Spins of primordial binary black holes before coalescence

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Abstract. Primordial stellar-mass black holes, which may contribute to dark matter and to the observed LIGO binary black hole coalescences, are expected to be born with very low spins. Here we show that accretion mass gain by the components of a primordial black hole binary from the surrounding matter could lead to noticeable spins of the components prior to the coalescence provided high initial orbital eccentricities.

Keywords: astrophysical black holes, GR black holes

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Contents

1 Introduction 1

2 Accretion mass gain by binary BH components 2
  2.1 Circular orbits 2
  2.2 Elliptical orbits 3

3 Discussion 5

4 Conclusion 8

1 Introduction

The discovery of coalescing binary black holes (BHs) heralded the advent of gravitational wave (GW) astronomy [1]. Presently, the GWTC-1 catalog of binary coalescences detected by LIGO/Virgo GW interferometers includes 10 BH+BH binaries and one NS+NS (GW170817) binary [2]. A statistical analysis of properties of coalescing binary BHs [3] suggests that the spin distribution of BHs prior to coalescence favors low spins of the components.\(^1\)

The origin of the observed BH binaries is not fully clear. While the evolution of massive binary systems [5–7] is able to reproduce the observed masses and effective spins of the LIGO BH+BH sources [8–11], the alternative (or additional) mechanisms of the binary BH formation is not yet excluded. These channels include, in particular, the dynamical formation of close binary BH in dense stellar clusters [12, 13] or coalescences of primordial black hole binaries which can constitute substantial fraction of dark matter [14–20]. Primordial BHs may form clusters (see [21] for a review) facilitating the formation of binary BHs.

In this note, we focus on the last possibility in order to understand whether primordial binary BHs formed in the early Universe can have noticeable spins before the coalescence. Originally, the spins of primordial BHs should be close to zero (at a percent level at most, see e.g. recent studies [22, 23]). However, when in a binary system, accretion of matter will inevitably bring angular momentum, and the components of a binary BH should acquire spins.

The spin of a BH with mass \(M\) and angular momentum \(J\) is characterized by the dimensionless parameter \(a = J/(GM^2/c)\), where \(G\) and \(c\) are the Newtonian gravity constant and speed of light, respectively. Below we will use geometrical units \(G = c = 1\). We will measure masses in solar mass units, \(M_\odot = 2 \times 10^{33} \text{g}\), \(m = M/M_\odot\), so that the length unit is \(1[\text{cm}] = 2/3 \times 10^{-5} m\), the time unit is \(1[\text{s}] = 2 \times 10^5 m\), etc.

It is easy to estimate the final spin of an initially Schwarzschild BH. Assuming that no accreted mass \(\Delta M = M_f - M_0\) is radiated away, the BH spin after acquiring mass \(\Delta M\) reads [24]:

\[
a^* = \sqrt{\frac{2}{3}} \left(\frac{M_0}{M_f}\right) \left[ 4 - \sqrt{18 \left(\frac{M_0}{M_f}\right)^2 - 2} \right].
\]

\(^1\)An independent analysis of LIGO O1 data discovered one more possible BH+BH binary, GW151216, which may have rapidly spinning aligned components, but with a low astrophysical probability \(\sim 0.71\) [4].
This formula is valid insofar as \( M_f/M_0 < \sqrt{6} \). For a larger final BH mass, \( a^* = a_{\text{max}}^* = 1 \) (more precisely, \( a^* \approx 0.998 \), if one takes into account photon drag from accretion-generated radiation, [25]). If \( \Delta M \ll M_0 \), the acquired BH spin is \( a^* \approx \frac{9}{\sqrt{6}}(\Delta M/M_0) \).

It is easy to estimate the accretion mass gain for a single BH. Suppose it is immersed in a medium with sound velocity \( c_s \). Typically, in the interstellar medium \( c_s \sim \sqrt{T} \frac{T}{1\text{eV}} \) or less (here \( T \) is the temperature of the medium). Assuming a Bondi-Hoyle-Lyttleton accretion onto the BH, we find

\[
\Delta M/M_0 \approx \frac{4\pi \rho m^2}{v^2 + c_s^2} a_{\text{max}}^* \frac{t_0}{M_0^3/2},
\]

where \( \rho \) is the density of the medium, \( v \) is the proper velocity of the BH relative to the medium, \( t_0 \) is the duration of the accretion. For example, for the typical ISM density \( \rho \sim 10^{-24} \text{g cm}^{-3} = (27/16) \times 10^{-42} m^{-2} \) and a maximum possible Hubble time \( t_0 = t_H = 4 \times 10^{17} s = 8 \times 10^{22}m \), by neglecting the BH motion, from eq. (1.2) we obtain \( \Delta M/M_0 \approx 1.710^{-3} m \ll 1 \) and the final spin \( a^* \approx 3.76 \Delta M/M \approx 0.006 m \), i.e. fairly small. For a 30–50\( M_\odot \) BH this would give a noticeable value but it is hard to measure the mass and spin of a single BH.

The situation is less certain for the initially non-rotating components of a binary BH that is able to coalesce over the Hubble time. Below we calculate the accretion mass gain by the components of such a binary and show that the acquired spins can be interesting only if the initial orbital eccentricity of the binary is large.

2 Accretion mass gain by binary BH components

Consider a binary system consisting of two point-like masses \( m_1, m_2 = m_1/q \) (\( q \) is the binary mass ratio). The total mass is \( M = m_1 + m_2 = m_2(1+q) \), the orbital period \( T \) is found from the 3d Kepler’s law \( 4\pi^2/T^2 = M/a^3 \), where \( a \) is the orbital semi-major axis.

2.1 Circular orbits

Let us start with the simplest case of a circular orbit. For typical BH+BH binaries with \( m \sim 10–50 \), orbital velocities even at the maximum initial separations allowing for the coalescence over the Hubble time are much larger than the ISM sound velocity, so we will neglect \( c_s \) in the Bondi-Hoyle-Lyttleton formula. For the \( i \)-th component \( (i = 1, 2, j = 3 - i) \) moving with the velocity \( \upsilon_i \), the mass accretion rate (see section 3 for the discussion of the numerical coefficient) is

\[
\dot{M}_i = 4\pi \rho m^2 \frac{a^3}{\upsilon_i^3} = \frac{4\pi \rho m^2}{m_j^3} \frac{a^{3/2} M^{3/2}}{M^{3/2}},
\]

where we have used the expression for the Keplerian orbital velocity of the \( i \)-th component \( \upsilon_i = \sqrt{m_j^2/aM} \). The binary loses the energy and angular momentum due to emission of gravitational waves, and during the time before the coalescence the mass gain by the 1-st component (for definiteness) will to good accuracy read

\[
\Delta M_1 = \int_0^{t_0} \dot{M}_1 dt = \int_0^{a_0} \frac{dM_1}{da} \frac{dt}{da} da,
\]
Figure 1. The fractional accretion mass gain by a coalescing binary BH system in a circular orbit over the Hubble time in a cold medium with density $1 \text{ cm}^{-3}$.

where the initial orbital separation $a_0$ of the binary is uniquely determined from the GW-driven coalescence time

$$t_0 = \frac{5a_0^4}{256Mm_1m_2} \quad (2.3)$$

and $dt/da$ is found from the quadrupole GW formula for a circular binary system:

$$dt = -\frac{5a^3}{64m_1m_2M}da. \quad (2.4)$$

After taking the integral in eq. (2.2) and substituting $a_0$ through $t_0$ from eq. (2.3), we arrive at:

$$\frac{\Delta M_1}{M_1} \bigg|_0 = \frac{5\pi \rho M^{1/2}a_0^{11/2}}{88m_2^3} = \frac{5}{88} \left( \frac{256}{5} \right)^{11/8} \pi \rho t_H^{11/8} m_1^{5/8} q^{3/4} (1 + q)^{15/8}. \quad (2.5)$$

The plot of $\Delta M_1/M_1$ as a function of $m_1$ is shown in figure 1 for different mass ratios $q = m_1/m_2$ for the fiducial ISM density $1 \text{ g cm}^{-3}$. Clearly, the effect increases both with $m_1$ and $q$ but even for large $q > 1$ (i.e., when we consider the mass gain by the heaviest binary component) is desperately small to enable astrophysically interesting BH spins, even for larger densities.

2.2 Elliptical orbits

The case of initially elliptical orbits is more interesting. Elliptical orbits of binary BHs are possible in both the dynamical channel of binary BH formation in dense stellar clusters and for primordial BHs.
Consider a Keplerian binary in an elliptical orbit with eccentricity $e_0$. The mass accreted over one orbital revolution with period $T$ reads:

$$
\delta M_1 = \int_0^T \dot{M}_1 dt = 2 \int_0^\pi \dot{M}_1 \left( \frac{dt}{d\theta} \right) d\theta = 8\pi \rho q^2 (1+q)(1-e^2)a^3I_1(e)
$$

(2.6)

where

$$
I_1(e) = \int_0^\pi \left( 1 + e \cos \theta \right)^2 \left( 1 + 2e \cos \theta + e^2 \right)^{3/2} - 1 d\theta .
$$

(2.7)

Here we have used the expressions for the orbital velocity $v_i(\theta) = \sqrt{M/a(1-e^2)(1+2e \cos \theta+e^2)} \left( \frac{m_j}{M} \right)$, the orbital angular momentum conservation $r^2 \frac{d\theta}{dt} = \sqrt{Ma(1-e^2)}$ and $r = a(1-e^2)/(1+e \cos \theta)$ for the Keplerian motion.

The mass accretion rate averaged over one orbital period $T$ is

$$
\langle \dot{M}_1 \rangle = \frac{\delta M_1}{T} , \quad T = 2\pi \sqrt{\frac{a^3}{M_2(1+q)}} .
$$

(2.8)

In a way similar to the circular case, we find the accretion mass gain by the 1-st component of a binary BH with initial orbital eccentricity $e_0$ coalescing over the Hubble time:

$$
\Delta M_1(e_0) = \int_0^{e_0} \langle \dot{M}_1 \rangle \left( \frac{dt}{de} \right) de ,
$$

(2.9)

where $de/dt$ reads [26]

$$
\frac{de}{dt} = -\frac{304m_1m_2Me}{15a^4(1-e^2)^5/2} \left( 1 + \frac{121}{304} e^2 \right) .
$$

(2.10)

For the coalescing binary, the expression $a(e)$ reads [26]

$$
a(e) = \frac{C_0e^{12/19}}{(1-e^2)^{5/2}} \left( 1 + \frac{121}{304} e^2 \right)^{870/2299},
$$

(2.11)

where the constant $C_0$ is determined by substituting $a(e_0)$ into the formula for the binary coalescence time $t_0$ in the case of elliptical orbit [26]:

$$
t_0 = \frac{5a(e_0)^4}{256Mm_1m_2} \frac{48(1-e_0^2)^4}{19e_0^{48/19}} \left( 1 + \frac{121}{304} e_0^2 \right)^{-3480/2299} I_2(e_0) ,
$$

$$
I_2(e_0) = \int_0^{e_0} \left( 1 + \frac{121}{304} e^2 \right)^{1181/2299} e^{29/19} (1-e^2)^{3/2} de .
$$

(2.12)

Substituting eq. (2.8) into eq. (2.9) with an account of eq. (2.11) and eq. (2.12), we finally obtain the accretion mass gain for the elliptical orbit:

$$
\frac{\Delta M_1}{M_1} \bigg|_{e = e_0} = \frac{15}{76} \left( \frac{304}{15I_2(e_0)} \right)^{11/8} \rho_0^{11/8} m_1^{5/8} q^{3/4}(1+q)^{15/8} \int_0^{e_0} I_1(e)e^{47/19} \left( 1 + \frac{121}{304} e^2 \right)^{226/209} de .
$$

(2.13)
It can be written in the form

\[
\frac{\Delta M_1}{M_1} \bigg|_e = \frac{\Delta M_1}{M_1} \bigg|_0 \cdot K(e_0),
\]

where the enhancement factor \( K(e_0) \) reads:

\[
K(e_0) = \frac{66}{19\pi} \left( \frac{19}{48I_2(e_0)} \right)^{11/8} \int_0^{e_0} I_1(e) e^{47/19} \left( 1 + \frac{121}{304} e^2 \right)^{226/209} \, de.
\]

Clearly, the enhancement factor with respect to the circular orbit is a function of the initial orbital eccentricity only, and is shown in figure 2, left panel. In the \( e_0 \to 0 \) limit, \( I_1(e_0) \sim \pi \) and \( I_2(e_0) \sim (19/48)e_0^{48/10} \), and \( K_{e_0} \to 1 \). In the more interesting limit of large eccentricities \( e_0 \sim 1 \), we find from numerical integration \( I_1(e_0) \sim (1 - e_0^2)^{-1.27} \). Therefore, in this limit \( K(e_0) \sim (1 - e_0^2)^{11/16} \times (1 - e_0^2)^{-3.27} = (1 - e_0^2)^{-2.58} \). This power-law asymptotic is clearly seen on the plot \( \log K(e_0) - \log 1/(1 - e_0^2) \) shown in figure 2, right panel. Therefore, in the limit of high initial orbital eccentricities, we find approximately

\[
\frac{\Delta M_1}{M_1} \bigg|_e \approx 10^{-5} \left( \frac{\rho}{10^{-24} \text{g cm}^{-3}} \right) \left( \frac{M_1}{30M_\odot} \right)^{5/8} q^{3/4} (1 + q)^{15/8} \left( \frac{0.1}{1 - e_0^2} \right)^{2.58},
\]

\[
\approx 10^{-5} \left( \frac{\rho}{10^{-24} \text{g cm}^{-3}} \right) \left( \frac{\mathcal{M}}{30M_\odot} \right)^{5/8} q(1 + q)^2 \left( \frac{0.1}{1 - e_0^2} \right)^{2.58}.
\]

In the last equality, we have introduced the chirp mass of the binary system \( \mathcal{M} \equiv (M_1M_2)^{3/5}/M_1^{1/5} = M_1(q^2(1 + q))^{-1/5} \) that is directly read off the chirp GW signal from binary coalescences. It is seen that for eccentric orbits with \( e_0 \gtrsim 0.95 \) (\( 1/(1 - e_0^2) \gtrsim 10 \)) this factor can bring the mass accretion gain into astrophysically interesting region for BH spin, especially for more massive component of a binary with large mass ratio \( q > 1 \). As an example, in figure 3 we show the fractional accretion mass gain by a BH with mass \( m_1 = 30 \) as a function of the initial orbital eccentricity \( e_0 \) for different binary mass ratios \( q \). Formally, for the assumed ISM density, a noticeable spin of the primary BH component before the coalescence, \( a^* \sim 3.76(\Delta M_1/M_1) \), could be achieved only for very eccentric orbits with \( e_0 \sim 1 \). The effect is stronger for more massive BHs and higher surrounding densities.

3 Discussion

In our analysis, we have neglected the sound velocity \( c_s^2 \) in the Bondi-Hoyle-Lyttleton formula eq. (1.2). This needs to be justified in the case of strongly eccentric orbits because the accretion rate is determined by the maximal of the orbital velocity and the sound velocity \( c_s \). The orbital velocity of star \( M_1 \) at apastron is

\[
v_a(M_1) = \sqrt{\frac{M(1-e)^2}{a(1-e^2)} \frac{M_2}{M}} = \sqrt{\frac{M_1(1-e^2)}{a(1+e)^2q(1+q)}}.
\]

Then the condition \( v_a(M_1) > c_s \) can be written as

\[
1 - e^2 > c_s^2(1+e)^2q(1+q) \left( \frac{a}{M_1} \right).
\]
Figure 2. Left: the enhancement factor $K(e_0)$ of the fractional mass accretion gain in elliptical orbit by a component of a coalescing binary BH relative to the circular case as a function of the orbit eccentricity $e_0$. Right: log $K(e_0) - (-\log(1-e_0)^2)$ plot manifestly showing the asymptotic power-law behaviour at large $e_0 \approx 1, K(e_0) \sim (1-e_0^2)^{-2.58}$.

Figure 3. The fractional accretion mass gain by a $30M_\odot$ black hole in an elliptical binary as a function of the initial orbital eccentricity $e_0$. The cold medium density is $1 \text{ cm}^{-3}$.

Clearly, if the initial orbital eccentricity $e_0$ satisfies this inequality, it will hold always true in the subsequent binary evolution due to GW losses. Making use of eq. (2.12) to express $a(e_0)/M_1$, in the limit $e_0 \to 1$ of interest here we find

$$1 - e_0^2 > 4c_s^2q^{5/4}(1+q)^{5/4}\left(\frac{256}{5} \frac{tH}{m_1}\right)^{1/4}\left(\frac{19}{48(1-e_0^2)^4}\right)^{1/4}\left(\frac{425}{304}\right)^{870/2299}\left[I_2(e_0 \to 1)\right]^{-1/4}$$
Noticing that \( I_2(e_0 \to 1) \approx \left( \frac{425}{30^4} \right)^{1181/2299} (1 - e_0^2)^{-1/2} \), plugging \( t_H/m_1 \) for the fiducial mass ratio \( m_1 = 30 \), after making arrangements, we arrive at the inequality

\[
1 - e_{0,\text{max}}^2(m_1) > 0.01 \left( \frac{c_s}{10^{-5}} \right)^{16/15} \left( \frac{m_1}{30} \right)^{-2/15} (q(1 + q))^{2/3}
\]  

(3.3)

that restricts the applicability of our approximation. Proceeding exactly in the same way as for \( m_1 \), we obtain the restriction for the initial orbital eccentricity for \( m_2 \):

\[
1 - e_{0,\text{max}}^2(m_2) > 0.01 \left( \frac{c_s}{10^{-5}} \right)^{16/15} \left( \frac{m_1}{30} \right)^{-2/15} q^{-6/5}(1 + q)^{2/3}.
\]  

(3.4)

This limit for the initial binary eccentricity, \( e_{0,\text{max}} < 1 - 0.005 \) for the fiducial parameters, leaves quite a room for a significant enhancement factor \( K(e_0) \). To see this, let us estimate the maximum possible effective spin of a coalescing binary BH, \( \chi_{\text{eff}} = (m_1 a_1^* + m_2 a_2^*)/M_c \), which can be inferred from GW observations [2]. Note that in our setup fully aligned BH spins are expected. As \( a_1^* \sim (\Delta M/M)_1 \), we need to calculate also the mass gain by the secondary component, \( (\Delta M/M)_2 \). This is obviously done by substituting \( M_1 \to M_2 \) in eq. (2.16), i.e. simply changing \( q \to 1/q \):

\[
\frac{\Delta M_2}{M_2} \approx 10^{-5} \left( \frac{\rho}{10^{-24} \text{g cm}^{-3}} \right) \left( \frac{M_2}{30M_\odot} \right)^{5/8} q^{-3}(1 + q)^2 \left( \frac{0.1}{1 - e_0^2} \right)^{2.58}.
\]  

(3.5)

If there would be no initial eccentricity restrictions (e.g., in the limit of a cold medium with very low sound velocities \( c_s \)), the effective spin of the coalescing BH binary with \( M = 30M_\odot \) would be

\[
\chi_{\text{eff}} = \frac{q}{1 + q} a_1^* + \frac{1}{1 + q} a_2^*
\]

\[
\approx 3.76 \times 10^{-5} \left( \frac{\rho}{10^{-24} \text{g cm}^{-3}} \right) \left( \frac{M_2}{30M_\odot} \right)^{5/8} \left( \frac{0.1}{1 - e_0^2} \right)^{2.58} (1 + q)(q^2 + q^{-3})
\]

\[
> 5.3 \times 10^{-4} \left( \frac{\rho}{10^{-24} \text{g cm}^{-3}} \right) \left( \frac{M_2}{30M_\odot} \right)^{5/8} \left( \frac{0.1}{1 - e_0^2} \right)^{2.58}
\]  

(3.6)

for any mass ratio \( q \) because the function \( f(q) = (1 + q)(q^2 + q^{-3}) \) reaches the minimum \( f(q_{\text{min}}) = 4 \) at \( q_{\text{min}} = 1 \). However, taking into account the initial eccentricity limits, eq. (3.3) and eq. (3.4), due to finite sound velocity of the medium, we find

\[
\chi_{\text{eff}} < \chi_{\text{eff, max}} = \frac{q}{1 + q} a_1^*(e_{0,\text{max}}(m_1) + \frac{1}{1 + q} a_2^*(e_{0,\text{max}}(m_2)) \propto M^{0.97/c_s^{2.75}} \Psi(q),
\]  

(3.7)

where \( \Psi(q) \) is a function of the mass ratio that can be readily calculated by substituting the factors \( 1 - e_{0,\text{max}}^2(m_{1,2}) \) for \( a_1 \) [eq. (3.3)] and \( a_2 \) [eq. (3.4)], respectively, into eq. (3.6). The plot of \( \chi_{\text{eff, max}} \) for the fiducial values \( c_s = 10^{-5} \), \( \rho = 10^{-24} \text{ g cm}^{-3} \) and \( M = 30M_\odot \) is shown in figure 4. Roughly, we can take \( \chi_{\text{eff, max}} \approx 0.01(\rho/10^{-24} \text{ g cm}^{-3})(M/30M_\odot)^{0.97/c_s^{2.75}} \) for any mass ratio \( 0.1 < q < 10 \). This estimate shows that coalescing primordial binary BHs can acquire measurable values of \( \chi_{\text{eff}} \) a few percents due to accretion mass gain in galactic ISM.
4 Conclusion

Here we presented the results of calculation of the mass gain by components of a BH+BH binary system, which can coalesce over the Hubble time, due to the Bondi-Hoyle-Lyttleton accretion from a relatively cold ($c_s \sim$ a few km s$^{-1}$) surrounding medium. The angular momentum by the accreted material can spin up an the initially Schwarzschild BH up to noticeable values if the initial binary orbit had a high eccentricity $e_0 \sim 1$. Such eccentricities are in principle possible in primordial BH binaries that can be formed in the early Universe and coalesce at the present time.

In our calculations we have assumed the simplest formula for the accretion rate onto a binary components, which is, of course, a rough estimate. For example, recent 3D simulations of Bondi-Hoyle accretion [27] suggest an orbital-averaged reduction of the Bondi-Hoyle accretion efficiency by a factor of $\sim 1/4$ in circular binaries. However, for our purposes this
reduction is not very important in view of much more uncertain density of matter surrounding the coalescing binary. This density can be an order of magnitude higher or smaller depending on the location of the binary in a galaxy or in the galactic halo. Moreover, primordial binary BHs are thought to have high velocity dispersion $\sim 300 \text{ km s}^{-1}$, which drastically reduces the efficiency of matter accretion. Still, some BH binaries could have rather small velocities and can be found inside the galactic ISM. Therefore, in principle, the components of such BH binaries can acquire noticeable aligned spins prior to the coalescence.

The accretion-gained spins should be higher in more massive binaries ($a^* \sim \Delta M/M_0 \sim m_5^{5/8}$, see figure 1). Interestingly, the most massive LIGO binary BH, GW170729 [2] and (not very reliable) recently reported BH binary GW151216 [4] show appreciable and likely aligned spins of the components. Of course, presently it is difficult to separate different formation channels of the observed coalescing binary BHs, and increased statistics of binary BH coalescences in the ongoing O3 LIGO/Virgo run could help disentangling various scenarios of binary BH formation and evolution. However, we stress that even primordial binary BHs could have appreciable aligned spins before the coalescence due to matter accretion in galaxies.

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