ESTIMATING THE PROMPT ELECTROMAGNETIC LUMINOSITY OF A BLACK HOLE MERGER

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ABSTRACT

Although recent work in numerical relativity has made tremendous strides in quantifying the gravitational wave luminosity of black hole mergers, very little is known about the electromagnetic luminosity that might occur in immediate conjunction with these events. We show that whenever the heat deposited in the gas near a pair of merging black holes is proportional to its total mass, and the surface density of the gas in the immediate vicinity is greater than the (quite small) amount necessary to make it optically thick, the characteristic scale of the luminosity emitted in direct association with the merger is the Eddington luminosity independent of the gas mass. The duration of the photon signal is proportional to the gas mass, and is generally rather longer than the merger event. At somewhat larger distances, dissipation associated with realigning the gas orbits to the new spin orientation of the black hole can supplement dissipation of the energy gained from orbital adjustment to the mass lost in gravitational radiation; these two heat sources can combine to augment the electromagnetic radiation over longer timescales.

Key words: black hole physics – galaxies: nuclei – gravitational waves

1. INTRODUCTION

The merger of two supermassive black holes has been a topic of lively astrophysical speculation for many years (Begelman et al. 1980). Recent developments in galaxy formation theory have made the prospect more plausible and suggest an environment for such events: the centers of galaxies that underwent major mergers a few hundred million years in the past (Haehnelt & Kauffmann 2002; Volonteri et al. 2003). Mergers may be particularly likely when the galaxy contains a relatively rich supply of interstellar gas, which may help binary black holes overcome the “last parsec problem” and approach each other close enough for gravitational wave emission to compress the orbit to merger within a Hubble time (Gould & Rix 2000; Armitage & Natarajan 2002; Kazantzidis et al. 2004; Escala et al. 2005; Cuadra et al. 2009; Dotti et al. 2009; but for a contrary view see Lodato et al. 2009). The presence of sizable quantities of interstellar gas in the parsec-scale environment then raises the question of how much gas might find itself even closer to the pair at their moment of merger.

This is a subject of great uncertainty. It has been argued, for example, that there should be little gas closer to the merging pair than \(\sim 100–1000 r_g \) \( (r_g \equiv GM/c^2) \), where \( M \) is the total mass of the system because eventually the timescale of shrinkage of the binary orbit by gravitational wave radiation becomes shorter than the timescale for mass inflow due to internally generated fluid stresses (Milosavljević & Phinney 2005). On the other hand, the mass of such a circumbinary disk might be as large as \( \sim 100 M_\odot \) or more (Milosavljević & Phinney 2005; Armitage & Natarajan 2002; Rossi et al. 2009; Corrales et al. 2009); if even \( 1 M_\odot \) were close enough to the merging black holes to be given heat equal to 1% of its rest mass, the total energy—\( \sim 10^{60} \) erg—might well be large enough to produce observable radiation. It is therefore a worthwhile exercise to estimate what sort of light might be generated if even a small fraction of the surrounding gas were able to make its way in close to the merging black holes.

Because the amount of mass near the merging black holes is so difficult to estimate at present, the plan of this paper is to explore prompt electromagnetic radiation in a way that is scaled to whatever gas mass is there. Thus, we will first estimate the heat per unit mass that might be deposited in this gas, then, in order to find the luminosity, estimate the timescale on which the energy is radiated. Next, the more model-dependent subject of the spectrum will be broached. Lastly, having seen how the light emitted depends on gas mass, we will discuss the issues related to whether an “interesting” amount of mass may be present. In order to avoid additional complications, we will ignore any luminosity due to accretion through the circumbinary disk.

2. EFFICIENCY OF ENERGY DEPOSITION

Let us begin, then, with the supposition that immediately before a merger of two supermassive black holes there is at least some gas orbiting over a range of distances not too far from the system’s center of mass. To discuss the effect of the merger on this relatively nearby gas, it is useful to distinguish two regions: the inner gas \( (r/r_g \lesssim 10) \) and gas farther away \( (10 < r/r_g < 10^3) \). These regions are distinguished both by the magnitude of the heating they are likely to experience and by the time at which it occurs.

Because the inner gas is in the “near-field” regime, its gravitational environment during the merger is better described as a nonlinear time-dependent distortion of spacetime, rather than a passage of gravitational waves. The amplitude and extent of the distortions are, in some sense, proportional to the binary mass ratio, reaching a maximum when the two black holes have similar mass. Strong shearing can deflect fluid orbits, provoking shocks, and can also stretch magnetic field lines, driving magnetohydrodynamic (MHD) waves whose energy can ultimately be dissipated. Because the characteristic timescale of the spacetime variability is \( \sim O(10)r_g/c \) (Buonanno et al. 2007), while the dynamical timescale of particle orbits in this region is only a little longer \( \sim (r_g/c)(r/r_g)^{3/2} \), the fluctuations can be very efficient in transferring energy to the gas (cf. the test-particle calculations of van Meter et al. 2009). The same near-coincidence of merger timescale with the gas’s dynamical timescale means that the dissipation timescale (via shocks,
etc., which develop on a dynamical timescale) should also be comparable to the merger duration. Unfortunately, without detailed general relativistic MHD calculations, it is very difficult to make a quantitative prediction of just how much energy might thus be given to the gas.

Nonetheless, to the extent that gravitational dynamics dominate, the Equivalence Principle suggests that the energy left in the gas should be proportional to its mass. It is therefore convenient to scale the gas heating in familiar rest-mass efficiency terms. For a fiducial value, one might imagine that this efficiency \( \epsilon \) could be as large as \( \sim 0.1 \) for matter subjected to a truly nonlinear dynamical spacetime. Gravitational shear acting on magnetic fields may supplement this energy deposition (Moortgat & Kuipers 2004; Duez et al. 2005; Palenzuela et al. 2009).

At greater distances, the perturbations to spacetime are much smaller and the contrast between the wave frequency and the dynamical frequency is much greater. In addition, there can be a considerable delay between the time of the merger and the time at which the principal heating occurs. Three mechanisms can cause heating at these larger distances, one resulting from the sudden loss of mass from the merged black hole due to its emission of gravitational waves (Bode & Phinney 2007), another due to its sudden loss of angular momentum, and a third the result of the merged black hole’s recoil as a result of asymmetric gravitational wave radiation (Lippai et al. 2008). We stress, however, that the radial scales whose radiation is under consideration here are not those responsible for the longer-term afterglow that has been the focus of prior work (Milosavljević & Phinney 2005; Lippai et al. 2008; Shields & Bonning 2008; Schnittman & Krolik 2008; Rossi et al. 2009; Corrales et al. 2009); the afterglow is due to heating of mass in the circumbinary disk proper, at the relatively large radii (at least \( \sim 100 r_g \); Milosavljević & Phinney 2005) where conventional accretion dynamics were able to bring it during the time when gravitational wave evolution of the binary was faster than the typical inflow rate. Although the relevant heating mechanisms in the radial range considered here are very similar to those acting at larger radii, the focus of this work is on matter interior to the disk proper, where any gas present arrived as the result of angular momentum loss faster than that acting on the bulk of the circumbinary disk. Note, too, that because we restrict our attention to \( r/r_g < 10^3 \), all the gas remains bound to the merged black hole for even the largest of recoil velocities.

As a result of the mass lost by the black hole in gravitational wave radiation, the binding energy of orbiting matter is immediately reduced by \( (\Delta M/M)(r/r_g)^{-3} \), where the fractional black hole mass loss \( \Delta M/M \sim 1\%-10\% \) (Berti et al. 2007; Rezzolla et al. 2008; Lousto et al. 2009) with the exact number depending on the mass ratio and spins of the merging pair. Relative to the dynamical timescale, this change in energy occurs almost instantaneously because the merger duration is so much shorter than an orbital period when \( r \gg r_g \). However, any heating due to this change in energy is delayed by \( \sim (r_g/c)(r/g)^{3/2}(\Delta M/M)^{-1} \) because the eccentricity induced in the orbits is only \( \sim \Delta M/M \) (Schnittman & Krolik 2008; O’Neill et al. 2009). In addition, it is possible that only a fraction \( \Delta M/M \) of the orbital energy gained is dissipated, as most of the energy may be used simply to expand the gas orbits (O’Neill et al. 2009; Corrales et al. 2009; Rossi et al. 2009).

There can also be dissipation due to the sudden change in angular momentum. Torques driven by the binary’s quadrupolar mass distribution acting on matter surrounding the pre-merger binary will cause any obliquely orbiting gas to precess around the direction of the binary’s total angular momentum (at large separation, the total angular momentum is almost exactly the orbital angular momentum); dissipation between intersecting fluid orbits should then lead to alignment of the gas’s orbital angular momentum with the binary’s angular momentum (Lubow & Ogilvie 2000). If there is any misalignment between either of the two spin directions and the orbital angular momentum, the angular momentum of the merged black hole could be in a different direction from the original total angular momentum (Schnittman 2004; Bogdanović et al. 2007; Rezzolla et al. 2008). After the merger, Lense–Thirring torques will act in a fashion closely analogous to the Newtonian torques acting during the binary phase, and dissipation should then reorient orbiting gas into the new equatorial plane (Bardeen & Petterson 1975). Because the kinetic energy of motion out of the new equatorial plane is a fraction \( \sim \sin^2(\Delta \theta) \) of the orbital energy for misalignment angle \( \Delta \theta \), the amount of energy dissipated can be an order unity fraction of the orbital energy. The delay from the time of merger to when this mechanism acts will be of order the precession time, \( \sim (r/r_g)^3(r_g/c), \) which is rather longer than the delay before dissipating the orbital energy gained from mass loss wherever \( r/r_g > (\Delta M/M)^{2/3} \). However, the degree of misalignment is highly uncertain and may depend strongly on details of the environment; for example, accretion during the inspiral may align both black hole spins with the orbital angular momentum, eliminating a change in angular momentum direction as a result of the merger (Bogdanović et al. 2007). Close alignment of both spins with the orbital angular momentum will, however, have the compensating effect of producing an especially large \( \Delta M/M \) (Lousto et al. 2009). Because both the mass loss and the Lense–Thirring effects scale with the local binding energy, we will combine them, writing their efficiency as \( \epsilon \sim \eta(r/r_g)^{-3} \), where we expect \( \eta \lesssim 0.01 \).

Lastly, response to the black hole recoil adds an energy per unit mass to the disk matter \( \sim v_{\text{recoil}}^2 \). Even if only a fraction \( \Delta M/M \) of the mass-loss energy leads to heating, the recoil energy becomes larger than the heating due to mass loss only at radii

\[
r > \left( \frac{\Delta M/M}{v_{\text{recoil}}/c} \right)^2 \approx 10^3 \left( \frac{\Delta M/M}{0.03} \right)^2 \left( \frac{v_{\text{recoil}}}{300 \text{ km s}^{-1}} \right)^{-2},
\]

(Schnittman & Krolik 2008; Rossi et al. 2009), so we neglect it in the estimates presented here.

3. COOLING TIME AND LUMINOSITY

Whatever heat is deposited in the gas, the rate at which this energy is carried away by radiation is determined by how rapidly electrons can generate photons and then by how rapidly those photons can make their way outward through the opacity presented by the material itself. We will consider the first issue later (Section 4) when we discuss the complications of estimating the radiating gas’s temperature. For the time being we will assume that photon diffusion is the slower of the two processes. Unless \( \epsilon \) is extremely small, the gas is likely to be so thoroughly ionized that electron scattering dominates its opacity, so the optical depth is simply proportional to the surface density.
To be optically thin, the surface density $\Sigma$ must then be very small: $\Sigma < 3 \, \text{gm cm}^{-2} \lesssim 3 \times 10^{-9} M_7^2 (M_\odot / r_g^2)$, where $M_7$ is the total black hole mass in units of $10^7 M_\odot$. In this case, if the gas is able to convert heat into photons at least as fast as the heat is delivered, the luminosity per logarithmic radius is the ratio of energy deposited to the time in which it is dissipated, roughly the dynamical time:

$$\frac{dL}{d\ln r} \simeq 2.5 \times 10^{14} \epsilon (r/r_g)_{1/2} M_7 \, \text{erg s}^{-1}, \quad (2)$$

for $\Sigma$ in gm cm$^{-2}$. The duration of such a flare should be only $\sim 50 (r/r_g)_{1/2} M_7$ s.

A more interesting regime is presented by the case of an optically thick region. Under optically thick conditions, if the vertical scale height of the disk is $h$, the cooling time $t_{\text{cool}} \sim \tau h / c \sim \tau (r/r_g)(r/r_g)(r_g/c)$, where the optical depth from the midplane outward is $\kappa = \kappa \Sigma / 2$ for opacity per unit mass $\kappa$. Because the duration of the merger event is only $\sim O(10) r_g / c$, unless the disk is able to stay very geometrically thin (which the following argument will demonstrate is unlikely), the heat given the gas during the merger will be radiated over a time much longer than the merger event proper.

The disk thickness is controlled by two elements: the initial heat content of the disk and the heat deposited as a result of the merger. We will neglect the former partly because it is so uncertain, partly because it seems plausible that it will be outweighed by the latter, and partly because this represents in some respects a conservative assumption.

Suppose, then, to begin that the optical depth is large enough to make the cooling time longer than the orbital time (i.e., $\tau (h/r) > (r/r_g)^{1/2}$). There will then be time for the gas to achieve a dynamical equilibrium (if one is possible). Radiation pressure will likely dominate gas pressure because the photon escape time is much longer than the photon radiation time (the reasoning behind this assertion is discussed in Section 4). Put another way, nearly all the thermal energy density initially given the matter is rapidly converted into photons; when the cooling time is longer than the dynamical time, they are still present for many dynamical times. Consequently, their pressure becomes much greater than the gas pressure.

The force exerted by the (slowly) diffusing photons is proportional to their flux times the opacity; because the flux is the energy per unit area per cooling time, we can determine $h$ by matching the vertical radiation force to the vertical component of gravity. Here, in order to obtain a rough estimate of the disk thickness, we suppose that $h \ll r$ and that Newtonian gravity applies:

$$\kappa \epsilon \Sigma c^2 \simeq h \mathcal{Q}^2, \quad (3)$$

which leads to

$$h/r \simeq \epsilon^{1/2} (r/r_g)^{1/2}. \quad (4)$$

In other words, the geometric profile of optically thick gas immediately post-merger depends only on the heating efficiency $\epsilon$ (if it were optically thin, $h/r \simeq (\tau r c / r_g)^{1/2})$. Moreover, if this equilibrium is achieved, the criterion that the cooling time exceed the dynamical time is easily achieved in optically thick disks if $\epsilon$ is not too small, for all that is required is $\tau > \epsilon^{-1/2}$.

Close to the black hole, where $\epsilon$ may be as much as $\sim O(0.1)$, the disk may be almost spherical. It is imaginable that $\epsilon$ is so large that no hydrostatic equilibrium is possible (i.e., $h/r > 1$); in that case, the radiation flux would drive the gas away from the black hole merger remnant. For the purposes of this order-of-magnitude treatment, we ignore that possibility; we also ignore further order-unity corrections that might result from time-dependent photon diffusion effects. Further from the black hole, where $\epsilon \sim \eta (r/r_g)^{-1}$, the disk should be thinner: $h/r \sim \eta^{1/2}$.

In these more distant regions, cancellation of the radial scalings leaves the thickness to be determined by the details of local heat dissipation (i.e., the effectiveness of dissipating the energy gained due to mass loss and whatever disk re-orientation takes place).

With the disk thickness estimated, the cooling time immediately follows:

$$\tau_{\text{cool}} = \tau h/c \simeq 50 \epsilon (r/r_g)^{3/2} M_7 \, \text{s}. \quad (5)$$

Equivalently, it is $\simeq (\tau / 2 \pi) \kappa^{1/2}$ orbital periods.

Although the radiating timescale is proportional to the surface density, the luminosity is independent of it so long as the heating time is shorter than the cooling time. This is simply because the energy to be radiated is proportional to the surface density, while the cooling time is likewise:

$$\frac{dL}{d\ln r} \simeq 1.5 \times 10^{15} \epsilon (r/r_g)^{1/2} M_7 \, \text{erg s}^{-1} = \epsilon^{1/2} (r/r_g)^{1/2} L_E. \quad (6)$$

The luminosity scale is the Eddington luminosity because the time to cross the heated (and inflated) radiating region is proportional to how well the radiation flux can resist gravity. How close the luminosity approaches to Eddington is determined by the efficiency.

Farther from the black hole, the efficiency $\epsilon = \eta (r/r_g)^{-1}$, so if $t_{\text{cool}} > t_{\text{heat}}$ and $\tau > 1$,

$$\frac{dL}{d\ln r} \simeq \eta^{1/2} L_E. \quad (7)$$

That is, the luminosity from regions where $r/r_g \gg 1$ should scale with the Eddington luminosity, but may be a fairly small fraction of it.

However, in these more distant regions (i.e., $10 \lesssim r/r_g < 10^3$), the heating time can be so extended that it might be longer than the cooling time, particularly if the optical depth is not very large:

$$t_{\text{cool}} \sim \tau_\eta^{1/2} \left( \frac{r/r_g}{\eta (r/r_g)^2} \right)^{-1/2} \text{ mass loss} \left( \frac{r/r_g}{\eta (r/r_g)^2} \right)^{-1/2} \text{ Lense–Thirring}. \quad (8)$$

When the dynamical response of the disk is so slow that it exceeds the cooling time, the luminosity becomes

$$\frac{dL}{d\ln r} = \eta \tau L_E \left( \frac{r/r_g}{\eta (r/r_g)^2} \right)^{-1/2} \text{ mass loss} \left( \frac{r/r_g}{\eta (r/r_g)^2} \right)^{-1/2} \text{ Lense–Thirring}. \quad (9)$$

To find the total observed luminosity, we must assemble the luminosity from different regions, making proper allowance for their different time dependences. Initially, the radiative output will be dominated by the inner radii, so that $L \simeq (\epsilon r_{in} / r_g)^{1/2} L_E$, where $r_{in}$ is the scale of the region subject to the truly dynamical spacetime. After a time $50 \epsilon (r/r_g)^{3/2} M_7$ s, this light decays, to be replaced over longer timescales by the signals due to mass loss and the Lense–Thirring re-orientation: $\sim 50 (r/r_g)^{3/2} M_7 (\Delta M / M)^{-1}$ s for the former, $\sim 50 (r/r_g)^3 M_7$ s for the latter. As estimated above, the luminosity from these regions at somewhat larger radius is likely rather less than the luminosity issuing from the innermost region.
4. TEMPERATURE AND SPECTRUM

Prediction of the output spectrum is much more model dependent. If there is enough absorptive opacity to thermalize the radiation, its characteristic temperature would be

\[ T \sim 1 \times 10^6 e^{-1/8} (r/r_g)^{-3/8} M_7^{-1/4} \text{ K}, \]

similarly universal, decreasing only slowly with increasing black hole mass.

Whether thermalization can be achieved, however, may be sensitive to conditions. The effectiveness of free–free opacity can be enhanced by the additional path length to escape imposed upon the photons by scattering. At the temperature just estimated, the effective optical depth (i.e., the geometric mean of the free–free and Thomson optical depths) for photons near the thermal peak is

\[ t_{\text{eff}} \sim 3 \times 10^{-4} r^{3/2} \epsilon^{-15/32} (r/r_g)^{-3/2} M_7^{-1/16} \text{ s}. \]

Thus, if free–free is the only absorption mechanism, \( \tau > 100 e^{-0.15}\epsilon^{-3.2} (r/r_g)^{-3.2} M_7^{-0.16} \) is required to achieve thermalization. On the other hand, where the temperature is \( \sim 10^5 \text{ K} \) or less, atomic transitions substantially enhance the absorptive opacity, making thermalization much more thorough. Thus, the detailed character of the emitted spectrum could vary considerably from case to case.

With an estimate of the temperature, we can now estimate the timescale on which electrons radiate most of the heat given the gas. Considering only free–free emission, it is

\[ t_{\text{rad}} \sim 60 e^{1/2} (r/r_g)^{3/2} M_7 \tau^{-1/2} T_5^{1/2} \text{ s}. \]

In other words, \( t_{\text{rad}}/t_{\text{cool}} \approx 12 \tau^{-2} T_5^{1/2} \), so that thermal balance at a temperature \( \lesssim 10^5 \text{ K} \) is a self-consistent condition for an optically thick region.

However, unless \( \epsilon \) is exceedingly small, the gas’s temperature immediately upon being heated will be far higher than \( 10^5 \text{ K} \). If the initial shocks driven by the dynamical spacetime have speeds comparable to the orbital speed, the post-shock electron energies will more likely be well in excess of 1 MeV, and the characteristic radiation rate will be much slower. In this initial high-temperature state, the free–free radiation time is

\[ t_{\text{rad}} \sim 1 \times 10^4 e^{1/2} (r/r_g) M_7 \tau^{-1} (\ln \Theta)^{-1} \text{ s}, \]

where \( \Theta \) is the electron energy in rest-mass units (assumed to be \( \sim 1 \) in this expression), and we have estimated the disk aspect ratio \( h/r \sim \epsilon^{1/2} \) (by assumption, at this stage radiation pressure is not yet important). Comparing this estimate of \( t_{\text{rad}} \) to \( t_{\text{cool}} \), we find that

\[ \frac{t_{\text{rad}}}{t_{\text{cool}}} \approx 200 \tau^{-2} (\ln \Theta)^{-1}, \]

entirely independent of \( h/r \) (as well as \( r/r_g \) and \( M \)). Thus, if free–free radiation is the only photon-creation mechanism, optical depths \( \gtrsim O(10) \) would be required in order to create photons carrying most of the heat in a time shorter than it takes for those photons to escape.

There are, however, other processes that can also likely contribute. Suppose, for example, that the magnetic field energy density is a fraction \( q \) of the plasma pressure. Synchrotron radiation would then cool the gas on a timescale

\[ t_{\text{rad}} \sim 50q^{-1} \tau^{-1} \Theta^{-2} (r/r_g) M_7 \text{ s}. \]

wherever \( \Theta > 1 \). Relative to the cooling, this photon production timescale is

\[ \frac{t_{\text{rad}}}{t_{\text{cool}}} \sim q^{-1} \tau^{-2} \Theta^{-2}. \]

Inverse Compton radiation would be equally effective if the energy density in photons of energy lower than \( \sim \Theta_{\text{cm}} c^2 \) is that same fraction \( q \) of the plasma pressure. As shown by Equation (14), relativistic electron free–free radiation is able to radiate at least \( \sim 10^{-2} \) of the heat during a photon diffusion time if \( \tau > 1 \); we therefore expect the \( q \) in photons to be at least this large. Thus, even a small initial cooling by free–free radiation, particularly when supplemented by a modest magnetic field, should provide enough seeds for inverse Compton cooling to allow the gas to radiate the great majority of its heat content in a time shorter than the photon diffusion time. All that is required is \( \tau > 10 (q/0.01)^{-1/2} \Theta^{-1} \).

Even if the majority of the dissipated energy is given to the ions, so that their temperature is larger than the electrons’, rapid radiation is still likely to occur. The ratio between the timescale for thermal coupling between ions and electrons by Coulomb collisions and the photon diffusion time is

\[ \frac{t_{\text{rad}}}{t_{\text{cool}}} \sim 60 \tau^{-2} \Theta^{3/2}. \]

Thus, provided \( \tau > 8 \Theta^{3/4} \), the plasma should achieve a one-temperature state more rapidly than the photons can leave.

5. WHAT IS \( \Sigma(r) \)?

Lastly, we turn to the hardest question to answer at this stage: how much gas there should be as a function of radius, here parameterized as \( \tau(r) \). Particularly in the inner region, it would take very little gas to create a large optical depth: even integrated out to \( r/r_g = 100 \), a disk with constant \( \tau = 100 \) would require only \( 10^{-4} M_7^2 M_\odot \).

Even though there is no reason to think the disk is anywhere near a conventional state of inflow equilibrium, one could use the optical depth of such a disk as a standard of comparison. The thermodynamics of equilibrium disks creates a characteristic scale for the surface density: the maximum at which thermal equilibrium can be achieved. One of the predictions of the Shakura & Sunyaev (1973) model is that in a steady-state disk in which the vertically integrated \( r-\phi \) stress is \( \alpha \) times the vertically integrated total pressure, the accretion rate at any particular radius increases as the surface density increases, but only up to a point. Larger surface density (and accretion rate) leads to a larger ratio of radiation to gas pressure. If radiation pressure exceeds gas pressure, increasing accretion rate can only be accommodated by a decreasing surface density. In other words, there is a maximum possible surface density. Although recent work on explicit simulation of disk thermodynamics under the influence of MHD turbulence driven by the magnetorotational instability has shown that this phenomenological model’s prediction about the thermal stability of disks is wrong (Hirose et al. 2009b), they also show that vertically integrated disk properties averaged over times long compared to a thermal time do match those predicted by the \( \alpha \) model (Hirose et al. 2009a): when radiation pressure dominates, the surface density and accretion rate are inversely related. The Thomson optical depth corresponding to this maximum surface density is

\[ \tau \sim 2.5 \times 10^4 (\alpha/0.1)^{-7/8} M_7^{1/8} (r/10 r_g)^{3/16}, \]
where we have scaled the stress/pressure ratio to 0.1. Close to the black hole, it occurs at a comparatively low accretion rate in Eddington units:

$$\dot{m} \simeq 1.4 \times 10^{-3} (\alpha/0.1)^{-1/8} (r/10r_g)^{21/16} \dot{M}_\odot^{1/8}.$$ 

Such a state might be consistent with a gas supply rate at large radius capable of feeding an active galactic nucleus (AGN; i.e., $\dot{m} \sim 0.1$), but reduced 2 orders of magnitude by the effects of binary torques and the inability of internal stresses in the disk to drive its inner edge inward as fast as gravitational wave emission compresses the black hole binary. It is significant in this respect that even such a strong suppression of accretion still yields an inner disk optical depth that is quite large. A smaller accretion rate would produce a smaller optical depth, but only $\propto \dot{m}^{3/5}$, when gas pressure dominates and the disk remains radiative.

These estimates can also serve as a springboard to gain a sense of what might occur in states of inflow non-equilibrium. For example, if gas accretes at larger radii but is held back by binary orbital torques at radii several times the binary separation, its surface density at the point where it is held back should be larger than what would be expected on the basis of equilibrium inflow at the accretion rate farther out. In such a case, although the preceding estimates might be reasonable at larger radii, the optical depth could be substantially enhanced closer to the merging binary.

Another uncertainty is presented by the question of whether the optical depth in a given region remains the same over the entire radiating period. In the inner disk, for example, there could be significant radial motions, both inward or outward, engendered directly by the dynamical spacetime during the merger event. Because the cooling time in the inner disk is rather longer than an orbital period, there might be time for the magneto-rotational instability to stir MHD turbulence that could drive accretion and restore some of the merger heat carried off by radiation. If the optical depth is relatively small, so that $\tau \propto (r/r_g)^{-1/2} \propto (r/r_g)^{3/4}$, the luminosity would gradually taper off as the accretion luminosity extends the bright period, but disk cooling causes the accretion rate to diminish. Alternatively, if the optical depth is larger, the inflow time would actually be shorter than the cooling time. In this case, the luminosity would be greater than earlier estimated, but the lifetime of bright emission from the inner disk would be reduced to the inflow time. In either case, the total energy released due to accretion could increase the total emission by a factor of order unity or more because the radiative efficiency of accretion near a black hole is generically $\sim O(0.1)$.

6. SUMMARY

To summarize, when a pair of supermassive black holes merge, provided only that the gas very close to the merging pair has at least a small electron scattering optical depth, we expect the prompt electromagnetic (EM) signal to likely have a luminosity comparable to the Eddington luminosity of the merged system, $\sim 10^{39} \propto M_\odot$ erg s$^{-1}$. The duration of this bright phase is proportional to the path length of gas at $\sim 10r_g$ from the merged black hole, a quantity that is at present extremely difficult to estimate. If there is only the very small amount of gas necessary to be optically thick, the duration would be only $\sim 10 \dot{M}_\odot$ minutes; on the other hand, quantities several orders of magnitude larger are well within the range of plausibility. Gas at somewhat greater distance ($10 < r/r_g < 10^3$) should continue to radiate at a level possibly as high as $\sim 0.1 L_E$, but for a longer time.

Because gravitational wave observatories like the Laser Interferometric Space Antenna (LISA) are expected to give approximate source localization days or even weeks in advance of merger (Lang & Hughes 2009), one could hope for synergistic observing campaigns that might catch the entire EM signal. Alternatively, it may be feasible to use EM surveys with large solid-angle coverage to search for source flaring having the characteristics predicted here in order to identify candidate black hole merger systems before any gravitational wave detectors are ready.

The spectrum of this light is, at present, much more difficult than its luminosity to predict with any degree of confidence. It may peak in the ultraviolet, but if it does, it is likely to be rather bluer than the familiar UV-peaking spectra of AGNs (if the spectrum does peak in the UV, extinction in the host galaxy may obscure some number events). While the inner region still shines, its luminosity will likely be greater than that from greater radii, making the spectrum (to the degree it is thermalized) closer to that of a single-temperature system. If it is only incompletely thermalized, a still harder spectrum might be expected. At later times, when the outer disk dominates, the temperature corresponding to a given radius is $\propto (r/r_g)^{-1/2} M_\odot^{1/4}$. Compared to the temperature profile of an accretion disk in inflow equilibrium ($T \propto (r/r_g)^{-3}$), this is a slower decline outward, suggesting a spectrum that might be softer than during the initial post-merger phase, but still harder than that of typical AGNs. In addition, over time its high frequency cutoff will move to lower frequencies. Because our understanding of what determines the spectra of accretion disks around black holes is far less solid than our understanding of their gross energetics (consider, for example, the still-mysterious fact that ordinary AGNs radiate a significant fraction of their bolometric luminosities in hard X-rays), these spectral predictions are far shakier than the predicted bolometric luminosity scale.

After the merger heat has been radiated, the disk should revert to more normal accretion behavior and, beginning at the smallest radii, should display standard AGN properties (Milosavljević & Phinney 2005). Unfortunately, the pace of fading as a function of radius depends on the run of surface density with radius, which at the moment is highly uncertain. It is therefore far beyond the scope of even this speculative paper to guess how rapidly that may occur.

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