A Second Decoupling Between Merging Binary Black Holes and the Inner Disc–Impact on the Electromagnetic Counterpart

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ABSTRACT

The coalescence of two supermassive black holes (SMBHs) produces powerful gravitational-wave (GW) radiation and, if gas is present in the vicinity, also an electromagnetic (EM) counterpart. In the standard picture, an EM outburst will be produced when the binary “decouples” from the circumbinary disc and starts “squeezing” the disc inside the secondary orbit, resulting in its quick accretion on to the primary black hole. Here we use analytical arguments and numerical simulations to show that the disc within about 20 $R_S$ of a SMBH survives the merger without being depleted. The reason is a “second decoupling”: the inner disc thickens due to tidal heating and inefficient cooling, effectively decoupling from the interaction of the binary. We show that this second decoupling quenches the heating sources in the disc within 10$^2$ days before coalescence. This will render the peak UV/X-ray luminosity significantly weaker than previously thought. After the merger, the residual disc cools down and expands, merging with the outer disc rather than being completely accreted. This results in continuous EM emission, hindering the detection of the cut-off and re-brightening proposed in earlier studies.

Key words: accretion, accretion discs – methods: analytical – methods: numerical – black hole physics – hydrodynamics – gravitational waves

1 INTRODUCTION

Supermassive black hole binaries (SMBHBs) form in galactic nuclei following hierarchical galaxy mergers (Begelman et al. 1980; Volonteri et al. 2003). Depending on the conditions of the surrounding stellar and gaseous environment, some SMBHBs efficiently lose orbital energy and angular momentum, managing to coalesce within a Hubble time (see Merritt & Milosavljević 2005; Colpi & Dotti 2009, for reviews). These mergers produce powerful gravitational-wave (GW) radiation whose frequency falls in the sensitive bands of the ongoing Pulsar Timing Array project1 and the planned Laser Interferometer Space Antenna2.

Electromagnetic (EM) radiation could be produced simultaneously with GW emission if a merger happens in an gas-rich environment (see Centrella et al. 2010; Schnittman 2013, for reviews). Such “EM counterparts” contain rich information about the distribution and evolution of matter around supermassive black holes (SMBHs). This information, combined with those extracted from GW signals, would deliver a more comprehensive picture in the strong-gravity regime. Such a picture will greatly enrich our knowledge of black hole physics and astrophysics.

A particular place that can produce EM counterparts is active galactic nucleus (AGN). SMBHBs in AGNs are likely embedded in gaseous accretion discs. Theoretical models predict that after a binary shrinks to a separation of $a \sim 10^2$ Schwarzschild radius ($R_S$) it will “decouple” from an “outer disc”–the part of the disc outside the binary orbit– because the merger time-scale of the SMBHs due to GW radiation becomes shorter than the viscous time-scale of that disc (Armitage & Natarajan 2002). More importantly for this work, since the GW radiation time-scale diminishes quickly as the binary separation shrinks, it will become shorter than the viscous time-scale of the “inner disc”–the disc enclosed by the binary orbit. The consequence is a strong tidal “squeezing” of the inner disc by the coalescing binary (Armitage & Natarajan 2002).

It is generally accepted that the squeezing would heat up the inner disc and produce a luminous “precursor”–an enhancement of the EM radiation– days to hours before the SMBHB merger (Lodato et al. 2009; Chang et al. 2010; Tazzari & Lodato 2015; Cerioli et al. 2016). It is also suggested that this process could drive the material of the inner disc either inward into the SMBHs or outside the binary orbit, so that eventually there is no material near the post-merger SMBH (Armitage & Natarajan 2002). As a result, AGN activity is halted (such as jet formation Liu et al. 2003; Liu 2004) and no UV/X-ray radiation could be detected until the outer disc either refills the cavity (Milosavljević & Phinney 2005; Dotti et al. 2006; Chang et al. 2010; Tanaka et al. 2010; Shapiro 2010) or gets shock-heated either by the sudden change in the gravitational potential (Schnittman & Krolik 2008; Megevand et al. 2009; O’Neill et al. 2009) or the recoil of the central SMBH (Lippai et al. 2009).

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1 http://www.ipta4gw.org

2 https://www.elisascience.org
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2008; Shields & Bonning 2008; Schnittman & Krolik 2008; Rossi et al. 2010).

However, an important point has been often overlooked in the previous studies: the squeezing mechanism is effective only when the inner disc is geometrically thin. If the disc becomes thick, fluid elements can cross the binary orbit and leak to the outer disc, either due to a large effective viscosity (Lin & Papaloizou 1979; Papaloizou & Lin 1984) or through horseshoe trajectories (Baruteau et al. 2012). Armitage & Natarajan (2002) envisaged that the inner disc inevitably becomes geometrically thick based on the observation that in their numerical simulations when \( a \lesssim 10 R_\odot \), the accretion rate exceeded the Eddington limit. Tazzari & Lodato (2015) also mentioned the effect of tidal heating on the thickness of the disc, but they did not include it in their models.

In this letter we use both analytical arguments and numerical simulations to prove that at a relatively large binary separation, \( a \gtrsim 20 R_\odot \), the inner disc always becomes geometrically thick. We also discuss the impact on the detectability of the EM counterparts.

2 HEATING SOURCES AND DISC THICKNESS

The thickness of an accretion disc can be characterized by the scale height \( h \) (or “half thickness”) and is normally supported by gas pressure \( p_{\text{gas}} \) and radiation pressure \( p_{\text{rad}} \). We are interested in the case when radiation pressure dominates since we are considering a disc that is at a distance of \( r \lesssim 10^7 R_\odot \) from the SMBH. In this case, the condition for vertical hydrostatic equilibrium reduces to

\[
h = \frac{\kappa F}{\Omega(2\Sigma^2)},
\]

where \( F \) is the radiation flux at the disc surface, \( \Omega \) the Keplerian angular velocity, \( \kappa \) the speed of light and \( \Sigma \) the opacity.

To quantify \( h \), we first consider a standard thin disc where radiative cooling is balanced by viscous heating. This condition leads to the equation of energy equilibrium (Frank et al. 2002)

\[
F = \dot{D}_s = 9\Sigma \Sigma v^2 / 8,
\]

where \( \dot{D}_s \) is the viscous dissipation rate per unit surface area, \( v \) the viscosity and \( \Sigma \) the surface density of the disc. Assuming a stationary disc, we can substitute \( v \Sigma \) with a constant fraction of the mass accretion rate, \( M/(3\pi) \) (Frank et al. 2002). From Equations (1) and (2) we find that the scale height has a constant value:

\[
h_0 = 3k \Sigma \Omega / (8\pi) \approx 7.\dot{m}R_\odot,
\]

depending only on \( \dot{m} \), the accretion rate normalized by the Eddington rate \( \dot{M}_\text{Edd} \), assuming a mass-to-radiation coefficient of 0.1.

The above analysis indicates that when the accretion rate of a disc approaches the Eddington limit, i.e. \( \dot{m} \sim 1 \), the inner part of it, e.g. \( r \lesssim 10 R_\odot \), will have an aspect ratio of \( h_0 / r \sim 1 \) and the disc will become thick. The same conclusion applies to the squeezing phase, only that the accretion rate and the heating rate will be determined by the tidal interaction, not viscosity.

The squeezing phase starts when the GW radiation time-scale of a SMBHB, \( t_{\text{GW}}(a) \), becomes shorter than the viscous time-scale of the inner disc, \( t_{\text{vis}} \). For a circular binary, \( t_{\text{GW}}(a) = 5a^3 / [8c R_\odot^2 q(1 + q)] \) (Peters 1964), where \( q = M_i / M_p \) is the mass ratio of the binary, \( M_p \) is the mass of the primary (bigger) black hole, \( M_i \) is the mass of the secondary one. For simplicity, we only consider the disc surrounding the primary SMBH, but our conclusions can also be applied to the disc around the secondary.

We can, without loss of generality, consider the fluid elements of a thin annulus between the radii \( r \) and \( r + \Delta r \), where \( \Delta r \ll r \) is the width of the annulus. When \( t_{GW} \ll t_{vis} \), the squeezing mechanism will force the radius of this ring to shrink on a time-scale similar to the shrinking time-scale of the SMBHB. Therefore, if \( v \) is the radial velocity of the annulus, we have \( |v| \lesssim t_{GW} \). To satisfy this relationship, the fluid elements in the ring must lose their orbital energy (by shocks, Lin & Papaloizou 1986b) at a rate of \( (\pi \Sigma \Delta r)(GM_p/2r) \), with \( G \) the gravitational constant. This dissipation will heat up the disc at a rate, per unit surface area, of

\[
D_\lambda \sim -GM_p \Sigma v^2 / (4r^2).
\]

This means that in this phase we have two heating terms in the energy equation: \( \dot{D}_s \) due to the tidal force and \( \dot{D}_\lambda \) due to viscosity. Taking into account the fact that \( t_{\text{vis}} = 2r^3 / (3v) \), we can write

\[
\dot{D}_s / \dot{D}_\lambda \sim t_{\text{vis}} / (3t_{\text{GW}}).
\]

This ratio is independent of the detailed structure of the disc and is valid as long as the fluid elements of the inner disc do not cross the orbit of the SMBHB. Since we already know that \( t_{\text{GW}}(a) \ll t_{\text{vis}}(r) \) during the squeezing phase, it becomes clear that \( D_\lambda \gg D_s \).

Therefore, the appropriate equation for thermal equilibrium is

\[
F = \dot{D}_s + \dot{D}_\lambda \sim D_\lambda.
\]

From this we derive that

\[
h_0 \sim 2.5\dot{m}R_\odot,
\]

where \( \dot{m} = 2\pi \Sigma \Delta M / M_\text{Edd} \). Again we find that the inner part of an accretion disc becomes thick when \( \dot{m} \sim 1 \).

Although the dependence of \( h \) on \( \dot{m} \) is the same (linear) for both viscosity- and tidally-dominated discs, the accretion rate \( \dot{m} \) entails very different physics in these two cases. For a standard disc, \( \dot{m} \) is determined by viscosity and hence proportional to \( D_s \). In the case of a squeezed disc, \( \dot{m} \) is driven by the tidal force so it scales with \( D_\lambda \). If one mistakenly uses \( D_s \) to calculate energy dissipation during the squeezing phase, one would significantly underestimate the accretion rate as well as the scale height, wrongly considering the disc to be thin.

For this reason, earlier works that neglected \( D_\lambda \) (Armitage & Natarajan 2002; Baruteau et al. 2012; Tazzari & Lodato 2015; Cerioli et al. 2016) inevitably have underestimated \( \dot{m} \). Lodato et al. (2009) and Chang et al. (2010) included \( D_\lambda \) in their energy equations. However, they do not appear to have paid attention to the aspect ratio, and hence overlooked that \( h/r \) will be \( O(1) \) during the squeezing phase. In the following sections we will calculate \( h/r \) and show that it gets close to unity for \( a \sim 20 R_\odot \).

3 ANALYTICAL MODEL

To find out at which binary separation is the condition \( h/r \sim 1 \) satisfied, we replace \( v \) in Equation (4) with \( -r/t_{\text{GW}} \), where \( t_{\text{GW}} \) is a function of \( a \). Then from Equation (1) and (3) we derive

\[
h/r = \sqrt{t_{\text{GW}} / (2t_{\text{GW}})},
\]

where \( t_{\text{cool}} = \tau h/c \) is the cooling time-scale due to radiation and \( \tau = \Sigma v / 2 \) the optical depth of the disc. Equation (7) indicates that the thickness of the inner disc closely correlates with its ability to cool. Only when \( t_{\text{cool}} \ll t_{\text{GW}} \) is the disc geometrically thin.

To proceed, we need to express \( t_{\text{cool}} \) as a function of \( r \) as well. This requires knowledge of \( \Sigma \). Although we do not know yet the surface density during the squeezing phase (this will be calculated in the next section), we notice that it should be greater than the surface density of an unperturbed standard accretion disc, because
the squeezing process generally increases $\Sigma$ (see e.g. Figure 3 of Armitage & Natarajan 2002). This lower limit is

$$\Sigma_0 = 1.6 \times 10^5 \alpha^{-4/5} m^{3/5} M_t^{1/5} r_0^{3/5} \text{g cm}^{-2},$$  \hspace{1cm} (8)

(adapted from Kocsis et al. 2012), where $r_0 = r/(10^2 R_*)$, $M_t = M_\odot/(10^5 M_\odot)$ and $\alpha$ is the standard viscosity parameter (Shakura & Sunyaev 1973, also see Section 4 for more details). Combining this $\Sigma_0$ and the scale height $h_0$ derived earlier, we find a lower limit $t_{cool} = \frac{\alpha \Sigma h_0}{(2c_s) = 0.25 \alpha^{-3/5} m^{8/5} M_t^{1/5} r_0^{3/5} \text{yr}},$  \hspace{1cm} (9)

for the cooling time-scale.

From Equation (7) and the fact that $t_{cool} \gg t_{cool0}$, we derive

$$h/r > 0.025 \alpha^{-2/5} m^{8/5} M_t^{1/5} r_0^{-3/5} \alpha_2^3 (q(1 + q))^{1/2},$$  \hspace{1cm} (10)

where $\alpha_2 = a/(10^2 R_*)$. Equation (10) can be simplified further: Earlier works showed that the inner disc truncates at a radius of $r_\text{cool} = n^{-2/5} a$ due to the tidal effect, where $n > 2$ is an integer determined by the strongest resonance (Artymowicz & Lubow 1994; Liu et al. 2003). Replacing $r$ in Equation (10) with $r_\text{cool}$, we derive, for the outer boundary of the disc,

$$h/r_\text{cool} > 0.025 \alpha^{-2/5} m^{8/5} M_t^{1/5} r_0^{-2/5} \alpha_2^3 (q(1 + q))^{1/2}.$$  \hspace{1cm} (11)

By equating the right-hand-side (RHS) of Equation (11) to 1, we find a critical separation $a_{ crit}$ for the SMBHB, $a_{ crit} = 20 R_5 \alpha^{-4/3} m^{8/3} M_t^{1/3} r_0^{-1/3} (q(1 + q))^{-2/3}.$  \hspace{1cm} (12)

When $a$ reduces to about $a_{ crit}$, the disc aspect ratio will exceed unity, first at $r \sim r_\text{cool}$. It is interesting that $a_{ crit}$ depends only weakly on each of the model parameters.

4 NUMERICAL SIMULATION

We now use a one-dimensional numerical simulation to derive more accurately $\Sigma$ and $h$ during the squeezing phase, taking into account both tidal and viscous dissipation in the energy equilibrium.

Following Armitage & Natarajan (2002), we shrink the binary separation according to Equation (3) and evolve the surface density of the inner (circum-primary) disc by numerically integrating

$$\frac{\partial \Sigma}{\partial t} = -r^{-1} \frac{\partial}{\partial r} (\Sigma \nu v) / \partial r,$$

where

$$\Sigma \nu v = -3r^{1/2} \partial \left(\frac{v \Sigma^{1/2}}{\partial r} + 2 \Sigma \Lambda/\Omega \right)$$

is the mass advection rate and $\Lambda$ the injection rate of specific angular momentum due to the tidal torque of the binary (Lin & Papaloizou 1986a). To solve $\Lambda$, we use

$$\Lambda = -0.5 f q \Omega^2 r^2 (r/\Delta)^3$$

(Armitage & Natarajan 2002), with $f = 10^{-2}$ a dimensionless parameter that constrains the strength of the torque, $\Delta = \text{max}(R_0, h, r - a)$ and $R_0 = a(q/3)^{1/5}$ is the Hills radius of the secondary black hole. We note that the above scheme allows us to compute $v$ numerically without assuming the relationship $v_r = -r/t_{cool}$. We also implemented a smoothing of the tidal torque (Tazzari & Lodato 2015) but found little difference in the results.

The difference between our approach and the one from Armitage & Natarajan (2002) lies in the calculation of $h$. While in that work the authors did not allow $h$ to vary with time, we evolve $h$ according to the so-called “$\beta$-disc model” (Shakura & Sunyaev 1973), also used by Lodato et al. (2009) and Chang et al. (2010).

In this model, the viscosity $\nu$ is proportional to $h$ as well as the gas pressure, which is only a fraction $\beta = p_{\text{gas}}/(p_{\text{gas}} + p_{\text{rad}})$ of the total pressure. As a result, $\nu = c_s h^2$, where the sound speed $c_s$, is defined as $c_s^2 = (p_{\text{gas}} + p_{\text{rad}})/\rho$ and $\rho = \Sigma/(2h)$ is the volume density. To simplify the calculations of $p_{\text{gas}}$ and $p_{\text{rad}}$, we assume that both quantities are determined by the mid-plane temperature $T_\ast$ of the disc, such that $p_{\text{gas}} = k T_\ast/(\mu m_p)$ and $p_{\text{rad}} = 4\pi T_\ast^4/(3c)$, with $k$ the Boltzmann constant, $\sigma$ the Stefan-Boltzmann constant, $m_p$ the proton mass and $\mu = 0.615$ the mean particle mass in unit of $m_p$ for a plasma of solar metallicity.

The computation of $h$ relies on three conventional assumptions which are also valid in the system of our interest. (i) Hydrostatic equilibrium in the vertical direction, i.e. $c_s = \Omega$. (ii) Heat is dissipated locally in the form of radiation, i.e. $F = D_r^\gamma + D_\gamma$ with $D_r$ and $D_\gamma$ as described in Section 2. (iii) Photons in the mid-plane are transported to the disc surface by diffusion, so $F = 4rT_\ast^4/(3\pi)$. These assumptions give us a system of three equations:

$$T_\gamma = \frac{3\kappa \Sigma^2 \Omega (9\alpha c^2 \beta - 4 \Lambda) / (64 \pi c)}{1/4},$$

$$\beta = \left[1 + 8r \mu m_p T_\gamma^4 c_s/(3c \kappa \Sigma \Omega)^{1},$$

$$c_s = 8r \mu m_p T_\gamma^4/(3c \kappa \Sigma (1 - \beta)).$$

Using these, as well as the surface density computed from the partial differential equation, we solve $T_\gamma$, $\beta$ and $c_s$ at each radius and time step. More specifically, we derive a quartic function of $T_\gamma$ and find the one real and positive solution for our system. We then use this solution to derive $\beta$, $c_s$, and finally $h$.

The simulation starts with a SMBHB of $M_\ast = 10^2 M_\odot$ and $q = 0.1$ at a separation $a = 100 R_5$, where the squeezing phase is expected to start (Armitage & Natarajan 2002). We set up an accretion disc around the primary black hole using the surface density from Equation (8) with $n = 0.01$. A disc with a lower accretion rate would become radiatively inefficient and already be thick (Narayan & Yi 1995). An initial $h$ higher than ours will result in a higher $h$, as can be seen in Equation (10). For this reason, this simulation provides a lower limit on the scale height. The initial disc is truncated to mimic the tidal interactions (Armitage & Natarajan 2002), and we consider a zero-torque inner boundary.

Figure 1 shows the evolution of $\Sigma$ and $h/r$. Both quantities, in general, increase over time. We checked the aspect ratio $h/r$ in each time step and stopped our simulation when the condition $h/r = 1$ is met. This happens at $a \approx 19 R_5$. By this point, the accretion rate, calculated as $\dot{M} = 2 \pi r v \Sigma$ and Equation (14), has increased to almost $M_{\text{Edd}}$ in most of the disc.

The numerical result agrees remarkably with our analytical prediction: If we calculate $a_{ crit}$ using Equation (12) assuming $\dot{m} = 1$ and $n = 2$, we find that $a_{ crit} = 20 R_5$. Moreover, throughout our numerical simulation $\beta$ is much smaller than one, justifying the assumption of a radiation-supported disc, as is adopted in Sections 2 and 3. We also find that only near the end of the simulation does the cooling time-scale become comparable to the GW-radiation time-scale. Therefore, our assumption of local energy dissipation is valid.

5 DISC EMISSION

The thickening of the inner disc will impact the EM counterpart in several ways. We divide the following discussion into four parts because the disc goes through four consecutive phases where the dominant source powering the radiation changes. For each phase, we first identify the dominant power source, then describe our
method of calculating the radiation, and finally discuss the results. The results are shown in Fig. 2. The top panel shows the bolometric light curve, while the bottom panel shows the expected spectral energy distribution (SED) at selected times.

(i) Squeezing phase: This is the phase that we modelled in the last two sections. During it, tidal heating dominates. The simulation gives us the surface temperature $T$, and since the disc is optically thick ($r \gg 1$), we can calculate the SED of each disc annulus using a black-body model (Frank et al. 2002). The resulting bolometric luminosity $L_{\text{bol}}$ is shown as a function of time in the top panel of Figure 2, as the black solid line labeled ‘i’. The cyan and the blue dots, respectively, refer to the initial condition and the last snapshot of the above numerical simulation. We can see that due to tidal heating the luminosity increases, until it reaches the Eddington luminosity $L_{\text{Edd}} = 0.1 M_{\odot} c^{2}$.

(ii) Decoupling phase: When the disc becomes thick, the squeezing process is ineffective ($D_{A} \approx 0$) because fluid elements of the inner disc can cross the binary orbit through the horseshoe orbits or from outside the equatorial plane, due to large effective viscosity and pressure gradient (Lin & Papaloizou 1979; Papaloizou & Lin 1984; Kocsis et al. 2012; Baruteau et al. 2012). The SMBHB goes through a second decoupling—this time from the inner disc. Our code is unable to capture the three-dimensional behaviour of the disc. However, we know that the material ending up outside the binary orbit should remain hot and thick due to the end of the merger, because cooling time-scale is longer than the GW-radiation time-scale (Section 3). Moreover, this material should maintain the same radial distance as that shown in Figure 1, because tidal evolution is no longer important. For these reasons, we assume that at the time of the merger, $t = t_{m}$, the remnant inner disc has an aspect ratio of $h/r = 1$ and an outer boundary at $10 R_{S}$, the same as in the last snapshot of our simulation. Now using Equation (1) we can calculate $F$ and derive $L_{\text{bol}}$, even though we do not know the exact surface density at this point. The result is shown in the top panel of Figure 2 as the black dot. The black dashed line connecting the end of the squeezing phase and the time of merger (with the label ‘ii’) is a rough estimate of the luminosity of the decoupling disc. It increases to a level above three times the Eddington limit.

(iii) Cooling phase: Immediately after the merger, the only heating source is viscosity, which is comparable to $D_{c}$ in the last evolutionary phase. The cooling rate, on the other hand, does not significantly change before and after the merger because the temperature is similar. Therefore, $F \approx D_{c}$. Then we have a situation in which cooling is more important than heating. This imbalance will lead to a drop of the temperature and a decrease of the disc scale height, on a time-scale of $t_{\text{cool}}$, the average cooling time of the disc at the moment of the merger. By the time $t_{m} + t_{\text{cool}}$, the disc would have cooled down such that $F$ becomes comparable to $D_{c}$ again. Considering $F = D_{c}$ and a constant scale height, with the same mass and size than the last snapshot of our simulation, we can apply the standard-disc model and obtain $h$, $T$, $\Sigma$ as a function of $r$. Those assumptions are motivated by noticing that the total luminosity never significantly exceeds the Eddington limit. Consequently, the mass loss due to disc wind, which relies on a super-Eddington luminosity (e.g. Lodato et al. 2009; Tazzari & Lodato 2015), would be much weaker than previously has been thought.

The luminosity at this time is shown as the red dot in the top panel of Figure 2. The evolution of the bolometric luminosity during the cooling phase is represented by the black dashed line connecting the black and the red dots (with the label ‘iii’). We can see
that it drops by more than one order of magnitude, indicating that the disc cools down significantly.

(iv) Recovering phase: After the condition $F = D$, is re-established, the following evolution of the disc is dominated by viscosity. We simulate it with our one-dimensional code setting $\Lambda = 0$. The initial condition is the same as the solution derived in the previous cooling phase. Since the viscous time-scale of this new inner disc is comparable to that of the original outer disc (at $r > 10^3 R_S$), we no longer can assume that the outer disc is invariant. Therefore, we include the outer disc in the simulation, with an initial condition given by Equation 8 and $\dot{m} = 0.01$, as well as a constant outer boundary. The resulting luminosity is shown in the top panel of Figure 2 as the black solid line to the RHS of the red dot (with the label ‘iv’). It first decreases, because the effective accretion rate of the inner disc drops as it expands. By the time the inner and the outer discs overlap (blue dot), the luminosity has reached a minimum. Afterwards, fresh material from the outer part of the disc refills the inner disc, driving the effective accretion rate towards the equilibrium, i.e. $\dot{m} = 0.01$. As a result, the entire disc returns to the thin-disc solution and the luminosity recovers the original value of 0.01$L_{\text{Edd}}$ (purple dot).

Having understood the light curve of a thick disc, let us now compare it with that of a thin disc, which is presented in the top panel of Figure 2 as the grey curves. Three important differences appear.

First, the luminosity of the thin disc rises sharply during the last $O(10^2)$ days of the merger, from $L_{\text{Edd}}$ to more than 10 times higher. It has been proposed that this “precursor” could be used to alert GW detectors for follow ups (Chang et al. 2010). We now see that such a large enhancement is unphysical because the power source–tidal heating–would already have shut down. The luminosity of the thick disc, according to our calculation, increases only 2–3 times during this period.

Second, the luminosity of the thin disc drops immediately after the merger by about four orders of magnitude. This behavior is caused by a complete depletion of the inner disc by the squeezing mechanism, leaving only the emission of the outer disc. It has been pointed out that this cut-off and the later recovering of the luminosity (due to a refilling of the inner cavity) can be used to identify black-hole mergers (Milosavljević & Phinney 2005). Besides, it also has been suggested that the disappearance of the inner disc could explain the interruption of jet activity seen in a sample of radio galaxies (Liu et al. 2003; Liu 2004). Both proposals would have difficulties in the light of our new results, because the inner disc never completely disappears.

The third difference is related to the second one, but is more clearly seen in the evolution of the SED, which is shown in the lower panel of Figure 2. If the inner disc is completely depleted, the SED after the merger would come entirely from the outer disc (grey solid curve) where the corresponding UV and X-ray luminosities are negligible. This part of the SED re-brightens on a time-scale of $10^9$ days because the outer disc refills the inner cavity on the viscous time-scale. In our model, however, the UV/X-ray emission remains present even after the merger (black SED) because the inner disc does not disappear. Only on a time-scale of $O(10^6)$ days after the merger does the UV/X-ray radiation fade away because of cooling (compare the black and red SEDs). This behavior is opposite to that suggested by the thin-disc models.

Therefore, we have seen that the thickening and the second decoupling of the inner disc has important implications for the detectability of the EM counterparts. It is worth noting that the residual inner disc remains gravitationally bound to the merged SMBH even though the black hole will receive a recoil velocity due to anisotropic GW radiation (Centrella et al. 2010). During this recoil the disc could be shock-heated to an even higher temperature, because of the shear induced by a passing GW (Kocsis & Loeb 2008), a loss of gravitational mass (Milosavljević & Phinney 2005; Schnittman & Krolak 2008; Megevand et al. 2009) and/or the orbital change relative to the recoiling SMBH (Lippai et al. 2008; Shields & Bonning 2008; Schnittman & Krolak 2008). Existence of such a hot disc also opens many possibilities of detecting recoiling SMBHs.

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