosphere radius and the bolometric luminosity not as $R_{\text{ph}} \propto L_{\text{bol}}^{1/2}$. For the photosphere fixed at the edge of the envelope, its radius is easy to get. Arnett (1980, 1982) uses the Eddington surface boundary condition on the fixed photosphere to determine the spatial and time structure of the temperature below the photosphere and then obtains the luminosity. In this case, an eigenvalue $\alpha$ in separating the spatial and time structures governed by partial differential equations is found to be constant. However time-dependent $\alpha$ is necessary when the recombination effect is included (Arnett & Fu 1989; Popov 1993). The recombination front, moving through the envelope, plays a role as the pseudophotosphere. Based on Arnett & Fu (1989), Nagy et al. (2014a) added the magnetar as an extra energy input to explain the Type-IIP supernovae.

These previous works did not consider the photosphere recession effect. Actually, the photosphere will recede into the envelope as its expands in any case. In this paper, we impose the Eddington boundary condition on the receding photosphere to determine its position and then study the luminosity evolution. We assume a homologous expansion and spherically symmetric supernovae ejecta, the same as previous papers did. And we also treat the radiation transport by the diffusion approximation.

This paper is organized as follows: In Section 2 we obtain the radius evolution of the receding photosphere in a different way. In Section 3 we derive the luminosity evolution. The results are summarized in Section 4. Finally, we give a summary in Section 5.

2. Photosphere Radius Evolution

The radius of the photosphere in the ejecta is determined by

$$\int_{R_{\text{ph}}(t)}^{R(t)} \rho(r, t) \kappa dr = \frac{2}{3},$$

where $\rho(r, t)$ is the density of the ejecta, $\kappa$ is the opacity, and $R(t)$ is the surface radius of the ejecta (Arnett 1980). Because Thomson scattering dominates the opacity, we use a constant $\kappa$ throughout the evolution. In the coasting phase, the surface radius of the homologous expanding ejecta is

$$R(t) = R(0) + v_{\infty} t,$$

where $R(0)$ is the initial surface radius and $v_{\infty}$ is the velocity scale. The density can be separated into

$$\rho(r, t) = \frac{\rho(0, 0) \eta(x) R(0)^3}{R(t)^3},$$

where $0 \leq x \leq 1$ is a dimensionless radius $x = r/R(t)$, $\rho(0, 0)$ is the initial density at the innermost edge, and is written as $\rho_0$ hereafter.
From Equation (1) and Equation (3), we obtain an integration including the dimensionless photosphere radius

$$\int_{x_{ph}(t)}^{1} \rho_0 \eta(x) \frac{R(0)^3}{R(t)^2} \kappa dx = \frac{2}{3}. \quad (4)$$

As long as \( \eta(x) \) is given, the evolution of the receding photosphere can be obtained. We use a Paczyński red supergiant envelope as an example, i.e., \( \eta(x) = e^{Ax} \) with \( A = -1.732 \). Substitute it into Equation (4) and we obtain

$$x_{ph}(t) = \ln \left( \frac{e^A - (2\lambda(0) + \psi(0)^2)}{A} \right), \quad (5)$$

which is the same as the result of Liu et al. (2018).

3. Luminosity Evolution

In this section we briefly describe Arnett’s model (Arnett 1980, 1982) first. It assumed that photons come from the surface of the ejecta, which is the photosphere in their model. In fact as its expanding, the ejecta becomes thinner and the photosphere will naturally recede into the inner part of the ejecta. So next we propose a method to determine the position of the receding photosphere and to get the solution of the luminosity.

3.1. Arnett’s Model of Supernovae

In the diffusion approximation, the luminosity is

$$L/4\pi r^2 = -(\lambda c/3)\partial aT^4/\partial r, \quad (6)$$

where the mean free path is \( \lambda = 1/(\rho c) \) and the mass density is \( \rho = 1/V \). In the strictly adiabatic case, the temperature can be separated as

$$T(r, t)^4 = \frac{\psi(x) \phi(r) T(0,0)^4 R(0)^4}{R(t)^4}. \quad (7)$$

Thus the photosphere luminosity can be written as

$$L(x_{ph}, t) = -\frac{4\pi aT(0,0)^4 c R(0)}{3\rho(0,0) \kappa} \frac{\phi(r)}{\eta(x)} \left. \frac{x^2 d\psi(x)}{dx} \right|_{x=x_{ph}}, \quad (8)$$

where \( \phi(r) \) can be solved by the thermodynamics of the trapped radiation in the expansion ejecta and \( \left. \frac{x^2 d\psi(x)}{dx} \right|_{x=x_{ph}} \) is determined by the boundary condition. According to the first law of thermodynamics, the thermal state of the expanding matter evolves in time as

$$\dot{E} + P \dot{V} = -\frac{\partial L}{\partial m} + \varepsilon, \quad (9)$$

where \( E \) is the thermal energy per unit mass, \( P \) is the pressure, and \( \varepsilon \) is the energy release per unit mass from radioactive decay. For a radiation dominated gas, the energy and pressure are \( E = aT^4 V \) and \( P = aT^4 /3 \). Substitute Equation (7) and Equation (6) into Equation (9) and note that \( \partial L/\partial m = (1/4\pi r^2 \rho)\partial L/\partial r \), Equation (9) becomes

$$aT^4 \frac{\phi(t)}{\phi(r)} = \varepsilon + \frac{1}{r^2 \rho} \frac{\partial}{\partial r} \left( \frac{x^2 d\psi(x)}{dx} \right). \quad (10)$$

Now as in A80, let

$$V(r, t) = V(0, 0)[R(t)/R(0)]^3 / \eta(x), \quad (11)$$

and

$$\varepsilon = \frac{\varepsilon_0}{\kappa \eta(x)} e^{-\tau/\gamma_0}. \quad (12)$$

Using these expressions, Equation (10) reduces to

$$\frac{R(0) \dot{\phi}(t)}{R(t) \phi(t)} = \frac{\varepsilon_0}{aT(0,0)^4 V(0,0)} \left[ \frac{\psi(x) \eta(x)}{\psi(x)} \right] \right. \times e^{-\tau/\gamma_0} \frac{\phi(t)}{\phi(r)} = -\frac{\psi(0)}{3\kappa R(0)^2} \frac{\partial}{\partial x} \left( \frac{x^2 \partial \psi}{\eta(x) \partial x} \right). \quad (13)$$

As in A82, if we assume that

$$b = \frac{\eta(x) \xi(x)}{\psi(x)} \quad (14)$$

is constant for any \( x \). It’s obvious that Equation (13) is separable. Arnett (1982) defined two parameters

$$a \equiv -\frac{1}{\psi(x) x^2 \partial x} \left[ \frac{x^2 \partial \psi}{\eta(x) \partial x} \right], \quad (15)$$

and

$$\tau_0 = \frac{3\kappa R(0)^2}{a V(0, 0) c}. \quad (16)$$

Finally, the time evolution of \( \phi(t) \) is governed by Equation (13), which becomes

$$\dot{\phi}(t) + \frac{R(t)}{R(0) \tau_0} \phi(t) = \frac{b \psi_0}{aT(0,0)^4 V(0,0)} \left[ \frac{R(t)}{R(0)} \right] e^{-\tau/\gamma_0}. \quad (17)$$

This time solution and the spatial solution of Equation (15) can be solved if the density profile \( \eta(x) \) and the spatial boundary condition are given. Then the luminosity of Equation (8) will be obtained.

3.2. Boundary Conditions

As in A80, at the center we can impose the initial conditions \( \psi(x = 0) / dx = 0 \) and \( \eta(0) \equiv 1 \) to get the solution of Equation (15), which is a function of the parameter \( \alpha \). The previous literature uses a “radiative-zero” boundary condition \( \psi(1) = 0 \) to get the value of \( \alpha \). But this boundary condition is only valid for dense objects. In this paper, we consider that the photosphere recedes into the envelope as its expanding and getting thinner. So the boundary changes as it expands, which will produce a time-dependent \( \alpha \).

According to Eddington (1926), we have

$$a cT^4 = H(2 + 3\tau), \quad (18)$$

where \( H = \frac{1}{4\pi} \int J(\theta) \cos \theta d\omega \). When \( \tau \) equals to \( 2/3 \) the envelope becomes transparent for photons. So at the position of the photosphere we can obtain

$$a cT_{ph}^4 = 4H. \quad (19)$$
And thus
\[ T^4 = \frac{1}{4} T_{ph}^4 (2 + 3\tau). \] (20)

Since \( T^4 \propto \psi \), we can obtain
\[ \psi = \frac{1}{4} \psi_{ph} (2 + 3\tau). \] (21)

Take the derivative of both sides, we get
\[ \frac{\partial \psi(x)}{\partial x} = \frac{3}{4} \psi_{ph} \frac{\partial R(t)}{\partial x} = -\frac{3}{4} \psi_{ph} \frac{R(t)}{\lambda(x)}. \] (22)

Then the boundary condition of the photosphere gives
\[ \psi(x_{ph}) = -\frac{4}{3} \frac{\lambda(x_{ph})}{R(t)} \frac{\partial \psi(x)}{\partial x} \bigg|_{x = x_{ph}}. \] (23)

For a dense object, \( \lambda(x_{ph})/R(t) \ll 0 \), so \( \psi(x_{ph} \rightarrow 1) \rightarrow 0 \), which is just the “radiative-zero” boundary condition \( \psi(1) = 0 \). However, in our treatment the outside of the photosphere is no longer dense but thin enough to be optically transparent. Then the “radiative-zero” boundary condition is no longer valid. The boundary condition of Equation (23) means that the photosphere is no longer fixed but time dependent, which leads to a time-dependent \( \alpha \). We stress that all the differences come from the photosphere recession. The term \( R(0)/R(t) \) in Equation (3) makes the density thinner and thinner so that the photosphere will naturally recede. In this case, \( \psi(x) \) is also time dependent. At any moment, we always have Equation (15) and Equation (13) established. We just need to make \( \psi(x) \) satisfy the boundary condition of Equation (23) during the recession. And this method is also used in Popov (1993) and Arnett & Fu (1989). They use a time-dependent \( \alpha \) to identify the position of the moving recombinant front.

### 3.2.1. Luminosity with Fixed Boundary

To simplify the luminosity of Equation (8) further. A82 defined the initial mass of \( ^{56}\text{Ni} \) in the envelope
\[ M_{Ni}^0 = \frac{4\pi R(0)^3}{V(0, 0)} \int_0^1 \eta(x) \xi(x) x^2 dx, \] (24)

and the initial internal energy
\[ E_{Th}^0 = 4\pi R(0)^3 aT(0, 0)^4 \int_0^1 \psi(x) x^2 dx, \] (25)

where
\[ I_{Th}^0 = \int_0^1 \psi(x) x^2 dx. \] (26)

Since
\[ b = \frac{\eta(x) \xi(x) \psi(x)}{\psi'(x)} = \frac{\int_0^1 \eta(x) \xi(x) x^2 dx}{\int_0^1 \psi(x) x^2 dx}, \] (27)

we use Equation (24), Equation (25), and Equation (27) to reduce Equation (17) to
\[ \dot{\phi} = \left[ \frac{E_{Th}^0 M_{Ni}^0 a^{-\gamma_{Ni}} - \phi}{E_{Th}^0 R(0) \infty} \right] R(t) \frac{\partial R(t)}{\partial t}. \] (28)

Integrating Equation (15) by part and evaluating at \( x = 0, 1 \) gives
\[ \alpha \int_0^1 \psi(x) x^2 dx = -\left( \frac{x^2}{\eta(x) \frac{d\psi(x)}{dx}} \right)_{x = 1}. \] (29)

Using Equation (25), Equation (29), and Equation (16), the luminosity of Equation (8) reduces to a very simple expression
\[ L(1, t) = E_{Th}^0 \frac{\phi(t)}{\tau_0}, \] (30)

which can be solved analytically (see A82).

### 3.2.2. Luminosity from the Receding Photosphere

The observed emission comes from the photosphere. The first law of thermodynamics of Equation (9) within the volume enclosed by the photosphere gives
\[ \dot{E}_{ph} + (P \dot{V})_{ph} = -\frac{\partial E_{ph}}{\partial m} + \varepsilon_{ph}, \] (31)

where the internal energy enclosed by the photosphere is
\[ E_{ph} = \int_0^{R_{ph}} 4\pi x^2 T^4 dx \]
\[ = 4\pi R(0)^3 aT(0, 0)^4 \psi(0) \int_0^{R_{ph}} \psi(x) x^2 dx. \] (32)

By defining a new parameter
\[ I_{ph} = \int_0^{R_{ph}} \psi(x) x^2 dx, \] (33)

Equation (32) is rewritten as
\[ E_{ph} = E_{Th}^0 \frac{I_{ph}}{I_{Th}^0} R(0) \frac{\phi(t)}{\tau_0}. \] (34)

Therefore
\[ \dot{E}_{ph} = E_{ph} \left( \frac{\dot{\phi}(t)}{\phi(t)} - \frac{\psi_{sc}}{R(t)} + \frac{I_{ph}}{I_{Th}^0} \right) \]
\[ = E_{ph} \left( \frac{d \ln \phi(t)}{dt} - \frac{d \ln R(t)}{dt} + \frac{d \ln I_{ph}}{dt} \right). \] (35)

Using \( d \ln V/dt = 3d \ln R(t)/dt \), the pressure \( P = E/3 \), and Equation (35) to eliminate \( d \ln R(t)/dt \) from Equation (31) gives
\[ E_{ph} \left( \frac{\dot{\phi}(t)}{\phi(t)} + \frac{I_{ph}}{I_{Th}^0} \right) = -\frac{\partial L}{\partial m} \bigg|_{x_{ph}} + \varepsilon. \] (36)

Here we assume a central energy production and therefore \( \varepsilon_{ph} = \varepsilon \). According to the definition of Equation (33), the time derivative of \( I_{ph} \) is
\[ I_{ph} = \psi(x_{ph}) x_{ph}^2 \dot{x}_{ph}. \] (37)
Now with the same procedures as presented in Section 3.1, the temporal part can be obtained as follows:

\[
\frac{d\phi(t)}{dt} = \frac{R(t)}{\gamma(t)} \left[ p_1^2 \frac{I_{\text{ph}}}{I_{\text{th}}} e^{-\gamma(t)\phi(t)} - \phi(t) \right] \\
- \phi(t) \frac{\psi(x_{\text{ph}}) x_{\text{ph}}^2}{I_{\text{ph}}} \frac{dx_{\text{ph}}}{dt},
\]

where \( p_1 = \varepsilon_{\text{N}_0} M_{\text{N}_0} / E_{\text{Th}} \). Integrating Equation (15) by part and evaluating at \( x = 0 \), \( x_{\text{ph}} \) gives

\[
\left( \frac{x^2}{\eta(x)} \frac{d\psi(x)}{dx} \right)_{x=x_{\text{ph}}} = -\alpha \int_0^{x_{\text{ph}}} \psi(x) x^2 dx.
\]

The luminosity from the receding photosphere is therefore written as

\[
L(x_{\text{ph}}, t) = \frac{E_{\text{Th}}^3 I_{\text{ph}}^3}{\gamma_0^2 R_{\text{ph}}^3} \phi(t),
\]

which can only be solved numerically.

### 3.2.3. Luminosity from the Receding Photosphere: Broken Power-law Density Profile

To understand how different environments affect the recession, we also consider the broken power-law density profile, which is usually assumed as follows (Chevalier 1982; Matzner & McKee 1999; Kasen & Bildsten 2010; Moriya et al. 2013)

\[
\eta(x) = \begin{cases} 
(\frac{x}{x_0})^{-0.5} & 0 \leq x \leq x_0, \\
(\frac{x}{x_0})^{-n} & x_0 \leq x \leq 1,
\end{cases}
\]

where \( x_0 \) is a dimensionless transition radius that divides the inner part and the outer part of the envelope. For SN Ib/Ic and SN Ia progenitors, one has \( n \approx 10 \) (Matzner & McKee 1999; Moriya et al. 2013). We choose the parameters of \( x_0 = 0.1, \delta = 0 \) and \( n = 10 \), the same as in Nagy & Vinkó (2016). We can express \( x_{\text{ph}} \) analytically

\[
x_{\text{ph}}(t) = \left\{ \begin{array}{ll}
(1 - n) x_0^3 R_{\text{ph}}^3 \left( \frac{\varepsilon_{\text{N}_0} M_{\text{N}_0} (\alpha - 1) R_{\text{ph}}}{\varepsilon_{\text{N}_0} M_{\text{N}_0} (\alpha - 1) R_{\text{ph}}} \right)^{\frac{2}{\alpha - 1}}, & x_{\text{ph}} > x_0, \\
\frac{2 R_{\text{ph}}^2 (t)}{3 R_0^3 (\alpha - 1)^{\frac{2}{\alpha - 1}}}, & x_{\text{ph}} \leq x_0,
\end{array} \right.
\]

In this case, the density in the outer part of the envelope is significantly lower, which causes the temperature profile \( \psi(x) \) in that region to be very close to zero. Therefore the gradient of density and temperature is almost zero. According to Equation (23), the parameter \( \alpha \) is highly related to the temperature gradient. Thus \( \alpha \) does not change much and as a result the luminosity in our model is quite similar to the former work as shown in Figure 1.

### 3.2.4. Luminosity from the Receding Photosphere: Constant-density Profile

The constant-density profile is widely used in the research of Type-IIP supernovae (Zampieri et al. 2003; Chatzopoulos et al. 2012; Nagy et al. 2014a). We let \( \eta(x) = 1 \) to obtain

\[
x_{\text{ph}} = 1 - \frac{2 R_{\text{ph}}^2 (t)}{3 R_0^3 (\alpha - 1)^{\frac{2}{\alpha - 1}}},
\]

As shown in Figure 1, the luminosity decreases much faster with higher density in the outer region of the envelope. Such a result is understandable. In the broken power-law case, the whole envelope is like a dense core with a thin shell; the shell is almost transparent. Basically we are observing the inner part of the envelope into which the photosphere does not recede very deeply. In the exponential and constant case the temperature and density gradient is higher, thus the photosphere recession effect is stronger and causes the luminosity to be obviously lower than the fixed photosphere model.
4. Results and Discussions

We revise the luminosity evolution of homologous explosion by considering the photosphere recession. Now we compare our numerical results with the fixed photosphere model. We choose the parameters of \( R(0) = 1.59 \times 10^{14} \text{ cm}, \ M = 10 \ M_\odot, \ \kappa = 0.33 \text{ cm}^2 \text{ g}^{-1}, \ v_{sc} = 1.0 \times 10^9 \text{ cm s}^{-1}, \ E_{Th}^0 = 2.0 \times 10^{51} \text{ erg}, \) and adopt three different envelope environments. To see clearly the recession effect, we do not include the energy source of radioactive decay from \(^{56}\text{Ni}\), which just produces an exponential tail at late time. In Figure 1, the luminosity with a fixed and recession photosphere is presented. It is shown that the photosphere recession reduces the luminosity compared with the results from the photosphere fixed at the surface. We find that the density gradient is the most important factor to reduce the luminosity. Figure 3 shows the relation between eigenvalue \( \alpha \) and \( x_{ph} \). As the photosphere recedes \( \alpha \) increases with it, which leads to the decrease of photospheric radius. The consequence of such effect is what is shown in Figure 1, i.e. lower luminosity if photosphere recedes significantly.

Our luminosity formula of Equation (40) is general. The previous result of Equation (30) is just a special case. The ratio \( I_{ph}/I_{Th}^0 \) between Equation (40) and Equation (30) is equal to 1 for \( x_{ph} = 1 \), which means that our formula recovers the previous fixed photosphere result if the photosphere does not recede. We show the time evolution of the ratio for the exponential-density environment in Figure 2. The ratio reduction due to the recession is obvious. In Figure 2, we also show the numerical solution of \( \phi(t) \), which exhibits a late decrease too. The two factors combine to result in the whole behavior of the luminosity. Although photosphere recession is newly considered, we would like to stress that the idea behind introducing the time-dependent \( \alpha \) was already used in

![Figure 2. Evolution of \( \phi(t) \) and \( I_{ph}/I_{Th}^0 \) for the fixed and recession photosphere for the exponential-density environment.](image)

![Figure 3. Evolution of \( \alpha \) and \( x_{ph} \) for the fixed and recession photosphere for the exponential-density environment.](image)
Arnett & Fu (1989) and Popov (1993) to locate the position of the recombination front. While they used this method to include the recombination effect to explain the plateau of Type-IIP supernovae, the photosphere is still fixed at the surface, not at the recombination front. In the same way, we determine the position of the receding photosphere from which the luminosity radiates. The receding photosphere is a real photosphere not like the recombination front. We can furthermore conjecture that the photosphere should recede even as the recombination of hydrogen takes place, which is beyond the scope of this paper.

The decrease of luminosity caused by the photosphere recession is significant at late times as shown in Figure 1. In this period, the luminosity due to the nebula or the magnetar comes to exceed the photosphere emission (Wang et al. 2016). Therefore it is difficult to observe the recession directly. However, we find that the behavior of the luminosity considering the recession in about the initial 100 days after the burst is quite different from previous results if the recombination effect is included. The recombination of hydrogen highly affects the light curves of Type-II supernovae, especially for Type-IIP supernovae.

In summary, we solve the luminosity of the homologous explosion in supernovae considering the photosphere recession for the first time. Fitting supernovae and other optical transient data in our future work (in preparation) will be used to find the evidence of the photosphere recession.

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