Magnetically powered outbursts from white dwarf mergers

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ABSTRACT

Merger of a white dwarf binary creates a differentially rotating object which is expected to generate strong magnetic fields. Kinetic energy stored in differential rotation is partially dissipated in the magnetically dominated corona, which forms a hot variable outflow with ejection velocity comparable to $10^9$ cm s$^{-1}$. The outflow should carry significant mass and energy for hours to days, creating an expanding fireball with the following features. (i) The fireball is initially opaque and its internal energy is dominated by the trapped thermal radiation. The stored heat is partially converted to kinetic energy of the flow (through adiabatic cooling) and partially radiated away. (ii) Internal shocks develop in the fireball and increase its radiative output. (iii) A significant fraction of the emitted energy is in the optical band. As a result, a bright optical transient with luminosity $L \sim 10^{41} - 10^{42}$ erg s$^{-1}$ and a characteristic peak duration comparable to 1 day may be expected from the merger. In contrast to classical novae or supernovae, the transient does not involve nuclear energy. The decay after its peak reflects the damping of differential rotation in the merger remnant. Such outbursts may be detected in the local Universe with current and upcoming optical surveys.

Key words: binaries: close — magnetic fields — MHD — radiation mechanisms: general — stars: coronae, rotation, winds, outflows — supernovae: general — white dwarfs

1 INTRODUCTION

Population synthesis models suggest the birth rate of binary white dwarfs (WD) in our Galaxy comparable to $0.05$ yr$^{-1}$, and a large fraction of these binaries, perhaps a half, are expected to merge in less than a Hubble time (Nelemans et al. 2001). The possibility of detecting gravitational waves could make the WD mergers particularly interesting sources for future observations. About 30 tight WD binaries have already been found (Kilic et al. 2012); however, no mergers events have yet been identified. It was suggested that some of them may be associated with thermonuclear supernovae (SN Ia, Iben & Tutukov 1984; Weblink 1984). The estimated rate of WD mergers is comparable to that of SN Ia (Badenes & Maoz 2012), supporting their candidacy for SN Ia progenitors. Many of the mergers are, however, unable to ignite a thermonuclear explosion — the possibility of ignition depends on the mass ratio of the binary and its chemical composition (Dan et al. 2012).

In this paper, we discuss WD mergers that do not explode or explode with a significant delay (Raskin et al. 2009). Using simple estimates, we argue that even without the liberation of nuclear energy the mergers should eject bright fireballs detectable in optical surveys.

1.1 Post-merger object

The merger forms an axisymmetric, rapidly rotating central core surrounded by a less massive debris disc (e.g. Benz et al. 1990; Mochkovitch & Livio 1990; Rasio & Shapiro 1995; Guerrero et al. 2004; Yoon et al. 2007; Lorèn-Aguilar et al. 2009; Raskin et al. 2012). The energy budget of this nascent object is

$$E_0 = \frac{GM^2}{R} \sim 3 \times 10^{50} \left( \frac{M}{M_\odot} \right) \left( \frac{R}{10^9 \text{ cm}} \right)^{-1} \text{ erg},$$

where $M \sim M_\odot$ and $R \sim 10^9$ cm are the characteristic mass and radius of the merger remnant. The remnant has three important features:

1. Fast rotation. The orbital angular momentum of the binary system is inherited by the post-merger object. Its characteristic angular velocity $\Omega$ is comparable to the maximum (break-up) angular velocity $\Omega_{\text{max}} = (GM/R^3)^{1/2}$. The outer parts of the core and the surrounding disc have significant differential rotation, i.e. $\Omega$ varies with cylindrical radius by $\Delta \Omega \sim \Omega$.

2. The deep interior of the merger core can remain...
degenerate, however its upper layers and the surrounding debris disc are strongly heated by shocks (e.g. Raskin et al. 2012), which lifts the electron degeneracy and creates significant thermal pressure. A large contribution to the pressure is made by radiation, $P_{\text{rad}} = \alpha T^4/3$, where $T$ is the temperature and $\alpha$ is the radiation constant. Convection is likely to develop in the upper layers of the core and the debris disc.

(3) The remnant is expected to be strongly magnetized. Even if the pre-merger magnetic fields are weak, the differential rotation and convection quickly generate strong fields. This process was recently studied by Ji et al. (2013) whose magneto-hydrodynamic simulations show fields up to $10^{10} - 10^{11}$ G.

The generated field is convenient to express using the “magnetization parameter” — the ratio of magnetic energy $E_B \sim B^2 R^3$ (omitting a numerical factor $\sim 0.1$) to the total energy of the object $E_0 \sim GM^2/R$, $\epsilon_B = B^2 R^4 / GM^2$. (2)

In this paper we focus on the low-density corona formed around the merger, and hereafter $\epsilon_B$ corresponds to the magnetic field in the corona.

1.2 Analogies with other astrophysical differential rotators

One close analogy is provided by neutron-star (NS) mergers. They received significant attention as major sources of gravitational waves; they are also thought to produce short gamma-ray bursts (GRBs, e.g. Piran 2004). A strongly magnetized corona and outflows are believed to develop in NS mergers. Although they differ from WD mergers in many respects — e.g. they have a smaller size (by a factor of $\sim 10^3$) and a much higher cooling rate due to neutrino emission, — basic physics of magnetic field amplification is similar. Both WD and NS mergers create an object that can be described as an excellent conductor with fast differential rotation and convection; the Rossby numbers for the two cases are comparable. Similar dynamo processes are also expected in proto-neutron stars formed in stellar collapse. Differential rotation generates a strong toroidal magnetic field, which is buoyant and forms a magnetically dominated corona (e.g. Ruderman et al. 2000; Spruit 2008). Unlike proto-neutron stars, the WD merger is not cooled by neutrino emission and does not become solid; the lifetime of its differential rotation and coronal activity is controlled by the effective viscosity due to magnetic stresses and turbulent diffusion.

Rich observational data are available for another class of fast magnetized rotators — pre-main-sequence stars (protostars). They have a similar, fast-rotating core surrounded by a Keplerian disk. These objects are significantly bigger in size than WD mergers, by a factor of $\sim 10^2$. Their characteristic age is $\sim 10^7$ yr and their inferred magnetization is modest. Nevertheless, the observed rotation-powered activity of protostars (e.g. Montmerle et al. 2000; Getman et al. 2008) can provide a useful analogy. They show a strong coronal activity — X-ray flares triggered by reconnection events. Their average X-ray luminosity is a factor $10^3 - 10^5$ larger than that of the sun (e.g. Feigelson et al. 2007). More generally, accretion discs are observed in protostars, X-ray binaries, and quasars. These canonical differential rotators are known to produce magnetized jets and strong nonthermal emission.

Amplification of magnetic fields is also suspected in stellar mergers, one of which has been caught by recent observations and gave rise to the red nova V1309 Scorpii (Tylenda et al. 2011). It provides a possible mechanism for formation of magnetic Ap/Bp stars (Soker & Tylenda 2007; Ferrario et al. 2009; Tutukov & Fedorova 2010). The WD merger is intermediate between the stellar mergers and neutron-star mergers in terms of size, active lifetime and energy budget of the remnant.

2 ACTIVE CORONA

2.1 Formation

The merger remnant maintains hydrostatic equilibrium on the sound crossing timescale,

$$t_0 = \frac{R}{c_s} \approx \left( \frac{R^3}{GM} \right)^{1/2} = \frac{\Omega_{\text{max}}}{\alpha} \sim 3 \text{ s}.$$ (3)

Its density profile is controlled by the distribution of entropy generated by shocks in the merger and subsequent viscous dissipation. The low-density upper layers are expected to have higher entropy per unit mass, and supported mainly by radiation pressure. The object is in differential rotation and will tend to redistribute its angular momentum on a viscous timescale $t_{\text{visc}} \gg 10^3$ s (Shen et al. 2012; Ji et al. 2013). This timescale depends on the value of viscosity created by magnetic fields (and turbulence) in the remnant; viscous stress may be parameterized as $T_{\text{visc}} = \alpha P$ where $P \lesssim GM^2/R^4$ is the characteristic pressure. The timescale $t_{\text{visc}} \sim 10^3$ s corresponds to $\alpha \sim 0.1$. In Section 3 below we mainly focus on early times $t < t_{\text{visc}}$, when differential rotation is still strong. At this stage, magnetic fields are amplified and buoyantly emerge from the remnant on a timescale that is intermediate between $t_0$ and $t_{\text{visc}}$. An active magnetically dominated corona must be sustained around the differential rotator.

Strong magnetic fields should be generated in the merger debris disc as well as in the upper layers of the central core. The disc has the initial scale-height $H/r \sim 0.1$ (Raskin et al. 2012). Its heating and further evolution develops on the viscous timescale $t_{\text{visc}} \sim \alpha^{-1}(H/r)^{-2} \Omega^{-1}(r) \sim 10^4$ s. The resulting super-Eddington accretion disc becomes thick and prone to outflow formation. The net energy released on the viscous timescale is comparable to the kinetic energy of the rotating matter. This matter, which was initially gravitationally bound, will remain bound if the generated heat is distributed strictly in proportion to mass density, giving a sound speed $c_s \sim v_K = (GM/r)^{1/2}$. Then the disc can evolve into a bound, radiation pressure-supported, quasi-spherical structure, growing in size and slowing down its rotation (Shen et al. 2012). However, the uniform heating is unlikely. Magnetic fields generated by the magneto-rotational instability are buoyant and expected to deliver and dissipate energy in the upper layers of lower density. As a result, a fraction $f$ of the disc mass is heated to $c_s \gtrsim v_K$ and becomes

$^1$ The average entropy per unit mass in WD mergers corresponds to $P_{\text{rad}} \sim P_{\text{gas}}$. Upper layers with higher entropy have $P_{\text{rad}} \gg P_{\text{gas}}$. 

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gravitationally unbound. For instance, $f \sim 10^{-2}$ will give an outflow of total mass $\sim 10^{-3} M_\odot$ and energy $\sim 10^{48}$ erg.

Magnetic activity of the post-merger object is demonstrated by recent numerical simulations of Ji et al. (2013). They find that the magnetic energy of the remnant at its peak exceeds $10^{48}$ erg (which corresponds to a space-average $B \sim 10^{11}$ G) and a significant mass $M \sim 10^{-3} M_\odot$ is ejected from the system over the run time of their simulations, $t = 2 \times 10^4$ s.

Hereafter “corona” refers to the magnetically dominated region of radius $r \sim R$ coupled to the central core or the disc. We will assume that the magnetization parameter of the corona, as defined in Equation (2), satisfies $\epsilon_B > 10^{-5}$ which approximately corresponds to the magnetic field $B > 2 \times 10^9$ G.

Unlike the solar corona, the coronae of WD mergers are opaque to radiation. The photosphere lies at a large radius in the outflow zone which will be described below.

2.2 Dissipated energy

Shearing of the footprints of coronal magnetic field lines by differential rotation and convective motions in the remnant repeatedly twists the field lines to the threshold of instability, leading to magnetic flares and field-line opening similar to coronal mass ejection in solar flares. An upper limit for the dissipation rate in the corona is given by $E_B \Omega$, where $E_B \sim 0.1 \epsilon_B G M^2 / R$ is the magnetic energy of the corona,

$$L_c < L_{\text{max}} \sim 0.1 \epsilon_B G M^2 / R \Omega \sim 10^{49} \epsilon_B \text{ erg s}^{-1}. \quad (4)$$

The corona may be approximately described as a force-free magnetic configuration with dissipation localized in current sheets. Numerical simulations of dissipation are challenging, with results depending on the artificial (numerical) resistivity. Simulations by Ji et al. (2013) suggest strong coronal dissipation during $t_{\text{visc}} \sim 10^4$ s, with the average $L_c \sim 10^{44}$ erg s$^{-1}$.

The magnetic energy around the remnant is continually pumped by the shearing motions of the field-line footprints frozen in the differential rotator, and a quasi-steady balance between pumping and dissipation (flares) is maintained. As a result, a fraction of the rotational energy stored in the merger, $E_0 \sim GM^2 / R$, is gradually lost through the corona and outflow. Assuming that the most active phase of this process ends together with strong differential rotation at $t \sim t_{\text{visc}}$, the lost energy may be written as

$$E \sim L_c t_{\text{visc}}. \quad (5)$$

For numerical estimates we will adopt $E \sim 10^{48}$ erg, similar to the results of Ji et al. (2013); note that it is less than one per cent of the merger energy $E_0$ given by Equation (1).

The outflow energy is carried by matter and magnetic fields; the typical ejection velocity $v$ is a few times the escape velocity $v_0$ (see Section 2.4 below), and the characteristic ejected mass is

$$M \sim 10^{-3} M_\odot \epsilon_{\text{th}} v_0^2. \quad (6)$$

2.3 Density and temperature

The ejection of mass $M \sim 10^{-3} M_\odot$ during $t_{\text{visc}} \sim 10^4$ s corresponds to the average mass loss rate $\dot{M} \sim 2 \times 10^{26}$ g s$^{-1}$.

Let $t_c$ be the residence time of gas in the corona before it is ejected; $t_c$ cannot be shorter than the sound crossing time $t_0 = R/v_0 \sim 3$ s or the timescale for field-line twisting $\Delta \Omega^{-1} \gtrsim t_0$. The characteristic mass density of the corona is related to $t_c$ by

$$\rho \sim \frac{\dot{M} t_c}{4 \pi R} \sim 1 \frac{M_{26}}{10^3} \left(\frac{t_c}{10^4 \text{ s}}\right) \text{ g cm}^{-3}. \quad (7)$$

It is much smaller than the interior density of the merger remnant $\rho \sim M/R^3 \sim 10^6$ g cm$^{-3}$.

The hydrostatic scale-height of the corona $H$ is comparable to the remnant radius $R \sim 10^9$ cm. The plasma confined in the magnetic field is free to expand along the field lines, and the condition $H \sim R$ corresponds to the sound speed $c_s \sim (P/\rho)^{1/2}$ being comparable to the virial velocity,

$$c_s \sim v_0 = \left(\frac{GM}{R}\right)^{1/2} \sim 4 \times 10^8 \text{ cm s}^{-1}. \quad (8)$$

This implies a relation between the characteristic mass density $\rho$ and pressure $P = U/3$ in the corona: $P \sim \rho c_s^2$. The pressure is dominated by radiation, $P = a T^4 / 3$, which gives

$$T \approx \left(\frac{3 \rho v_0^2}{a}\right)^{1/4} \sim 10^8 \rho^{1/4} \text{ K}, \quad (9)$$

where $\rho$ is in units of g cm$^{-3}$.

Using these estimates one can verify that the corona is completely opaque to radiation, and its thermal energy density and pressure are dominated by blackbody photons. Gas pressure $P_{\text{gas}} \approx p kT/m_p$ (where $m_p$ is proton mass) contributes a small fraction to the total pressure $P$,

$$P_{\text{gas}} \approx \frac{p kT}{m_p} \sim \frac{kT}{m_p \rho v_0^2} \sim 5 \times 10^{-2} \left(\frac{T}{10^8 \text{ K}}\right). \quad (10)$$

The thermal energy density of the corona $U = 3P$ is supplied by dissipation of magnetic energy, and may be written as

$$U = \epsilon_{\text{th}} \frac{B^2}{8 \pi} \sim \epsilon_{\text{th}} \epsilon_B \frac{GM^2}{R^3}, \quad (11)$$

where $\epsilon_{\text{th}} < 1$. In a local dissipative region (a flare generated by an unstable current sheet), $\epsilon_{\text{th}}$ is not much below unity; the average value of $\epsilon_{\text{th}}$ is much smaller. Note that $\rho \sim 1$ g cm$^{-3}$ and $U \sim \rho v_0^2$ roughly corresponds to $\epsilon_{\text{th}} \epsilon_B \sim 10^{-6}$.

2.4 Ejection velocity

The outflow is gravitationally unbound and its minimum expected velocity is comparable to the virial velocity $v_0$. The outflow can be additionally accelerated by the magnetic Lorentz force (equivalent to centrifugal acceleration if viewed in the frame co-rotating with the open magnetic field lines), as described by the standard theory of magnetized winds (e.g. Lamers & Cassinelli 1999). The acceleration is significant if the magnetic flux penetrating the outflow $\Psi_{\text{op}}$ (open flux) is sufficiently strong to enforce the flow co-rotation with angular velocity $\Omega$ to a large radius $R_A > R$. Then the flow can be ejected with the velocity $v \sim v_0 + \Omega R_A$. Here $R_A$ is the Alfvén radius at which $B^2 \sim 4 \pi \rho v_0^2$; this condition can be rewritten using the open magnetic flux $\Psi_{\text{op}} \sim Br_A^2$ and mass flux $\dot{M} \sim \rho v_0 r_A^2$ at $r \sim R_A$,

$$\frac{\Psi_{\text{op}}^2}{R_A} \sim 4 \pi \frac{M}{R_A} v_A. \quad (12)$$
Substitution of $R_A \sim v_A/\Omega$ gives the equation for $v_A$, and one finds that the outflow is ejected with
\[ v \sim \left( \frac{\Omega^2 \Psi_{op}^2}{4\pi M} \right)^{1/3} \approx 10^9 \Omega^{2/3} \Psi_{op,28}^{2/3} M_{26}^{-1/3} \text{ cm s}^{-1}. \] (13)

This standard estimate assumes $v_0 < v < c$. If $v$ approaches $c$, the relativistic wind theory should be used (Michel 1969), however this does not occur in our fiducial model with $M \sim 10^{26}$ g s$^{-1}$. The active corona can only open a fraction of the total magnetic flux of the post-merger object, $\Psi_{t,a}$, which is a small fraction of the total angular momentum $J$. The condition $v > v_0$ requires $\Psi_{op} > 10^{27} M_{26}^{1/2} / G$ cm$^2$. In the numerical estimates below we adopt the outflow speed $v \sim 10^9$ cm s$^{-1}$, a few times larger than $v_0$. Ji et al. (2013) find a comparable $v \sim 2v_0$.

2.5 Spindown effect of the outflow

The model outlined above assumes that the energy stored in differential rotation is mostly dissipated inside the remnant, and roughly one per cent or less is dissipated in the corona and feeds the outflow. The outflow also carries away angular momentum,
\[ J \sim M R_A v \sim M \frac{v}{\Omega}, \] (14)
which is a small fraction of the total angular momentum of the object, $J \sim M R^2 \Omega$.

\[ \frac{J}{J} \sim \frac{M}{M} \left( \frac{v}{v_0} \right)^2 \left( \frac{\Omega}{\Omega_{\text{max}}} \right)^{-2} \ll 1. \] (15)

This ratio is $\sim 10^{-2}$ for the typical parameters of our model.

The remnant can slowly lose its angular momentum to a weaker wind at much longer times $t \gg t_{\text{visc}}$, when the mass loss rate is reduced. Note that even with conserved angular momentum the angular velocity of the remnant can decrease on the timescale $t_{\text{visc}}$ as it viscously spreads into a bigger object (Shen et al. 2012).

3 RADIATION FROM THE FIREBALL

Consider a fireball of mass $M \sim 10^{-3} M_{\odot} \approx 2 \times 10^{30}$ g ejected during time $T \sim 10^5 - 10^6$ s with velocity $v \sim 10^9$ cm s$^{-1}$. Its heat is comparable to its kinetic energy, $E \sim M v^2 / 2 \sim 10^{48}$ erg. The heat can be lost through radiative diffusion and adiabatic cooling; on the other hand, it can be generated by delayed internal dissipation, in particular by internal shocks in the fireball. These processes and the resulting emission are discussed below.

3.1 Diffusion radius and photospheric radius

Radiation tends to diffuse out of the fireball. In time $t$, radiation diffuses through the characteristic length defined by $l_D^2 = D t$, where $D = (3k\rho)^{-1}c$ and $k$ is the opacity. As a reference value for $k$ one can use Thomson opacity $\kappa_T \approx 0.2$ cm$^2$ g$^{-1}$; the actual $k$ will be discussed below. The fireball density at radius $r = vt$ may be written as
\[ \rho = \frac{\dot{M}}{4\pi r^2 v} \sim \frac{M}{4\pi r^2 v T}. \] (16)

where we assumed that the fireball is quasi-spherical (it is straightforward to extend the model to beamed outflows). This gives the diffusion length at radius $r$,
\[ l_D = \left( \frac{4\pi^3 c}{3k M} \right)^{1/2}. \] (17)

When $l_D$ approaches the fireball thickness $l_f = \min \{vT, r\}$ radiation is no longer trapped. The condition $l_D = l_f$ defines the characteristic “diffusion radius” $R_D$ where radiation escapes the fireball. This radius is given by
\[ R_D = \left\{ \begin{array}{ll}
\frac{3k T \dot{M} v^3}{4\pi c} \; & vT \rightarrow \frac{T}{R_D}, \; T < T_1, \\
\frac{3k M}{4\pi c} = vT_1 \left( \frac{T}{T_1} \right)^{-1}, \; & T > T_1
\end{array} \right. \] (18)

$R_D$ is maximum if $T = T_1$,
\[ T_1 = \left( \frac{3k M}{4\pi c v} \right)^{1/2} \approx 6 \times 10^4 \left( \frac{k}{k_T} \right)^{1/2} M_{-3}^{1/2} v_9^{-1/2} \text{ s}, \] (19)
where $M_{-3} = M / 10^{-3} M_{\odot}$ and $v_9 = v / 10^9$ cm s$^{-1}$. A typical diffusion radius is $R_D \sim 3 \times 10^{13}$ cm.

At a larger radius $R_\star$, the fireball becomes transparent to radiation,
\[ R_\star \approx \left( \frac{k M}{4\pi} \right)^{1/2} \approx 2 \times 10^{14} \left( \frac{k}{k_T} \right)^{1/2} M_{-3}^{1/2} \text{ cm}, \] (20)
where we assumed $T < R_\star/v$. At late times $t > R_\star/v$ the photosphere shrinks; its evolution is controlled by the decreasing $M$: $R_\star \approx \kappa M / 4\pi v$.

In the above numerical estimates for $R_D$ and $R_\star$, we normalized the opacity $k$ to its Thomson value $k_T \approx 0.2$ cm$^2$ g$^{-1}$. The actual Rosseland mean opacity is comparable to $k_T$ at radii $r \lesssim R_D$ and reduced at $r \sim 10^{14}$ cm (Section 3.4), which slightly reduces the photospheric radius.

3.2 Effective temperature and optical luminosity

For a given luminosity $L$ emitted by the fireball at a radius $r$, the effective temperature of radiation is given by
\[ T = \left( \frac{L}{4\pi r^2 \sigma} \right)^{1/4} \approx 1.9 \times 10^4 L_{42}^{1/4} r_{14}^{-1/2} \text{ K}, \] (21)
where $\sigma$ is the Stefan-Boltzmann constant. Assuming a quasi-thermal radiation spectrum near the photosphere, one can see that it will peak near the optical band if $L \sim 10^{43} - 10^{44}$ erg s$^{-1}$. At smaller $r$ or higher $L$, one finds $kT > h\nu$ for the optical frequency $\nu \sim 6 \times 10^{14}$ Hz. Then the luminosity in the optical band $L_\text{O}$ can be significantly smaller than the total emitted luminosity $L$. Approximating the emission at $\nu < kT$ by the Raleigh-Jeans formula, one finds
\[ L_\text{O} \sim 8\pi^2 r^2 v^3 kT \approx 5 \times 10^{41} L_{42}^{1/4} r_{14}^{3/2} \text{ erg s}^{-1}. \] (22)

3.3 Adiabatic cooling and internal shock heating

In the absence of delayed internal dissipation, the expanding opaque fireball cools adiabatically with adiabatic index $\gamma = 4/3$. Its energy density is decreasing as $U \propto r^{-8/3}$, which corresponds to $T \propto r^{-2/3}$. As the fireball expands
from $r \sim R$ to the diffusion radius $R_D$, its volume is increased as $(r/R_D)^2$ and the total thermal energy is reduced as $(r/R_D)^{-2/3}$. Thus, the total energy emitted to distant observers is

$$\mathcal{E}_{\text{em}} \sim \left( \frac{R_D}{R} \right)^{2/3} E \sim 10^{-3} \mathcal{E}. \quad (23)$$

It is emitted on a timescale $t \sim R_D/v$ if this time is longer than $\tau_{\text{r}}$, which gives luminosity

$$L \sim \left( \frac{R_D}{R} \right)^{-2/3} \frac{v}{R_D} \sim 10^{39} - 10^{41} \text{ erg s}^{-1}. \quad (24)$$

A more detailed estimate takes into account the additional cooling that occurs when the fireball is magnetically accelerated to $v > v_0$; then there is an additional cooling factor $(v/v_0)^{-4/3}$.

The luminosity emitted by the fireball at earlier times, when its radius $r < R_D$, may be estimated as

$$L \sim \frac{4\pi^2 L_D U}{t}, \quad r < R_D. \quad (25)$$

One then finds that the bolometric luminosity is slightly higher at earlier times, $L \propto t^{-1/6}$. Luminosity in the optical band is smaller at $t < R_D/v$; it reaches its peak at $t \sim R_D/v$. The rise of the optical luminosity toward the peak can be estimated using Equation (22) and $L^{1/4} \propto t^{-1/24} \propto \text{const}$, which gives $L_0 \propto t^{3/4}$.

These estimates assume passive adiabatic cooling of the expanding fireball. In reality, it is likely to experience internal heating in a broad range of radii. The outflow is created by the variable corona of the differential rotator, and its ejection velocity $v$ can vary by $\Delta v \sim v$ on a broad range of timescales $\Delta t_{\text{var}}$, from $R/v_0$ to the age of the remnant.

As the outflow cools, its variable velocity profile leads to internal supersonic motions and formation of shocks at radius $r \sim v^2 \Delta t_{\text{var}}/\Delta v$. Then part of the energy lost to adiabatic cooling is converted back to heat. The shock-heated plasma again adiabatically cools, and new shocks can form. The shock heating can continue to large radii, even approaching the photospheric radius $R_*$. Note the possibility of a gradual increase of the ejection velocity on the long timescale $\tau \sim t_{\text{var}}$, as at late times the mass loading of the outflow may be reduced, and the Alfvén radius may increase, leading to a stronger centrifugal acceleration of the flow. As the faster parts of the fireball catch up with the earlier ejected slower part, a strong shock develops at a radius $r \sim v\tau \sim 10^{13} v_0 T_{\text{b}} \text{ cm}$. The development of shocks can be affected by the magnetic field carried by the fireball. The field is transverse to the outflow velocity at radii $r > R_D$ and radial waves propagate in the plasma with the fast magnetosonic speed $v_m \approx \delta v \lesssim v$, where $v_m = B/(4\pi \rho)^{1/2}$. Shock dissipation may be suppressed if variations in the ejection velocity are smaller than $v_m$, which can be satisfied when a large fraction of the fireball energy is carried by the magnetic field. The field itself can, however, dissipate through magnetic reconnection in the outflow (Drenkhahn & Spruit 2002), providing an alternative source of heat.

Internal dissipation at radii $r \gg R$ significantly increases the expected luminosity. The adiabatic cooling factor at the diffusion radius $(R_D/R)^{-2/3} \sim 10^{-3}$ is offset by heating and should be replaced by a larger factor $\sim 10^{-2} - 10^{-1}$. The dissipative fireball can easily emit $L \sim 10^{42} \text{ erg s}^{-1}$ at $R_D$ and possibly a comparable luminosity at $R_*$. The luminosity generated at the photospheric radius $R_*$ may also be estimated as

$$L_* \sim \frac{\eta_* \mathcal{E}}{\tau_r}, \quad (26)$$

where $\eta_*$ is the dissipation efficiency at $r \sim R_*$ and $\tau_r \sim R_*/v$ is the characteristic timescale of photospheric emission. This gives,

$$L_* \approx \eta^2 \frac{M v^2}{2 R_*} \sim \pi^{1/2} \eta^2 \left( \frac{M}{\kappa_*} \right)^{1/2} v^3$$

$$\approx 6 \times 10^{42} \eta \left( \frac{R_*}{\kappa T} \right)^{-1/2} M_4^{1/2} v_0^3 \text{ erg s}^{-1}. \quad (27)$$

A moderate dissipation efficiency $\eta_*$ of a few per cent provides a high photospheric luminosity $L_* \sim 10^{41} \text{ erg s}^{-1}$.

3.4 Detailed models with accurate opacity

The above estimates were scaled to Thomson opacity $\kappa_T \approx 0.2 \text{ cm}^2 \text{ g}^{-1}$. The actual opacity of the outflowing plasma is a function of temperature $T$ and density $\rho$. This dependence can be found e.g. in the OPAL tables (Iglesias & Rogers 1996) for helium or carbon-oxygen composition. Using the OPAL tables, we have calculated several simple models of the outburst with accurate opacities. The models assumed the following mass loss rate of the remnant,

$$\dot{M} = 10^{26} \text{ g s}^{-1} \times \begin{cases} 1 & t < T \tau \quad \text{g s}^{-1} \times \end{cases} \quad (28)$$

At radii $r < R_D$, radiation carried by the flow is described by $L = 4\pi^2 U v$, where $U = aT^4$. It is approximated by

$$L \approx \frac{\dot{M} v^2}{2} \frac{v_0}{r_0}. \quad r < R_D. \quad (29)$$

Outside the diffusion radius $R_D$, radiation density $U$ is quickly reduced so that luminosity $L \approx 4\pi^2 U v_{\text{diff}} \approx \text{const}$, where $v_{\text{diff}}$ increases from $v$ at $R_D$ to $\sim c$ at the photosphere $R_*$. We calculated the models with outflow velocity $v = 10^5 \text{ cm s}^{-1}$, $T = 10^4 - 10^5 \text{ s}$, and studied two cases for $\delta$: $\delta = 2/3$ (simple adiabatic cooling, no internal dissipation) and $\delta = 1/3$ (toy model where cooling is slowed down due to internal dissipation). Accurate opacities $\kappa(\rho, T)$ were used in the calculations, for two chemical compositions: pure helium or carbon-oxygen (with the carbon mass fraction $X_C = 0.4$). The results confirmed that $\kappa$ is close to $\kappa_T$ (within $\sim 50$ per cent) until the fireball reaches $r \sim (2 - 6) \times 10^{13}$ cm where its temperature decreases to $\sim 10^4$ K; then the opacity drops. The diffusion radius $R_D \lesssim 10^{14} \text{ cm}$ and photospheric radius $R_* \approx 10^{14} \text{ cm}$ are close to the estimates in Section 3.1. Note that all absorptive processes — free-free, bound-free, and bound-bound — contribute to the fireball opacity in the main emission region $10^{13} \lesssim r \lesssim 10^{14} \text{ cm}. This should help the thermalization of escaping radiation.
3.5 Decay of the outburst
At times $t > T$ the mass loss rate of the merger remnant $\dot{M}$ decreases and the outburst decays. At late times, the diffusion radius $R_D$ and the photospheric radius $R_\star$ are both reduced proportionally to $\dot{M}$. The bolometric luminosity scales as $L \propto \dot{M} R^2 \eta$. This gives the effective temperature of escaping radiation $T \propto \dot{M}^{-3/4} v^{1/2} \eta^{1/4}$, i.e. the temperature grows if the efficiency $\eta$ and the flow velocity $v$ remain approximately constant.

Then Equation (22) gives the optical luminosity,

$$L_\text{O} \propto \dot{M}^{7/8} v^{1/2} \eta^{-1/4}.$$  (30)

The optical luminosity is sensitive to the mass loss rate of the remnant and should drop when $\dot{M}$ is reduced.

3.6 Nonthermal emission
Internal shocks dissipate the velocity variations $\Delta v \lesssim v$ and heat the plasma to temperature $kT \sim 10^3$ keV. The plasma immediately converts its heat to radiation behind the shock (via inverse Compton scattering, bremsstrahlung, and line emission). Thus, an X-ray luminosity $L_X$ up to $\sim 10^{42} \text{erg s}^{-1}$ is generated inside the fireball. As long as the fireball is opaque to X-rays, $L_X$ is re-processed into quasi-blackbody radiation. Late internal shocks propagating outside the X-ray photosphere could produce observable X-ray emission. Internal shocks may also generate non-thermal particles, which produce high-energy (inverse Compton) photons and synchrotron radiation with a broad spectrum.

At early times the remnant age is smaller than the radiation escape time, $t < R_D/v$, so internal shocks only occur in the highly opaque zone. The longest timescale of the central engine variability $t_{\text{var}}$ is comparable to its age, and the growing age of the remnant helps formation of internal shocks at large radii. At the same time, the photospheric radius of the outflow is reduced $\propto \dot{M}$ as the mass loss rate decreases. Thus, propagation of shocks in the transparent zone becomes more likely at late times, after the peak of the outburst. Then highly variable nonthermal emission may be detected in addition to the quasi-thermal component.

4 OUTFLOW FROM A QUIET CORONA
This section briefly discusses the mass loss rate that could be expected from a “quiet” corona, i.e. in the absence of flares due to field-line twisting by footpoint motion. The corona may become relatively quiet as the merger remnant ages past $t_{\text{visc}}$. Shen et al. (2012) argue that the remnant should evolve into a quasi-spherical object in approximately solid-body rotation (although we do not exclude that the aging remnant is still surrounded by a low-mass Keplerian disc). The viscously heated remnant expands, and its maximum (break-up) angular velocity decreases as $\Omega_{\text{max}} \propto R^{-3/2}$ while the actual angular velocity decreases as $\Omega \propto R^{-2}$ (as long as its angular momentum is approximately conserved).

Even in the absence of differential rotation and coronal flares, the rotating magnetized remnant would lose mass along open magnetic field lines. Some of the field lines must be open by rotation. In the force-free approximation, the minimum open magnetic flux is $\Psi \sim \mu/R_{\text{LC}}$, where $R_{\text{LC}} = c/\Omega$ is the light cylinder radius and $\mu$ is the magnetic dipole moment of the remnant. Then less than one percent of magnetic field lines are open. More field lines can become open if significant mass is lifted from the remnant and its corona is changed from the force-free configuration (Mestel & Spruit 1987).

4.1 Mass flux
Consider the quiet corona supported by radiation pressure in a fixed strong magnetic field. The field enforces co-rotation of the plasma with angular velocity $\Omega$. To demonstrate the reason of mass loss and estimate its magnitude consider a field-line bundle near the equatorial plane in the simplest, monopole-like geometry (radial field lines). The plasma at the base of the corona must be close to hydrostatic equilibrium. The hydrostatic balance for the radiation-dominated plasma reads

$$\frac{\kappa F}{c} = g(r) = \frac{GM}{r^2} - \frac{\Omega^2}{\kappa} r,$$  (31)

where $F$ is the radiation flux along the magnetic field line and $\kappa \approx \kappa T = \text{const}$ is the plasma opacity. The hydrostatic balance implies continual heating of the plasma,

$$\dot{Q} = -\frac{1}{r^2} \frac{\partial}{\partial r} (r^2 F) = 3\frac{\kappa c}{\Omega}.$$  (32)

Hence, a steady state is possible only if the plasma is gradually flowing away from the object, which modifies the hydrostatic picture and introduces advection of heat (dominated by radiation) by the opaque flow.

The hydrostatic approximation (31) is still useful where the flow velocity is subsonic. As the sound speed $c_s$ of the heated flow increases, it climbs the effective (gravitational − centrifugal) potential barrier, and its hydrostatic scale-height grows,

$$H(r) = \frac{c_s^2(r)}{g(r)}.$$  (33)

At sufficiently large altitudes, the heated flow passes through the sonic point and escapes with velocity $v \sim c_s \approx v_0$.

However, to estimate the mass loss rate one can consider the deep, subsonic, approximately hydrostatic region where $v \ll c_s \ll v_0$ and $H \ll R$, $r \approx R$. The heating timescale $\sim U/Q$ is comparable to the timescale for the change in $H$,

$$\frac{3pc}{Q} \sim \frac{H}{v}.$$  (34)

Equations (32)-(34) give the mass flux in the outflow,

$$F_m = \rho v \sim \frac{\dot{Q}}{3g} = \frac{c \Omega^2}{\kappa g} = \frac{c}{\kappa R} \left(\frac{\Omega_{\text{max}}^2}{\Omega^2} - 1\right)^{-1}.$$  (35)

The net mass loss rate is $\dot{M} \sim F_m A$, where $A$ is the area of the footprint of the open field-line bundle on the remnant. This mass loss rate is much smaller than that of the active, flaring corona discussed in Sections 2 and 3.

4.2 Strong field regime
In a sufficiently strong magnetic field, the outflow remains magnetically dominated and continues to accelerate centrifugally at $r \gg R$, reaching the light cylinder with $v \sim c$. © 0000 RAS, MNRAS 000, 000-000
In this regime, most of the energy lost by the rotating object is carried away by the Poynting flux and described by the standard pulsar spindown formula,
\[ \dot{E}_P \sim \frac{\mu^2 \Omega^4}{c^3} \sim 4 \times 10^{38} \frac{\mu_3^2}{c^2} \Omega^{-1} \text{erg s}^{-1} , \]
where \( \mu \sim BR^3 \) is the magnetic dipole moment of the remnant. The kinetic power of the matter ejected at the light cylinder is \( \dot{E}_{\text{mat}} \sim Mc^2 \); more energy is transferred from the Poynting flux to the plasma outside \( R_{\text{LC}} \) where \( \dot{E}_{\text{mat}} \) grows to its asymptotic value (Michel 1969),
\[ \dot{E}_{\text{mat}} \sim \sigma \frac{1}{3} M c^2 , \quad \sigma = \frac{\dot{E}_P}{Mc^2} > 1 . \]
Estimating \( A/4\pi R^2 \sim R/R_{\text{LC}} \) (which would roughly correspond to a dipole field), one obtains
\[ \dot{E}_{\text{mat}} \sim \frac{L_E \sigma^{1/3} c^2}{R_{\text{LC}} r_0^2 / \tau_0} \sim \frac{L_E \sigma^{1/3} R^2}{\tau g R_{\text{LC}}} \sim 3 \times 10^{39} \sigma^{1/3} R_0^3 \Omega^{-1} \text{erg s}^{-1} , \]
where \( \tau_g = GM/c^2 \sim 10^5 \) cm and \( L_E = 4\pi GMc/\kappa \approx 2.5 \times 10^{38} (M/M_\odot) \) erg s\(^{-1}\) is the Eddington luminosity with \( \kappa = \kappa_T \). Comparing Equations (36) and (37), one can see that the condition \( \sigma > 1 \) (the strong-field regime) is satisfied when \( \mu_3^2 \Omega_3^3 > 10 \).

Radiation is trapped and advected by the outflow below the sonic radius \( r_s \sim R \). At larger radii, radiation diffusion becomes faster than advection; here radiation diffuses through the outflow and escapes at its photosphere, which is located well inside of the light cylinder (roughly at \( r \sim 10^{10} \) cm). A moderate quasi-thermal luminosity is emitted from the object, comparable to the Eddington luminosity \( L_E \sim 10^{38} \) erg s\(^{-1}\), with the effective temperature comparable to \( 10^5 \) K. The kinetic and magnetic power of the outflow in the strong-field regime significantly exceeds the quasi-thermal photospheric luminosity. It may be partially dissipated and converted to radiation at large radii, producing nonthermal radiation with a broad spectrum.

### 4.3 Weak field regime

The regime \( \dot{E}_P > \dot{E}_{\text{mat}} \) discussed in Section 4.2 is satisfied if the surface magnetic field \( B \sim \mu R^{-3} \) exceeds the characteristic value
\[ B_1 \sim \left( \frac{2\pi F_m c^4}{R^4 \Omega^2} \right)^{1/2} \sim \left( \frac{2\pi c^5}{\kappa \gamma R^4 \Omega^2} \right)^{1/2} \sim 3 \times 10^{10} \frac{R_3^3}{\Omega_3^{3/2}} \text{G} . \]

If the surface field is weaker, the Alfvén radius \( R_A \) (at which \( B^2 \sim 4\pi \rho c^2 \)) becomes smaller than \( R_{\text{LC}} \), the open field-line bundle becomes broader (its footprint area \( A \) is increased) and the asymptotic outflow velocity is reduced below \( c \).

The calculation of \( F_m \) in Section 4.1 is valid if the surface magnetic field is strong enough to enforce co-rotation of the coronal plasma. This condition requires a minimum field,
\[ B_0 \sim \left( 4\pi F_m v_0 \right)^{1/2} \sim 10^{6} \frac{\omega}{(1 - \omega^2)^{1/2}} R_9^{3/2} \text{G} , \]
where \( \omega = \Omega/\Omega_{\text{max}} \). If \( B > B_0 \) then \( R_A \sim R, v_A \sim v_0 \), and a large fraction of magnetic field lines are open by the outflow, \( A \sim R^2 \). Then \( \dot{M} \sim R^2 F_m \sim cR/\kappa T \) and \( \dot{E}_{\text{mat}} \sim M_4 v_6^2/2 \).

The scaling of \( R_A, v_A, A, \dot{M} \) and \( \dot{E}_{\text{mat}} \) with \( B \) in the range of \( B_0 < B < B_1 \) may be estimated by comparing their values at \( B \sim B_0 \) and \( B \sim B_1 \). This gives a rough estimate, \( \dot{M} \sim B^{-1/2} \), \( E_{\text{mat}} \sim B^{1/2} \), \( (B_0 < B < B_1) \).

The maximum mass loss of the quiet corona, \( \dot{M} \sim cR/\kappa T \sim 10^{20} \text{g s}^{-1} \) is approached when \( B \sim B_0 \), with a modest energy output \( \dot{E}_{\text{mat}} \sim \dot{E}_P \sim 10^{38} \) erg s\(^{-1}\).

### 5 DISCUSSION

The active corona of the differentially rotating remnant produces an outflow with the velocity \( v \sim 10^9 \) cm s\(^{-1}\). The expected duration of the high mass-loss phase is \( T \gtrsim 10^4 \) s, comparable to the lifetime of strong differential rotation in the object. The outflow carries away a fraction of the energy stored in differential rotation. This fraction is uncertain. Numerical simulations of Ji et al. (2013) and simple estimates suggest that it can be comparable to one per cent, which corresponds to the ejected mass \( M \sim 10^{-3} M_\odot \).

In this paper we discussed the consequences of this mass ejection.

The outflow creates a dense fireball of characteristic thickness \( vT \sim 10^{13} \) cm, which becomes transparent to radiation at the photospheric radius \( R_\ast \sim 10^{14} \) cm. It is unlikely to be spherically symmetric, and may contain a faster jet near the rotation axis. The initial thermal energy of the fireball is dominated by radiation and comparable to its kinetic energy \( E \sim Mv^2/2 \); most of it is lost to adiabatic cooling. In the absence of dissipative processes in the fireball, only a small fraction \( \sim 10^{-3} \) of \( E \) is radiated away. Even this small fraction would give an interesting transient event with luminosity \( L \sim 10^{40} \) erg s\(^{-1}\).

We further argued that the outflow from the active corona must be highly variable in a broad range of timescales. The variable outflow is expected to develop internal shocks (in addition to possible delayed magnetic dissipation), which convert a fraction of the fireball kinetic energy to heat, offsetting adiabatic cooling. The dissipative fireball is expected to produce a higher luminosity \( L \sim 10^{41} \) erg s\(^{-1}\), comparable to the luminosities of core-collapse supernovae. The peak timescale of the produced outburst is between \( t_{\text{visc}} \) and \( R_\ast/v \), comparable to 1 day.

The emitted radiation has the effective temperature \( T \gtrsim 10^4 \) K, and a significant fraction of the outburst is emitted in the optical band. Estimates in this paper used the simplifying blackbody assumption up to the photospheric radius. A more realistic emission spectrum could be obtained with detailed transfer calculations.

The fireball emission is \( \sim 10^5 \) times brighter than classical novae, and may be called “kilonova”, similar to the transients expected from neutron-star mergers (e.g. Metzger et al. 2010; Kasen et al. 2013). The classical novae are emitted by ejecta of a smaller mass \( M \sim 10^{-5} \) \( 10^{-6} M_\odot \) moving with velocity \( v \sim 10^8 \) cm s\(^{-1}\). Their photospheric radii approach \( R_\ast \sim 10^{17} \) cm when the ejecta temperature decreases.

3 Results of more accurate calculations would depend on details of the magnetic configuration, in particular, the angle between the magnetic dipole moment and the rotation axis of the object, and the presence of multipoles.
to $\sim 10^4$ K; then the ejecta opacity quickly decreases, and the photosphere recedes (e.g. Gallagher & Starrfield 1978). Similar behavior, but with a larger $R_\ast \sim 10^{14}$ cm and higher $L$, may be expected for the outbursts from the WD mergers. We emphasize that the proposed outburst mechanism does not invoke nuclear reactions, in contrast to classical novae, supernovae or kilonovae from NS mergers. The main energy source is differential rotation that is partially converted to heat around the remnant through magnetic dissipation.

A special feature of this scenario is that the decay of the central engine activity occurs on a timescale comparable to the radiation diffusion time $R_D/v \sim 3 \times 10^3$ s. This coincidence implies that the observed light curve can be affected by both the fireball expansion and the central engine evolution. The long timescale variability of the central engine also makes it possible for internal dissipation to occur in an extended range of radii not much below the main emission zone $r \sim 10^{13} - 10^{14}$ cm. A similar mechanism is unlikely to work for the mildly relativistic outflows from NS mergers, which are ejected on a timescale $< 10^2$ s, much shorter than the emission time $t \sim 10^2$ s.

The outbursts with luminosities $L \sim 10^{41} - 10^{42}$ erg s$^{-1}$ (absolute visual magnitudes of $\sim 13$ to $\sim 16$) can easily be detected in optical surveys with sufficiently short cadence $\sim 1$ day. Adapting that the observed volume (and the expected number of detections) scales as $L^{3/2}$ and the rate of WD mergers is comparable to that of thermonuclear supernovae (Badenes & Maoz 2012), one can roughly estimate the expected number of detections as $N \sim (L/L_{SN})^{3/2} N_{SN}$, where $N_{SN}$ is the number of detected thermonuclear supernovae and $L_{SN} \sim 10^{43}$ erg s$^{-1}$ is their luminosity. The current survey by the Palomar Transient Factory may detect the outbursts, and the upcoming Large Synoptic Survey Telescope should routinely observe them.

Magnetically powered outbursts from WD mergers should differ from normal supernovae in several respects. (i) The optical light curve should peak early (one day timescale) and then show an unusual decay that reflects the decay of differential rotation in the remnant. (ii) The effective temperature of emission decreases as the optical luminosity rises towards its peak, reaches the minimum $T_{\text{min}} \approx 10^4$ K near the peak, and is expected to grow while the source is fading. (iii) Line features in the spectrum should differ from those in supernovae, as the fireball is dominated by unburned material, with chemical composition close to that of the merging WDs (dominated by carbon, oxygen or helium). (iv) Heating by internal shocks and delayed magnetic dissipation might extend to large radii where the fireball becomes transparent. Then a variable nonthermal component may be emitted with a broad spectrum, from radio to gamma-rays (Section 3.6).

Several transients with luminosities of $10^{41}$--$10^{42}$ erg s$^{-1}$, fast decay and puzzling chemical composition have recently been detected; they were interpreted as unusual, rare variations of supernovae, although their origin is not established (see e.g. Kleiser & Kasen 2013 and refs. therein). Magnetically powered fireballs from WD mergers should produce similar events, contributing to the diversity of observed transients.

The decay of differential rotation must significantly change the observational appearance of the merger remnant at late times $t \gg t_{\text{dec}}$. The corona becomes less active and the mass loss should significantly decrease (Section 4). The strong magnetic fields generated by differential rotation, $B \sim 10^{10} - 10^{11}$ G, are expected to decay. The surviving fields may still be relatively strong, $B \sim 10^8 - 10^9$ G, providing a possible formation scenario for magnetic white dwarfs (García-Berro et al. 2012; Külebi et al. 2013). The remnant can temporally increase in size, due to viscous heating, and correspondingly slow down its rotation (Shen et al. 2012); then the remnant will cool down on the long Kelvin-Helmholtz timescale and shrink.

The ejected fireball will eventually be decelerated by the surrounding medium. The characteristic deceleration radius $R_{\text{dec}}$ is where the swept-up external mass becomes comparable to the ejecta mass $M$,

$$R_{\text{dec}} \sim 7 \times 10^{37} M_*^{1/3} n^{-1/3} \text{ cm},$$

where $n$ is the number density of the external medium in units of cm$^{-3}$. The ejected mass $M$ reaches $R_{\text{dec}}$ in time $t_{\text{dec}} \sim R_{\text{dec}}/v \sim 20 v_0^{-1} M_*^{1/3} n^{-1/3} \text{ yr}$. This interaction is accompanied by a strong shock wave, which is expected to produce nonthermal particles and synchrotron radiation. Significant radio emission may be produced at this stage.

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