Measuring coalescing massive binary black holes with gravitational waves: 
The impact of spin-induced precession

Ryan N. Lang and Scott A. Hughes
Department of Physics and MIT Kavli Institute, 
MIT, 77 Massachusetts Ave., Cambridge, MA 02139

The coalescence of massive black holes generates gravitational waves (GWs) that will be measurable by space-based detectors such as LISA to large redshifts. The spins of a binary’s black holes have an important impact on its waveform. Specifically, geodetic and gravitomagnetic effects cause the spins to precess; this precession then modulates the waveform, adding periodic structure which encodes useful information about the binary’s members. Following pioneering work by Vecchio, we examine the impact upon GW measurements of including these precession-induced modulations in the waveform model. We find that the additional periodicity due to spin precession breaks degeneracies among certain parameters, greatly improving the accuracy with which they may be measured. In particular, mass measurements are improved tremendously, by one to several orders of magnitude. Localization of the source on the sky is also improved, though not as much — low redshift systems can be localized to an ellipse which is roughly $10 \sim 10 \times 10$ arcminutes in the long direction and a factor of 2 smaller in the short direction. Though not a drastic improvement relative to analyses which neglect spin precession, even modest gains in source localization will greatly facilitate searches for electromagnetic counterparts to GW events. Determination of distance to the source is likewise improved: We find that relative error in measured luminosity distance is commonly $\sim 0.1 - 0.4\%$ at $z \sim 1$. Finally, with the inclusion of precession, we find that the magnitude of the spins themselves can typically be determined for low redshift systems with an accuracy of about 0.1% - 10%, depending on the spin value, allowing accurate surveys of mass and spin evolution over cosmic time.

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I. INTRODUCTION

A. Background to this analysis

Observations have now demonstrated that massive black holes are ubiquitous in the local universe. It appears that all galaxies with central bulges contain black holes whose masses are strongly correlated with the properties of the bulge [1, 2]. Hierarchical structure formation teaches us that these galaxies assembled over cosmic history through the repeated coalescence of the dark matter halos in which they reside [3]. Taken together, these suggest that coalescences of massive black holes should be relatively frequent events, especially at high redshift when halo coalescences were common [4].

Massive black hole coalescences are extremely strong gravitational wave (GW) sources. In the relevant mass band — thousands to millions of solar masses — the GWs these binaries generate are at low frequency ($f \sim 10^{-4.5} - 10^{-1}$ Hz) where ground-based GW antennae have poor sensitivity due to geophysical and other terrestrial noise sources. Measuring GWs from massive black holes requires going into the quiet environment of space. LISA, the Laser Interferometer Space Antenna, is being designed as a joint NASA-ESA mission to measure GWs in this frequency band; cosmological massive black hole coalescences are among its highest priority targets. By measuring these GWs, one can infer the properties of the source that generated the waves. Some particularly important and interesting properties are the masses of the binary’s members, their spins, the binary’s location on the sky, and its distance from the solar system barycenter. Measuring a population of coalescence events could thus provide a wealth of data on the cosmological distribution and evolution of black hole masses and spins.

Most analyses of how well binary black hole parameters can be determined by LISA measurements have ignored the impact of spin-induced precession [5, 6, 7]. Under such an assumption, subsets of parameters can be highly correlated with each other, increasing the errors in parameter estimation. One such subset comprises the binary’s “chirp mass” $M$, its reduced mass $\mu$, and the spin parameters $\beta$ and $\sigma$ (which are written out explicitly in Sec. II A). These four parameters influence the GW phase $\Phi$. As discussed in [6, 7], the correlation coefficient between $\mu$ and $\beta$ is nearly 1. It is thus difficult to “detangle” these parameters from one another in a measurement.

Another such subset consists of a binary’s sky position, orientation, and distance. To see why these parameters are strongly correlated, consider the form of the two polarizations of the strongest quadrupole harmonic of the gravitational waveform:

$$h_+(t) = 2M^{5/3}(\pi f)^{2/3} D_L (1 + \cos^2 \iota) \cos \Phi(t), \quad (1.1)$$

$$h_\times(t) = -4M^{5/3}(\pi f)^{2/3} D_L \cos \iota \sin \Phi(t). \quad (1.2)$$

(We work in units with $G = 1 = c$; a convenient conversion factor in this system is $10^9 M_\odot = 4.92$ seconds.) The quantity $\iota$ is the binary’s inclination relative to the
line of sight: \( \cos \ell \equiv \mathbf{\hat{L}} \cdot \mathbf{\hat{n}} \), where \( \mathbf{\hat{L}} \) is the direction of the binary’s orbital angular momentum, defines its orientation and \( \mathbf{\hat{n}} \) is the direction from observer to source. The quantity \( D_L \) is the luminosity distance to the source, and \( f(t) \equiv (1/2\pi) d\Phi/dt \).

One does not measure the polarizations \( h_+ \) and \( h_\times \) directly; rather, one measures a sum \( h_M(t) \) in which the two polarizations are weighted by antenna response functions as follows:

\[
h_M(t) = F_+(\theta_N, \phi_N, \psi_N) h_+(t) + F_\times(\theta_N, \phi_N, \psi_N) h_\times(t).
\]

(This equation should be taken as schematic; see [11] for a more detailed and definitive description.) The angles \( \theta_N \) and \( \phi_N \) denote the location of the source on the sky in some appropriate coordinate system. The angle \( \psi_N \), known as the “polarization angle,” fixes the orientation of the component of \( \mathbf{\hat{L}} \) perpendicular to the line of sight. (In other words, \( \mathbf{\hat{L}} \) is fixed by \( \ell \) and \( \psi_N \).)

Measuring the phase determines chirp mass with high accuracy; the fractional error in \( M \) is often \( \sim 10^{-3} - 10^{-4} \). As far as amplitude is concerned, the chirp mass can be regarded as measured exactly. What remains is to determine, from the measured amplitude and the known \( M \), the angles \( \theta_N \), \( \phi_N \), \( \psi_N \), \( \ell \), and the distance \( D_L \).

As Eqs. (1.1), (1.2), and (1.3) illustrate, these five parameters are strongly correlated. The motion of LISA around the sun\(^1\) breaks these degeneracies to some extent — the angles \( \theta_N \) and \( \phi_N \) appearing in Eq. (1.3) can be regarded as best defined in a coordinate system tied to LISA. As the antenna orbits the sun, these angles become effectively time-dependent. The one year periodicity imposed by this motion makes it possible to detangle these parameters. Analyses typically find that the position of a merger event at \( \sim 1 \) can be determined, on average, to an ellipse which is a degree across in the long direction and a few \( \times 10 \) arcminutes across in the short direction\(^2\). The distance to such a binary can be determined to 1% accuracy on average (less in some exceptional cases)\(^3\).

B. Black hole spin and spin precession

The preceding discussion ignores an important piece of relativistic physics: the precession of each binary member’s spin vector due to its interaction with the spacetime in which it moves. In general relativity, the spacetime of an isolated object can be regarded as having an “electric piece,” arising from the object’s mass and mass distribution, and a “magnetic piece,” arising from the object’s mass currents and their distribution\(^3\). Spin precession consists of a geodetic term, arising from the parallel transport of the spin vector in the gravitoelectric field of the other hole, and Lense-Thirring terms, caused by the gravitomagnetic field of the other hole. The basic physics of gravitomagnetic precession can be simply understood by analogy with a similar (and closely related) electromagnetic phenomenon — the precession of a magnetic dipole \( \mu \) immersed in an external magnetic field \( \mathbf{B} \). An object’s spin angular momentum \( \mathbf{S} \) can be regarded as a gravitational “magnetic dipole.” When immersed in a “gravitomagnetic field,” one finds that \( \mathbf{S} \) feels a torque, just as a magnetic dipole \( \mu \) experiences a torque when immersed in magnetic field \( \mathbf{B} \).

In a binary black hole system, the gravitomagnetic field arises from the binary’s orbital motion and the spins of its members. Precession thus includes both spin-orbit (geodetic and orbital gravitomagnetic) and spin-spin effects\(^4\). (The major goal of the “Gravity Probe B” experiment is to measure the effects of geodetic and spin-spin Lense-Thirring precession upon a gyroscope in low Earth orbit\(^5\).)

As the spins precess, they do so in such a way that the total angular momentum \( \mathbf{J} = \mathbf{L} + \mathbf{S}_1 + \mathbf{S}_2 \) is held constant; the orbital angular momentum \( \mathbf{L} \) precesses to compensate for changes in \( \mathbf{S}_1 \) and \( \mathbf{S}_2 \). As a consequence, the inclination angle \( \ell \) and polarization angle \( \psi_N \) become time varying (as do certain other quantities appearing in the GW phase function \( \Phi \)). Figure 11 shows the so-called “polarization amplitude,” defined in Sec. II C of the waveform measured by a particular detector. Without precession, this quantity is modulated by the orbital motion of LISA, helping to provide some information about the binary’s sky position. The polarization amplitude also depends on the angles \( \ell \) and \( \psi_N \), so it undergoes additional modulation when precession is included. Such precession-imposed time variations quite thoroughly break many of the degeneracies which have been found to limit parameter measurement accuracy in earlier analyses.

It is without a doubt that black holes in nature spin. Observations are not yet precise enough to indicate the value of typical black hole spins; the evidence to date does, however, seem to indicate that fairly rapid rotation is common. For example, the existence of jets from active systems seems to require non-negligible black hole spin — jets appear to be “launched” by the shearing of magnetic field lines (supported by the highly conductive, ionized material accreting onto the black hole) by the differential rotation of spacetime around a rotating black hole\(^6\). Also, observations of highly distorted iron K-\(\alpha\) lines — a very sharp fluorescent feature in the rest frame of

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1 LISA is being designed as a constellation of three spacecraft whose centroid orbits the sun with a period of one year; see \( \text{http://lisa.nasa.gov} \) for further details.

2 It is worth bearing in mind that the full moon subtends an angle of about 30 arcminutes.

3 This analogy is most apt in the weak field. In that limit, one can recast the Einstein field equations of general relativity into a form quite similar to Maxwell’s equations; see [11] for detailed discussion. Though the analogy does not fit quite so well in strong field regions, it remains accurate enough to be useful.
the emitting iron ions — indicate that this emission is coming from very deep within a gravitational potential (at radii less than the Schwarzschild innermost radius $6M$) and is smeared by near luminal relativistic speeds to boot. Though perhaps influenced somewhat by selection effects, these pieces of evidence are strong hints that the black holes which will form the binaries we hope to measure will be strongly influenced by spin.

The only limit in which spin precession can be neglected is that in which the spins of the binary’s members are exactly parallel (or antiparallel) to one another and to the orbital angular momentum $L$. Since the target binaries of this analysis are created by galactic merger processes, their members will almost certainly have no preferred alignment — random spin and orbit orientation is expected to be the rule. (This expectation is borne out by work showing that jets in active galaxies are oriented randomly with respect to the disks of their host galaxies.) Taking into account spin precession is thus of paramount importance for GW observations of merging black hole systems.

A great deal of work has gone into developing families of model waveforms (“templates”) sufficiently robust to detect GWs from spinning and precessing binaries, at least in the context of measurements by ground-based detectors. The key issue in this case is that the various modulations on the waveform imposed by the binary’s precession smear its power over a wider spectral range, making it much more difficult to detect at the (relatively) low signal-to-noise ratios (SNRs) expected for ground-based observations. Not as much work has gone into the complementary problem of measuring these waves — examining the impact precession has upon the precision with which binary properties may be inferred from the waves. To date, the most complete and important analysis of this type is that of Vecchio. Vecchio focuses (for simplicity) on equal mass binaries and only includes the leading “spin-orbit” precession term. This limit is particularly nice as a first analysis of this problem, since it can be treated (largely) analytically (cf. discussion in Sec. III B of Ref.).

Vecchio’s work largely confirms the intuitive expectation discussed above — the precision with which masses are measured is substantially improved; in particular, the reduced mass of the system can be measured with several orders of magnitude more accuracy. Parameters such as the sky location of the binary and the luminosity distance are also measured more accurately, but only by a factor of 2 – 10.

C. This analysis

Our goal here is to update Vecchio’s pioneering analysis by taking the precession equations and the wave phase to the next higher order and by performing a broader parameter survey (including the impact of mass ratio). By taking the precession equations to higher order, we include “spin-spin” effects — precessional effects due to one black hole’s spin interacting with gravitomagnetic fields...
from the other hole’s spin. By taking the wave phase to higher order, we include, among other terms, a time-dependent spin-spin interaction. Finally, when the mass ratio differs from one, the geodetic spin-orbit term causes the two spins to precess at different rates, even without the spin-spin corrections.

Including these effects means that we cannot model the precession with a simple, analytic rule — we are forced to integrate the equations of precession numerically as inspiral proceeds, incurring a significant performance cost. Fortunately, the basic “engine” on which this code is based runs extremely fast, thanks largely to the use of spectral integrators (which in turn is thanks to a suggestion by E. Berti 5), so total run time remains reasonable.

The cost in efficiency due to the inclusion of higher-order effects is offset by the more complete description of the signal they provide. An important consequence is that it now becomes possible from GW measurements to determine the spin of each member of the binary. With Vecchio’s approximations, only three components of the black holes’ vector spins can be determined — enough to constrain, but not determine, their spin magnitudes. Our more general approach allows us to measure all six vector spin components. To our knowledge, this is the first analysis indicating how well spin can be measured from merging comparable mass binary systems. (As Barack and Cutler have shown, spin is very well determined by measurements of GWs from extreme mass ratio binaries — those in which the system’s mass ratio \( m_2/m_1 \lesssim 10^{-4} \) or so.)

Our error estimates are computed using the maximum likelihood formalism first introduced in the context of GW measurements by Finn 28. A potential worry is that we are using a Gaussian approximation to the likelihood function. This approximation is very convenient since it allows us to directly compute a Fisher information matrix. Its inverse is the covariance matrix, which directly encodes the estimated 1-\( \sigma \) errors in measured parameters, as well as correlations among different parameters. The Gaussian approximation is known to be accurate when the SNR is “high enough” 28 52.

Unfortunately, it is not particularly obvious what “high enough” really means. In our case, we are estimating measurement errors on 15 parameters5 — a rather fearsome number to fit. The Gaussian approximation almost certainly underestimates measurement error, since it assumes the likelihood function is completely determined by its curvature in the vicinity of a maximum, missing the possibility of a long tail to large error. We thus fear that our estimates are likely to be optimistic, especially for events with relatively small SNR. It would be quite salubrious to “spot check” a few cases by directly computing the likelihood function in a few important corners of parameter space and comparing to the Gaussian predictions. This would both quantify the degree to which our calculations are too optimistic and help to determine how large SNR must be for this approximation to be reliable.

In addition to concerns about the Gaussian approximation, it must be noted that the waveform family we use for our analysis is somewhat limited. We use a post-Newtonian description of the GWs from these binaries. Since our analysis requires us to follow these binaries deep into the strong field where the usual post-Newtonian expansion is likely to be somewhat unreliable, it is likely that we are introducing some systematic error. In particular, the equations of spin precession that we are using are only given to the leading order needed to see spin-orbit and spin-spin precession effects 29. Higher spin-orbit corrections to the equations of motion and precession have recently been derived 30, which have their impact on the the waves’ phasing 31. Another analysis 32 has worked out higher order spin-spin corrections to the post-Newtonian metric, from which it would not be too difficult to work out equations of motion and precession and thence the modification to the waves’ phase. It would be interesting to see what effect the higher order corrections have on these results.

Finally, it should be noted that the frequency domain expression of the signal which we are using is derived formally using a “stationary phase” approximation. This approximation is based on the idea that the binary’s orbital frequency is changing “slowly.” The orbital frequency is thus well-defined over “short” timescales. Quantitatively, this amounts to a requirement that the timescale on which radiation reaction changes the orbital frequency, \( T_{\text{insp}} \), be much longer than an orbital period, \( T_{\text{orb}} \). Precession introduces a new timescale, \( T_{\text{prec}} \), the time it takes for the angular momentum vectors to significantly change their orientations. For the stationary phase approximation to be accurate, we must in addition require \( T_{\text{prec}} \gg T_{\text{orb}} \), a somewhat more stringent requirement than \( T_{\text{insp}} \gg T_{\text{orb}} \). No doubt, a certain amount of error is introduced due to the breakdown of this condition late in the inspiral.

Thus, the results which we present here should be taken as indicative of how well LISA is likely to be able to measure the parameters of massive black hole binaries, but cannot be considered definitive. We are confident however that the improvement in measurement accuracy obtained by taking spin precession into account is robust. Specifically, we see that errors in masses are reduced dramatically, from one to several orders of magnitude. Errors in sky position and distance are also reduced, but by a smaller factor. Such improvement may nonetheless critically improve the ability of LISA to interface with electromagnetic observatories 33 64 35. Finally, the added information in the precession signal allows us to

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5 2 masses; 2 angles specifying the initial orientation of the binary’s orbit; 4 angles specifying the initial orientation of the spins; 2 spin magnitudes; the time at which coalescence occurs; the phase at coalescence; 2 angles specifying the binary’s position on the sky; and the distance to the binary.
measure the spins of the holes. These improvements due to precession will certainly survive and play an important role even in an analysis which addresses the caveats we list above.

D. Organization of this paper

The remainder of the paper is organized as follows. In Sec. II we discuss the gravitational waveform generated by binary black hole coalescence, focusing upon the slow, adiabatic inspiral. Section II.A describes the “intrinsic” waveform produced by the motion of the orbiting black holes as given in the “restricted post-Newtonian expansion” of general relativity. Section II.B then describes the post-Newtonian precession equations which we use to model the evolution of the spins of a binary’s members, as well as how those precessions influence the waveform. Finally, in Sec. II.C we describe “extrinsic” effects which enter the measured waveform through its measurement by the LISA constellation.

In Sec. III we summarize our parameter estimation formalism; this section will be largely review to readers familiar with the literature on GW measurements. Section III.A first summarizes the maximum likelihood formalism we use to estimate measurement errors. In Sec. III.B we then describe the up-to-date model for the noise which we expect to accompany LISA measurements.

Section IV presents our results. After describing some critical procedural issues in the setup of our calculations in Sec. IV.A we summarize our results for parameters “intrinsic” to the binary (particularly masses and spins) in Sec. IV.B and for “extrinsic” parameters (particularly sky position and luminosity distance) in Sec. IV.C. In both cases, we compare when appropriate to results from a code which does not incorporate spin precession physics. (This code was originally developed for the analysis presented in Ref. [41].) The general rule of thumb we find is that the accuracy with which masses can be determined is improved by about one to several orders of magnitude when precession physics is taken into account. In addition, we find that for low redshift (\(z \sim 1\)) binaries LISA should be able to determine the spins of the constituent black holes with a relative precision of 0.1–10\%, depending (rather strongly) on the spin value. Likewise, we find improvement in the measurement accuracy of extrinsic parameters, though not quite as striking — half an order of magnitude improvement in source localization and distance determination is a good, rough rule of thumb.

An important consequence of these improvements is that LISA should be able to localize low redshift binaries — using GW measurements alone — to an elliptical “pixel!” that is perhaps 10 – a few x 10 arcminutes across in its widest direction and about a factor of 2 smaller along its minor axis. For higher redshift binaries (\(z \sim 3 – 5\)), this pixel is several times larger, perhaps a degree to a few degrees in the long direction and tens of arcminutes to a degree or two in the narrow one. These results suggest that it should not be too arduous a task to search for electromagnetic counterparts to a merging binary black hole’s GW signal — particularly at low redshift, these pixel sizes are comparable to the field of view of planned large scale surveys.

Concluding and summarizing discussion is given in Sec. V. Along with summarizing our major results and findings, we discuss future work which could allow us to quantitatively assess the consequences of some of the simplifying assumptions we have made.

At several points in this analysis, we need to convert between a source’s redshift \(z\) and luminosity distance \(D_L\). To make this conversion, we assume a flat cosmology \((\Omega_{\text{total}} = 1)\) with contributions from matter \((\Omega_M = 0.25)\) and from a cosmological constant \((\Omega_\Lambda = 0.75)\). We also choose a Hubble constant \(H_0 = 75\, \text{km s}^{-1} \text{Mpc}^{-1}\). These choices are in concordance with the latest fits presented by the WMAP team in their 3-year analysis of the cosmic microwave background [36]. The luminosity distance as a function of redshift is then given by

\[
D_L(z) = \frac{(1+z)c}{H_0} \int_0^z \frac{dz'}{\sqrt{\Omega_M(1+z')^3 + \Omega_\Lambda}}. \tag{1.4}
\]

II. GRAVITATIONAL WAVES FROM BINARY BLACK HOLE INSPIRAL

The coalescence of a black hole binary can be divided into three stages: (1) an adiabatic inspiral, (2) a merger, and (3) a ringdown, when the resulting black hole settles down to its final state. In this paper, we will focus on the inspiral. Ringdown waves have been analyzed in other work [2, 28, 37, 38]: the most comprehensive recent analysis was performed by Berti, Cardoso, and Will [39]. The merger waveform, describing the strong field and (potentially) violent process of the two black holes merging into a single body, has historically been poorly understood. Recent breakthroughs in numerical relativity may soon correct this [40, 41, 42].

The inspiral waveform which will be measured by LISA is a combination of the intrinsic waveform created by the source and extrinsic features related to its location on the sky and modulation effects caused by the motion of the detector. In this section we review the relevant physics involved in the construction of the waveform.

For sources at cosmological distances, all time scales redshift by a factor \(1+z\). In the \(G = c = 1\) units that we use, all factors of mass enter as timescales; thus, masses are redshifted by this \(1+z\) factor. [Likewise, quantities such as spin which have dimension \((\text{time})^2\) acquire a factor \((1+z)^2\), etc.] In the equations written below, we do not explicitly write out these redshift factors; they should be taken to be implicit in all our equations. When discussing results, we will always quote masses as they...
would be measured in the rest frame of the source, with redshift given separately.

A. Intrinsic waveform

We treat the members of our binary as moving on quasicircular orbits; eccentricity is very rapidly bled away by gravitational radiation reaction \[43\], so it is expected that these binaries will have essentially zero eccentricity by the time they enter LISA’s frequency band (at least at the mass ratios we consider in this paper, \(1 \leq m_1/m_2 \leq 10\)). We use the post-Newtonian formalism, an expansion in internal gravitational potential \(U\) and internal source velocity \(v\), to build our waveforms. A detailed review of the post-Newtonian formalism can be found in the article by Blanchet \[44\]; the key pieces which we will use can be found in Refs. \[44, 45, 46, 47, 48\].

The post-Newtonian equations of motion, taken to second post-Newtonian (2PN) order, yield the following generalization of Kepler’s third law relating orbital angular frequency \(\Omega\) and orbital radius (in harmonic coordinates) \(r\) \[10\]:

\[
\Omega^2 = \frac{M}{r^3} \left[ 1 - (3 - \eta) \left( \frac{M}{r} \right)^3 - \sum_{i=1}^2 \left( \frac{2m_i^2}{M^2} + 3\eta \right) \hat{L}_i \cdot S_i \left( \frac{M}{r} \right)^{3/2} + \left( 6 + \frac{41}{4} \eta + \eta^2 \right) + \frac{3}{2} \frac{\eta}{m_1^2m_2^2} \left( S_1 \cdot S_2 - 3(\hat{L}_i \cdot S_i)(\hat{L}_i \cdot S_2) \right) \left( \frac{M}{r} \right)^2 \right].
\] (2.1)

Here \(M = m_1 + m_2\) is the total mass of the system, and \(\eta = \mu/M\), where \(\mu = m_1m_2/M\) is the reduced mass. \(\hat{L}_i\) is the direction of the orbital angular momentum, and \(S_i\) is the spin angular momentum of black hole \(i\). The magnitude of the spin can be expressed as \(S_i = \chi_i m_i^2\), where \(0 \leq \chi_i \leq 1\). The leading term is the standard result from Newtonian gravity. The \(O(M/r)\) term is the first post-Newtonian correction; this is the same physics that, in solar system dynamics, causes the precession of the perihelion of Mercury. The \(O((M/r)^{3/2})\) term contains spin-orbit corrections to the equation of motion. Finally, the \(O((M/r)^2)\) term is a 2PN correction, which also includes spin-spin terms. From the equations of motion, the orbital energy of the binary \(E\) can also be computed; see \[10\].

The waveform that we will use is the “restricted” 2PN waveform. This approximation can be understood by writing the waveform (somewhat schematically) as \[8\]

\[
h(t) = \text{Re} \left( \sum_{x,m} h^x_{m}(t) e^{im\Phi_{orb}(t)} \right),
\] (2.2)

where \(x\) labels PN order, \(m\) is a harmonic index, and \(\Phi_{orb}(t) = \int \Omega(t) dt\) is orbital phase. In the restricted post-Newtonian waveform, we throw out all amplitude terms except \(h_0^3\) (the “Newtonian quadrupole” term) but compute \(\Phi_{orb}(t)\) to some specified PN order. The restricted PN approximation is motivated by the fact that matched filtering — matching a signal in noisy data by cross-correlating with a theoretical template — is much more sensitive to phase information than to the amplitude. Since the \(h_0^3\) harmonic contributes most strongly to the waveform over most of the inspiral, the restricted PN approximation is expected to capture the most important portion of the inspiral waveform. It is worth noting, however, that there is additional information encoded by the harmonics that we are neglecting. Especially for the SNRs expected for typical LISA binary black hole measurements, this additional information could play an important role in measuring source characteristics, as pointed out by Hellings and Moore \[42\].

At any rate, within the restricted PN approximation, the waveform can be written

\[
h_{ij}(t, \mathbf{x}) = -\frac{4M^{5/3}(\pi f)^{2/3}}{|\mathbf{x}|} \begin{bmatrix} \cos \Phi(t) & \sin \Phi(t) & 0 \\ \sin \Phi(t) & -\cos \Phi(t) & 0 \\ 0 & 0 & 0 \end{bmatrix},
\] (2.3)

where \(|\mathbf{x}|\) is the distance to the binary, \(M = \mu^{3/5}M^{2/5}\) is the “chirp mass” (so called because it largely determines the rate at which the system’s frequency evolves, or “chirps”), \(f = \Omega/\pi = 2f_{\text{orb}}\) is the GW frequency, and \(\Phi(t) = \int \Omega(t) dt\) is the GW phase. We have chosen a coordinate system oriented such that the binary’s orbit lies within the \(x-y\) plane; this tensor will later be projected onto polarization basis tensors to construct the measured polarizations \(h_+\) and \(h_\times\).

The rate at which the frequency changes due to the emission of gravitational radiation is given by \[10\]

\[
\frac{df}{dt} = \frac{96}{5\pi M^2} (\pi Mf)^{11/3} \left[ 1 - \left( \frac{743}{336} + \frac{11}{4} \eta \right) (\pi Mf)^{2/3} + (4\pi - \beta)(\pi Mf) + \left( \frac{34103}{18144} + \frac{13661}{2016} \eta + \frac{59}{18} \eta^2 + \sigma \right) (\pi Mf)^{4/3} \right].
\] (2.4)

Notice that the chirp mass \(M\) dominates the rate of change of \(f\); the reduced mass \(\mu\) and parameters \(\beta\) and \(\sigma\) have an influence as well. The parameter \(\beta\) describes spin-orbit interactions and is given by

\[
\beta = \frac{1}{12} \sum_{i=1}^2 \left[ 113 \left( \frac{m_i}{M} \right)^2 + 75 \frac{\mu}{M} \right] \frac{\hat{L}_i \cdot S_i}{m_i^2}.
\] (2.5)

The parameter \(\sigma\) describes spin-spin interactions:

\[
\sigma = \frac{\mu}{48M(m_1^2m_2^2)} \left[ 21(\hat{L}_1 \cdot S_1)(\hat{L}_2 \cdot S_2) - 247(S_1 \cdot S_2) \right].
\] (2.6)
Notice that $\beta$ and $\sigma$ depend on the angles between the binary’s angular momentum and the two spins. In previous analyses which have neglected precession, $\beta$ and $\sigma$ are constants; precession makes them time-dependent.

Using Eq. 2.4, we can now integrate to find
\[
t(f) = t_c - \frac{5}{256} M(\pi M f)^{-8/3} \\
\times \left[ 1 + \frac{4}{3} \left( \frac{743}{336} + \frac{11}{4} \eta \right) (\pi M f)^{2/3} \\
- \frac{8}{3} (4\pi - \beta)(\pi M f) \\
+ 2 \left( \frac{3058673}{1010664} + \frac{5429}{1008} \eta + \frac{617}{144} \eta^2 - \sigma \right) (\pi M f)^{1/3} \right].
\]

(2.7)
The parameter $t_c$ formally defines the time at which $f$ diverges within the post-Newtonian framework. In reality, we expect finite size effects to significantly modify the binary’s evolution as the members come into contact. The system evolves so quickly as the bodies come together that $t_c$ is nonetheless a useful surrogate for a “time of coalescence”. Finally, the wave phase $\Phi(t) = \int^t 2\pi f(t') dt'$ as a function of $f$ is given by
\[
\Phi(f) = \Phi(t(f)) = \Phi_c - \frac{1}{16} (\pi M f)^{-5/3} \\
\times \left[ 1 + \frac{5}{3} \left( \frac{743}{336} + \frac{11}{4} \eta \right) (\pi M f)^{2/3} \\
- \frac{5}{2} (4\pi - \beta)(\pi M f) \\
+ 5 \left( \frac{3058673}{1010664} + \frac{5429}{1008} \eta + \frac{617}{144} \eta^2 - \sigma \right) (\pi M f)^{1/3} \right],
\]

(2.8)
where $\Phi_c$ is the phase at time $t_c$. The restricted PN waveform is then constructed by inserting (2.8) into (2.3).

B. Precession equations

We next examine the effects of precession on the binary system. As discussed in the Introduction, spin-orbit and spin-spin interactions cause the black hole spins $S_1$ and $S_2$ to precess. Precession occurs, at leading order\(^6\), on a time scale $T_{\text{prec}} \propto r^{5/2}$ at large separations \[51\]. Since this is smaller than the inspiral time scale $T_{\text{insp}}$, we treat the total angular momentum $J = L + S_1 + S_2$ as constant over $T_{\text{prec}}$. The orbital angular momentum $L$ must then precess to compensate for changes in $S_1$ and $S_2$. Since $T_{\text{prec}}$ is longer than the orbital time scale $T_{\text{orb}}$, we use an orbit-averaged version of the precession equations \[29\]:
\[
\dot{S}_1 = \frac{1}{r^3} \left[ \left( 2 + \frac{3}{2} m_2 \right) \mu \sqrt{M r} \hat{L} \right] \times S_1 \\
+ \frac{1}{r^3} \left[ \frac{1}{2} S_2 - \frac{3}{2} (S_2 \cdot \hat{L}) \hat{L} \right] \times S_1,
\]

(2.9)
\[
\dot{S}_2 = \frac{1}{r^3} \left[ \left( 2 + \frac{3}{2} m_1 \right) \mu \sqrt{M r} \hat{L} \right] \times S_2 \\
+ \frac{1}{r^3} \left[ \frac{1}{2} S_1 - \frac{3}{2} (S_1 \cdot \hat{L}) \hat{L} \right] \times S_2,
\]

(2.10)
where\(^7 r = M^{1/3}/(\pi f)^{2/3}\). These equations each have two pieces \[12\]. Consider the equation for $\dot{S}_1$. The first piece, which contains no $S_2$ dependence, is the spin-orbit term. This term, which comes in at 1PN order, is due to the geodetic precession of $S_1$ as hole 1 orbits in the spacetime generated by the mass of hole 2 and to the Lense-Thirring precession of $S_1$ in the gravitomagnetic field generated by the orbital motion of hole 2. The second piece is the spin-spin term, which enters at 1.5PN order. This term can be understood as the Lense-Thirring precession of $S_1$ in the gravitomagnetic field generated by the spin of hole 2. Note that the magnitudes of the spins do not change at this order; see \[29\] for more details. From conservation of total angular momentum on short time scales, we have
\[
\hat{L} = -(\dot{S}_1 + \dot{S}_2).
\]

(2.11)
Over longer time scales, we must also consider the change in total angular momentum due to the radiation reaction, which is given by
\[
\dot{J} = -\frac{32 \mu^2}{5} \frac{M}{r} \left( \frac{M}{r} \right)^{5/2} \hat{L}
\]

(2.12)
to lowest order.

Considering only the spin-orbit terms and taking the limit $S_2 = 0$ or $m_1 = m_2$ leads to a system whose precession can described analytically; this is the “simple precession” limit described in \[29\]. Simple precession can be visualized as a rotation of $L$ and $S = S_1 + S_2$ around the total angular momentum $J$. (Since inspiral shrinks $J$, the precession is actually around a slightly different direction $J_0$; see \[29\] for further discussion.)

Since Vecchio restricts his analysis to $m_1 = m_2$ and does not include the spin-spin interaction, this limit is appropriate for his work \[29\]. As a consequence, Vecchio takes the quantities $\hat{L} \cdot \dot{S}_1$, $\dot{S}_1 \cdot \hat{S}_2$, and $\beta$ to be constant.

---

\(^6\) Several effects are built in to the precession equations discussed below, leading to precessions that occur on timescales scaling as $r^{5/2}$ and $r^3$. Since we integrate these equations numerically, all of these effects are included in our analysis. For the purposes of this discussion, we subsume these effects into the leading order timescale $T_{\text{prec}}$ discussed here.

\(^7\) We use only the lowest order Newtonian orbital quantities in these equations. Including more terms would introduce higher order effects into the precession.
(He does not include the spin-spin term in the analysis.) Here, we will study the impact of the full (albeit orbit-averaged) precession equations, including spin-spin terms, and include the impact of mass ratio. An analytic description is not possible in this case, so we must integrate these equations numerically. The behavior is qualitatively similar to the simple precession case, but with significant quantitative differences. For example, \( \beta \) now oscillates around an average value. For unequal masses (say \( m_1/m_2 \geq 2 \)), the difference due to precession can be substantial. Such cases are also astrophysically the most interesting — a mass ratio of roughly 10 is favored in binary black hole formation scenarios arising from hierarchical structure formation.

The precession of the orbital plane causes a change in the orbital phase \( \Phi_{\text{orb}}(t) \) 29. Multiplying by a factor of 2, the change in the wave phase is

\[
\delta_p \Phi(t) = - \int_t^{t'c} \delta_p \dot{\Phi}(t')dt',
\]

where

\[
\delta_p \dot{\Phi}(t) = \frac{2 \hat{L} \cdot \hat{n}}{1 - (\hat{L} \cdot \hat{n})^2} (\hat{L} \times \hat{n}) \cdot \hat{L}.
\]

At this point, we note that precession has several effects on the observed waveform. It changes the orbital phase, even at Newtonian order, by the amount 2.13, and it modifies the functions \( \beta \) and \( \sigma \) which appear in the post-Newtonian phase 2.3 and time-frequency relation 2.7. In the next section, we consider extrinsic effects on the waveform and find that precession of the orbital plane modifies them as well.

### C. Extrinsic effects

We have now constructed the intrinsic GWs emitted by a precessing binary in the restricted post-Newtonian approximation. The waveform measured by LISA will also include extrinsic effects due to the binary’s location on the sky and the motion of the detector.

We can write the wave as a combination of two orthogonal polarizations propagating in the \( -\hat{n} \) direction, where \( \hat{n} \) is the direction of the binary in the sky. Define \( \hat{p} \) and \( \hat{q} \) as axes orthogonal to \( \hat{n} \), with \( \hat{p} = \hat{n} \times \hat{L}/|\hat{n} \times \hat{L}| \) and \( \hat{q} = \hat{p} \times \hat{n} \). These are the principal axes for the wave; that is, they are defined so that the two polarizations are exactly 90° out of phase. The polarization basis tensors for these axes are \( H_{ij}^+ = p_i p_j - q_i q_j \) and \( H_{ij}^\times = p_i q_j + q_i p_j \):

\[
h_{ij}(t) = h_+(t)H_{ij}^+ + h_\times(t)H_{ij}^\times,
\]

where

\[
h_+(t) = \frac{2 M^{5/3} (\pi f)^{2/3}}{D_L} \left[ 1 + (\hat{L} \cdot \hat{n})^2 \cos[\Phi(t) + \delta_p \Phi(t)] \right],
\]

\[
h_\times(t) = -4 \frac{M^{5/3} (\pi f)^{2/3}}{D_L} (\hat{L} \cdot \hat{n}) \sin[\Phi(t) + \delta_p \Phi(t)].
\]

These expressions were first discussed in the Introduction (albeit without the precessional phase correction). Here \( D_L \) is the luminosity distance to the source. Notice that the weighting of the two polarizations depends upon the direction of the angular momentum vector relative to the sky position.

We now consider the GW as measured by the detector. All of this analysis is done using the long wavelength (\( \lambda \gg L \), where \( L \) is the LISA armlength) approximation introduced by Cutler 3; more details can be found there. This approximation is appropriate for our purposes since most of the signal accumulates at low frequencies where the wavelength is in fact greater than the armlength. The full LISA response function, including armlength effects, is discussed in 32, 33.

LISA consists of three spacecraft arranged in an equilateral triangle, 5 \times 10^6 km apart. The center of mass of the configuration orbits the sun 20° behind the Earth. The triangle is oriented at 60° to the ecliptic, so the orbits of the individual spacecraft will all be in different planes. This causes the triangle to spin around itself as it orbits the sun. Following Cutler, we define a barred “barycenter” coordinate system \( (\bar{x}, \bar{y}, \bar{z}) \), which is fixed in space with the \( \bar{x}\bar{y} \)-plane aligned with the ecliptic, and an unbarred “detector” coordinate system \( (x, y, z) \), which is attached to the detector. The \( z \) axis always points toward the sun, 60° away from vertical, while the \( x \) and \( y \) axes pinwheel around it. A particular binary will have fixed coordinates in the barycenter system, but its detector coordinates will be time-varying.

The three arms act as a pair of two-arm detectors. We are first interested in the strain measured in detector 1, that formed by arms 1 and 2:

\[
h_I(t) = \frac{\delta L_1(t) - \delta L_2(t)}{L},
\]

where \( \delta L_1(t) \) and \( \delta L_2(t) \) are the differences in length in arms 1 and 2 as the wave passes. \( L \) is the unperturbed length of the arms. Using the geometry of the detector and the equation of geodesic deviation 34, we find

\[
h_I(t) = \sqrt{3} \frac{1}{2} \left[ h_{xx} - h_{yy} \right].
\]

To obtain \( h_{xx} \) and \( h_{yy} \) for use in these equations, we must rotate the waveform from the principal axes into the detector frame. The result is that detector I measures both polarizations, modulated by the antenna pattern of that detector:
Detector I acts like a “standard” 90° GW interferometer (e.g. LIGO), with the response scaled by $\sqrt{3}/2$ (due to the 60° opening angle of the constellation). The antenna pattern functions are given by:

\[
F_I^+(\theta_N, \phi_N, \psi_N) = \frac{1}{2} \left[ (1 + \cos^2 \theta_N) \cos 2\phi_N \cos 2\psi_N - \cos \theta_N \sin 2\phi_N \sin 2\psi_N \right],
\]

(2.21)

\[
F_I^\times(\theta_N, \phi_N, \psi_N) = \frac{1}{2} \left[ (1 + \cos^2 \theta_N) \cos 2\phi_N \sin 2\psi_N + \cos \theta_N \sin 2\phi_N \cos 2\psi_N \right],
\]

(2.22)

where $\theta_N$ and $\phi_N$ are the spherical angles for the binary’s direction $\hat{n}$ in the (unbarred) detector frame and $\psi_N$ is the polarization angle of the wave in that frame:

\[
\tan \psi_N = \frac{\hat{q} \cdot \hat{z}}{\hat{p} \cdot \hat{z}} = \frac{\hat{L} \cdot \hat{z} - (\hat{L} \cdot \hat{n})(\hat{z} \cdot \hat{n})}{\hat{n} \cdot (\hat{L} \times \hat{z})}.
\]

(2.23)

In order to use these expressions, we must relate the time-dependent angles in the unbarred detector frame, $\theta_N, \phi_N,$ and $\psi_N,$ to quantities in the barred barycenter frame. Again, the details can be found in Cutler [5].

Similar expressions hold for detector II. Following Cutler, we construct the signal from detector II as

\[
h_{II}(t) = \frac{1}{\sqrt{3}} \left[ h_I(t) + 2h_{II'}(t) \right],
\]

(2.24)

where $h_I$ is the signal from detector I (2.18), and $h_{II'} = (\delta L_2(t) - \delta L_3(t))/L$ is the signal formed from the difference in the lengths of arms 2 and 3. This choice makes the noise in detector I uncorrelated with the noise in detector II; we will exploit this property in Sec. III to treat detectors I and II as independent detectors. From (2.22), we obtain

\[
h_{II}(t) = \frac{\sqrt{3}}{2} \left[ \frac{1}{2} (h_{xy} + h_{yx}) \right],
\]

(2.25)

The result is that detector II also behaves like a 90° interferometer (scaled by $\sqrt{3}/2$), but rotated by 45° with respect to detector I. Thus the antenna patterns for detector II are

\[
F_{II}^+(\theta_N, \phi_N, \psi_N) = F_I^+(\theta_N, \phi_N - \pi/4, \psi_N),
\]

(2.26)

\[
F_{II}^\times(\theta_N, \phi_N, \psi_N) = F_I^\times(\theta_N, \phi_N - \pi/4, \psi_N).
\]

(2.27)

We now rewrite the waveform in terms of an amplitude and phase. Letting $i = 1, 2$ label detector number, the waveform as measured by detector $i$ is

\[
h_i(t) = 2\sqrt{3} M^{5/3} (\pi f)^{2/3} A_{pol,i}(t) \cos[\Phi(t) + \delta_p \Phi(t)]
\]

\[
+ \varphi_{pol,i}(t) + \varphi_D(t) + \delta_p \Phi(t),
\]

(2.28)

where

\[
A_{pol,i}(t) = \frac{\sqrt{3}}{2} \left[ (1 + (\hat{L} \cdot \hat{n})^2) F_I^+(t) + 4(\hat{L} \cdot \hat{n}) F_I^\times(t) \right]^{1/2}
\]

(2.29)

is the “polarization amplitude” and

\[
\varphi_{pol,i}(t) = \tan^{-1} \left[ \frac{2(\hat{L} \cdot \hat{n}) F_I^\times(t)}{1 + (\hat{L} \cdot \hat{n})^2 F_I^+(t)} \right]
\]

(2.30)

is the “polarization phase” [5]. We have introduced the “Doppler phase” $\varphi_D(t)$, which arises from the detector’s motion around the sun and is given by

\[
\varphi_D(t) = 2\pi f(t) R_\oplus \sin \bar{\theta}_N \cos \bar{\Phi}_D(t) - \bar{\phi}_N,
\]

(2.31)

where $\bar{\Phi}_D(t)$ is the orbital phase of the detector and $R_\oplus = 1$ AU.

Much of our analysis is done in the frequency domain. We define the Fourier transform of the signal as

\[
\hat{h}(f) = \int_{-\infty}^{\infty} e^{2\pi i ft} h(t) dt.
\]

(2.32)

To evaluate the Fourier transform, we make use of the stationary phase approximation [3, 4]. This approximation relies on the fact that the orbital time scale $T_{orb}$ is much shorter than the precession time scale $T_{prec}$, as well as the inspiral time scale $T_{ins}$ and detector orbital time scale $T_D = 1$ yr. The result thus differs from the true Fourier transform by terms of order $T_{orb}/T_{prec}$ and $T_{orb}/T_{ins}$ [23]. The Fourier transform is thus likely to be inaccurate near the end of the inspiral, when all of these time scales become comparable. Using (2.27) and (2.28), we have

\[
\hat{h}_i(f) = \sqrt{\frac{5}{96}} \frac{M^{5/6}}{D_L} A_{pol,i}[\ell(f)] f^{-7/6} \times e^{i[\Psi(f) - \varphi_{pol,i}(f) - \varphi_D(f) - \delta_p \Phi(f)]},
\]

(2.33)
where the phase $\Psi(f)$ is given by

$$
\Psi(f) = 2\pi ft - \Phi_c - \frac{\pi}{4} + \frac{3}{128}(\pi Mf)^{-5/3} \\
\times \left[ 1 + \frac{20}{9} \left( \frac{743}{336} + \frac{11}{4}\eta \right) (\pi Mf)^{2/3} \\
- 4(4\pi - \beta)/(\pi Mf) + 10 \left( \frac{3058673}{1016064} + \frac{5429}{1008} \eta^2 \right) (\pi Mf)^{4/3} + \frac{617}{144}\eta^2 - \sigma \right] (\pi Mf)^{4/3}.
$$

(2.34)

In the work by Cutler, the separation of time scales that we used above leads to an interpretation of the polarization amplitude, polarization phase, and Doppler phase as modulations, in amplitude and phase, of an underlying carrier signal. These modulations make it possible to measure the sky position of the source, which also helps to measure the luminosity distance $D_L$. With the addition of precession, the polarization amplitude and polarization phase include additional modulations which further improve the measurement of these parameters. In conjunction with the purely intrinsic effects of precession, these effects also help us to better measure the masses and spins of the system.

### III. MEASUREMENT AND PARAMETER ESTIMATION WITH LISA

#### A. Theory

In the previous section, we constructed the expected form for the GW strain that LISA is being designed to measure. The signal $s_i(t)$ as measured by detector $i$ will of course also include noise $n_i(t)$:

$$
s_i(t) = h_i(t) + n_i(t). \tag{3.1}
$$

The LISA noise spectrum is discussed in section [III.B]. In this section, we discuss the theory of parameter estimation with a noisy signal. First, consider only one detector. We assume that the noise is zero mean, wide-sense stationary, and Gaussian. Wide-sense stationary means that the autocovariance function

$$
K_n(t, t') = \langle n(t)n(t') \rangle = \langle n(t) \rangle \langle n(t') \rangle \tag{3.2}
$$

depends only on the time difference $\tau = t - t'$. (Throughout this section, quantities within angle brackets are ensemble averaged with respect to the noise distribution.) A process is Gaussian if every sample of the process can be described as a Gaussian random variable and all possible sets of samples of the process are jointly Gaussian. However, the noise is colored, not white. A white noise process is defined to be a process which is uncorrelated with itself at different times; that is, its autocovariance is a delta function. Because the noise is colored, it has an interesting (nonflat) power spectral density (PSD), which is defined as the Fourier transform of the autocovariance function:

$$
S_n(f) = 2 \int_{-\infty}^{\infty} d\tau e^{2\pi i f \tau} K_n(\tau). \tag{3.3}
$$

The factor of 2 follows because we actually use the one-sided PSD. Since the noise is Gaussian, it is described entirely by its second moments. Therefore, we will only need the PSD, and not the full probability density function, to analyze the effect of the noise on the signal.

Incidentally, it can be shown that wide sense stationarity implies that the Fourier transform of $n(t)$ is a nonstationary white noise process in frequency:

$$
\langle \hat{n}(f)\hat{n}^*(f') \rangle = \frac{1}{2} \delta(f - f')S_n(f). \tag{3.4}
$$

The Fourier components are thus independent Gaussian random variables.

Now briefly consider both detectors. We explicitly constructed the second detector [23] with $h(t) \to s(t)$ so that the noise in it is uncorrelated with, and thus independent of, noise in the first detector. Thus we have

$$
\langle \hat{n}_i(f)\hat{n}_i^*(f) \rangle = \frac{1}{2} \delta_{ij} \delta(f - f')S_n(f). \tag{3.5}
$$

The uncorrelated nature of these two noises will allow us to easily generalize discussion from one detector to the full two effective detector system.

Let us write our generalized GW as $h(\theta)$, where the components of the vector $\theta$ represent the various parameters on which the waveform depends. We now assume that a GW signal with particular parameters $\tilde{\theta}$ is present in the data (i.e., “detection” has already occurred), and want to obtain estimates $\hat{\theta}$ of those source parameters. Finn shows that the probability for the noise to have some realization $n_0(t)$ is given by

$$
p(n = n_0) \propto e^{-(n_0|n_0)/2}, \tag{3.6}
$$

where the inner product used here is given by

$$
(a|b) = \Re \int_0^\infty df \frac{\hat{a}^*(f)b(f)}{S_n(f)} \tag{3.7}
$$

(3.7)

$$
= 2 \int_0^\infty df \frac{\hat{a}^*(f)b(f) + \hat{a}(f)\hat{b}^*(f)}{S_n(f)}. \tag{3.8}
$$

This product is a natural one for the vector space of (frequency domain) signals $a(f)$. (Note that this definition of the inner product differs from [23] by a factor of 2.)

Given a particular measured signal $s(t)$, the probability that the GW parameters are given by $\hat{\theta}$ is the same as the probability that the noise takes the realization $s - h(\theta)$:

$$
p(\tilde{\theta}|s) \propto e^{-(h(\tilde{\theta}) - s|h(\tilde{\theta}) - s)/2}, \tag{3.9}
$$

(3.9)
where the constant of proportionality may include prior probability densities for the parameters $\theta$. For simplicity, we take these to be uniform.

We can estimate the parameters $\hat{\theta}$ by the maximum likelihood (ML) method. This method involves finding the parameters $\hat{\theta}$ that maximize $\mathcal{L}$, or alternatively, minimize $(h(\hat{\theta}) - s|h(\hat{\theta}) - s)$, which can be considered a distance in signal space. A bank of template waveforms are correlated with the received signal and, assuming that any template produces a statistically significant correlation, the one with the highest correlation is the one with the ML parameters. The signal-to-noise ratio (SNR) for this signal is then given by $\rho$

$$\rho \approx (h(\hat{\theta})|h(\hat{\theta}))^{1/2} \approx (h(\hat{\theta})|h(\hat{\theta}))^{1/2}.$$  

(3.10)

To quantify the errors in the ML estimate, we expand around the most likely values $\hat{\theta}$. We can then write the probability density as $\mathcal{L}$:

$$p(\theta|s) \propto e^{-\Gamma_{ab}\delta\theta^a\delta\theta^b/2},$$  

(3.11)

where $\delta\theta^a = \hat{\theta}^a - \hat{\theta}^a$ and

$$\Gamma_{ab} = \left( \frac{\partial h}{\partial \theta^a} \mid \frac{\partial h}{\partial \theta^b} \right),$$  

(3.12)

evaluated at $\theta = \hat{\theta}$, is the Fisher information matrix. For small deviations from the ML estimate, the distribution is Gaussian. This expression holds for large values of the SNR $\mathcal{L}$. It is worth emphasizing at this point that in our evaluation of Eq. (3.12), most derivatives are taken numerically using finite differencing — the complicated nature of the signal (due to the inclusion of spin precession) makes it essentially impossible to evaluate all but a few of our derivatives analytically. This is another reason that the code we have developed for this analysis is substantially slower than those developed for analyses which do not include spin-precession physics.

Now we return again to the two detector case. Using $\mathcal{L}$, we can write a total Fisher matrix as the sum of the individual Fisher matrices for each detector:

$$\Gamma_{ab}^{\text{tot}} = \Gamma_{ab}^{I} + \Gamma_{ab}^{II}.$$  

(3.13)

The Fisher matrix is then inverted to produce the covariance matrix $\Sigma_{ab}^{\text{tot}} = (\Gamma_{ab}^{\text{tot}})^{-1}$. The diagonal terms of the covariance matrix represent measurement errors:

$$\Delta \theta^a \equiv \sqrt{\langle (\delta \theta^a)^2 \rangle} = \sqrt{\Sigma_{aa}}.$$  

(3.14)

The off-diagonal terms can be expressed as correlation coefficients, ranging from $-1$ to $1$:

$$c_{ab} \equiv \frac{\langle \delta \theta^a \delta \theta^b \rangle}{\Delta \theta^a \Delta \theta^b} \equiv \frac{\Sigma_{ab}}{\sqrt{\Sigma_{aa}\Sigma_{bb}}}. $$  

(3.15)

B. LISA detector and astrophysical noise

We turn now to a discussion of the noise we expect in LISA measurements. Our model for the instrumental noise spectrum, $S_{h}^{\text{inst}}(f)$, is based on that described in Ref. 52. (From now on, we use the notation $S_{h}$ for strain noise instead of $S_{n}$ for general noise.) In particular, we use the online sensitivity curve generator provided by S. Larson, which implements the recipe of 53 (see 19). The output of Larson’s webtool gives a sky averaged amplitude sensitivity curve, $h_{\text{Larson}}$. To convert to the noise we need for our analysis, we square this amplitude and insert two numerical factors:

$$S_{h}^{\text{inst}}(f) = \frac{1}{5} \times \left( \frac{\sqrt{3}}{2} h_{\text{Larson}} \right)^2 = \frac{3}{20} h_{\text{Larson}}^2.$$  

(3.16)

The factor of $1/5$ correctly accounts for the averaging of the antenna pattern functions over all sky positions and source orientations. This factor is only correct for measuring radiation with wavelength $\lambda \gg L$ (where $L$ is the LISA arm length). As a consequence, our instrumental noise will be inaccurate at high frequencies. This will have little impact on our analysis since, as already argued, the signal from merging binary black holes accumulates at low frequencies.

The factor $\sqrt{3}/2$ arises due to the $60^\circ$ opening angle of the interferometer arms; we have already accounted for this factor in our discussion of the interferometer’s interaction with a GW [cf. Eqs. (2.19) and (2.25)]. The numerical factor $3/20$ has been the source of some confusion; Berti, Buonanno and Will very nicely straightened this out. See Sec. IIC of Ref. 7 for further discussion of these factors.

Besides purely instrumental noise, LISA data will contain “noise” from a background of confused binary sources, mostly white dwarf binaries. An isotropic background of indistinguishable sources can be represented as noise with spectral density $\mathcal{L}$:

$$S_{h}^{\text{conf}}(f) = \frac{3}{5} f^{-3} \rho_c \Omega_{GW}(f),$$  

(3.17)

where $\rho_c = 3H_0^2/8\pi$ is the critical energy density to close the universe and $\Omega_{GW} = (f/\rho_c)d\rho_{GW}/df$ is the energy density in GWs relative to $\rho_c$ per logarithmic frequency interval. Using this form and the results of Farmer and Phinney 57, we model the confusion noise due to extragalactic binary sources by

$$S_{h}^{\text{exgal}}(f) = 4.2 \times 10^{-47} \left( \frac{f}{1 \text{ Hz}} \right)^{-7/3} \text{ Hz}^{-1}.$$  

(3.18)

\footnote{While surely noise when studying cosmological black holes, this background is signal to those interested in stellar populations.}
From Nelemans et al. [58], we take the galactic white
dwarf confusion noise to be
\[ S_{\text{h}}^{\text{gal}}(f) = 2.1 \times 10^{-45} \left( \frac{f}{1 \text{ Hz}} \right)^{-7/3} \text{ Hz}^{-1}. \] (3.19)

The combined instrumental and galactic confusion noise is
given by \[ S_{\text{h}}^{\text{inst+gal}}(f) = \min[S_{\text{h}}^{\text{inst}}(f)/\exp(-\kappa T^{-1}_{\text{mission}} dN/df), S_{\text{h}}^{\text{inst}}(f) + S_{\text{h}}^{\text{gal}}(f)]. \] (3.20)

The choice taken in Eq. (3.20) reflects the fact that, at
sufficiently high frequency, the number of binaries per
bin should be small enough that they are no longer truly
confused and can be subtracted from the data stream
(at least partially). The factor \( \exp(-\kappa T^{-1}_{\text{mission}} dN/df) \)
is the fraction of “uncorrupted” frequency bins. We choose
\( \kappa = 4.5 \) [59], \( T_{\text{mission}} \) is the mission duration (which we
take to be 3 years), and
\[ \frac{dN}{df} = 2 \times 10^{-3} \left( \frac{1 \text{ Hz}}{f} \right)^{11/3} \text{ Hz}^{-1} \] (3.21)
is the number density of galactic binaries per unit fre-
quency.

Finally, the total noise is given by
\[ S_{\text{h}}(f) = S_{\text{h}}^{\text{inst+gal}}(f) + S_{\text{h}}^{\text{exgal}}(f). \] (3.22)

IV. RESULTS

A. Procedure

1. Parameter space

Seventeen parameters describe the most general binary
black hole inspiral waveform [28]. Two of these are the
orbital eccentricity and the orientation of the orbital e-
ilipse; since we only consider circular orbits, we can ignore
these two. The other 15 parameters are all necessary to
describe the full post-Newtonian waveform with preces-
sion effects that we described in section 11.

We divide this set into intrinsic and extrinsic param-
eters. In our labeling system, intrinsic parameters are
those which label properties intrinsic to the binary it-
self; extrinsic parameters label properties which depend
upon the position and placement of the binary relative
to the observer. One can regard intrinsic parameters as
describing the physics or astrophysics of the binary sys-
tem, and extrinsic parameters as describing the binary’s
astronomical properties.

The intrinsic parameters we use are \( \ln m_1 \); \( \ln m_2 \);
\( \hat{\mu}_L(0) \equiv \cos \hat{\theta}_L(0) \) and \( \hat{\phi}_L(0) \), the initial direction of
the orbital angular momentum; \( \hat{\mu}_S(0) \equiv \cos \hat{\theta}_S(0) \), \( \hat{\phi}_S(0) \),
\( \hat{\mu}_S(0) \equiv \cos \theta_S(0) \), and \( \hat{\phi}_S(0) \), the initial directions of
the spins; \( \chi_1 \) and \( \chi_2 \), the dimensionless spin parameters;
t_; the time at coalescence; and \( \Phi_c \), the phase at coales-
cence. (Note that \( t_c \) and \( \Phi_c \) could very well be consid-
ered “extrinsic” in our labeling system, since they just
label the system’s state at some particular time. At any
rate, neither \( t_c \) nor \( \Phi_c \) is of much physical interest, so
their categorization isn’t too important.) Our extrinsic
parameters are \( \mu_N = \cos \theta_N \) and \( \hat{\phi}_N \), the sky position
in barycenter coordinates; and \( \ln D_L \), the luminosity dis-
tance to the binary. All of these parameters must be
fit in a measurement and thus must be included in our
Fisher matrix analysis. We are not necessarily interested
in all of them, however. In particular, we will focus on
the masses, the dimensionless spin parameters, the sky
position, and the luminosity distance.

It is worth noting that this choice of parameters is not
the same as that used in analyses which neglect preces-
sion. In that case, the direction of the angular moment-
\( \hat{L} \) is constant, and so the system’s orientation is
constant and fully described using two angles (e.g., \( \hat{\mu}_L \) and
\( \hat{\phi}_L \)). Including precession, \( \hat{L} \) is no longer constant, but
evolves according to (2.11). The solution to this differen-
tial equation requires two initial conditions, for instance
\( \hat{\mu}_L(0) \) and \( \hat{\phi}_L(0) \), which can be used as parameters of the
system. Since these initial conditions are taken at the
(somewhat arbitrary) starting point of our calculations,
they do not hold much physical interest (though they
must be fit for and thus included in our Fisher matrix).

Previous analyses, including the precursor to this work
[3], have used \( \beta \) (2.5) and \( \sigma \) (2.6) as parameters— these
are constants when precession is neglected. They are also
the only combinations of the spin magnitudes and spin
angles that enter into the expression for the waveform.
Boiling the six numbers which characterize \( S_1 \) and \( S_2 \)
down to two greatly simplifies the parameter space, but
also restricts us from being able to measure, for exam-
ple, the black holes’ spin magnitudes. When precession
is included, \( \beta \) and \( \sigma \) are no longer constants. In addi-
tion, they no longer fully characterize the signal, since
the precession equations (2.9) – (2.11) depend on all of
the components of the spins. We thus need six spin-
related parameters to fully describe the signal: the mag-
nitudes of the spins and their orientations at some initial
time. The orientations are again uninteresting, but the
fact that we can measure the magnitudes of the spins and
quantify their errors is quite interesting and new to this
analysis.

Finally, we break from tradition and use \( \ln m_1 \) and
\( \ln m_2 \) to parameterize our masses rather than \( \ln M \)
and \( \ln \mu \). The chirp mass and reduced mass have been used in
most previous work because of their appearance in the
waveform phase \( \Psi(f) \). However, the precession equa-
tions, as well as the spin parameters \( \beta \) and \( \sigma \), depend on
the individual masses of the black holes. It is a sim-
ple matter in principle to just solve for \( m_{1,2}(M, \mu) \) and
substitute into the precession equations. Unfortunately,
the Jacobian of the transformation between \( (M, \mu) \) and
\( (m_1, m_2) \) is singular when \( m_1 = m_2 \), leading to prob -lems
in evaluating the Fisher matrix.

These problems can be illustrated analytically. Consider how derivatives of some function \( f(m_1, m_2) \) with respect to \( M \) behave:

\[
\frac{\partial f}{\partial M} = \sum_{i=1}^{2} \frac{\partial f(m_1, m_2)}{\partial m_i} \frac{\partial m_i(M, \mu)}{\partial M}.
\]

When \( m_1 = m_2 \), the final derivative diverges — a behavior that we have seen numerically. The Fisher information is infinite, and the Gaussian approximation breaks down; the same problem occurs for \( \mu \). Thus, we argue that when precession is included, \( M \) and \( \mu \) are no longer a good choice of parameters to describe the system. Since we are still interested in the errors in \( \ln M \) and \( \ln \mu \) (which are determined to higher accuracy than the individual masses), we convert using the propagation of errors formulae

\[
\left( \frac{\Delta M}{M} \right)^2 = \left( \frac{m_1}{M} \right)^2 \left( \frac{\partial M}{\partial m_1} \right)^2 \left( \frac{\Delta m_1}{m_1} \right)^2 + \left( \frac{m_2}{M} \right)^2 \left( \frac{\partial M}{\partial m_2} \right)^2 \left( \frac{\Delta m_2}{m_2} \right)^2 + \left( \frac{m_1 m_2}{M^2} \right) \left( \frac{\partial M}{\partial m_1} \right) \left( \frac{\partial M}{\partial m_2} \right) \Sigma \ln m_1, \ln m_2,
\]

\[
\left( \frac{\Delta \mu}{\mu} \right)^2 = \left( \frac{m_1}{\mu} \right)^2 \left( \frac{\partial \mu}{\partial m_1} \right)^2 \left( \frac{\Delta m_1}{m_1} \right)^2 + \left( \frac{m_2}{\mu} \right)^2 \left( \frac{\partial \mu}{\partial m_2} \right)^2 \left( \frac{\Delta m_2}{m_2} \right)^2 + \left( \frac{m_1 m_2}{\mu^2} \right) \left( \frac{\partial \mu}{\partial m_1} \right) \left( \frac{\partial \mu}{\partial m_2} \right) \Sigma \ln m_1, \ln m_2.
\]

For unequal masses, we find that computing errors in \( m_1 \) and \( m_2 \) and then converting gives the same result as simply computing errors in \( M \) and \( \mu \) directly. We do not find good agreement in the equal mass case; for the reasons discussed above, however, we do not trust the \((M, \mu)\) parameterization in this case. At any rate, the case \( m_1 = m_2 \) is quite implausible in nature, so this is almost certainly a moot point as far as real measurements are concerned\(^9\). We note that Vecchio \(^24\) for simplicity considers the equal mass case exclusively but does not report any anomalous behavior such as we have seen. We are puzzled about this discrepancy.

\[\text{\(^9\) It is worth noting that even a slight mass difference (a few percent) is sufficient for the two approaches to match.}\]

2. Calculations

The code we use to calculate parameter measurement errors is based on that used in \[^3\]. It is written in C++ using several routines taken, sometimes with slight modification, from \[^4\]. As in \[^3\], we perform Monte Carlo simulations in which we specify rest-frame masses and redshift and then randomly choose sky position, initial angular momentum and spin directions, spin magnitudes, and time of coalescence within the three year LISA mission window. We specify spin magnitudes for some studies as well.

The primary function of the code is the calculation of the full gravitational waveform, including precessional effects. In order to effectively use the formulas of section \[^3\] we take the wave frequency \( f \) as the independent variable. The elapsed time is related to the frequency using (2.7). The calculation is started when the waveform enters LISA’s band (taken to be \( f_{\text{min}} = 3 \times 10^{-5} \) Hz throughout this paper) or when the LISA mission begins, whichever is later. (By treating the time of coalescence as a Monte Carlo variable, some signals will be partially cut off because they are in band when LISA begins observations.) The calculation proceeds until the binary reaches the Schwarzschild innermost stable circular orbit (ISCO) at orbital separation \( r = 6M \).

Though perhaps a somewhat crude choice, we use this criterion for simplicity. Choosing slightly different cut-off radii does not change our results very much; at any rate, the post-Newtonian phase formula and precession equations we use in this regime are unlikely to be very accurate. The frequency at this point, which we call the “merge frequency,” can be found using (2.7) for \( r = 6M \) (plus \( f = \Omega/\pi \)). Note that \( f \) depends on the directions of the angular momentum and spin vectors as a function of time; since we have not found them yet, this value can only be an estimate of where to stop. Any error due to this approximation is no doubt unimportant compared to the arbitrary selection of \( r = 6M \) for the transition.

Once the frequency range has been determined, the true work begins. We integrate the precession equations (2.10) – (2.11) using a Runge-Kutta routine to find the values of \( L, \dot{S}_1, \) and \( \ddot{S}_2 \) over the duration of the signal. The routine is a fifth-order adaptive step algorithm in the frequency domain. At each frequency, the code takes the results for the three orbital angular momentum components and six spin components and uses them to calculate \( \dot{\mu}_L, \dot{\phi}_L, \dot{\beta}, \) and \( \dot{\sigma} \). It also computes the integrated correction to the phase using the derivative (2.11).

As already discussed, our derivatives are taken numerically rather than analytically. We therefore must do the integration described above a total of 21 times: once for the given values of the parameters, and twice more for small shifts in each parameter which requires a numerical derivative. This repetition slows the code quite drastically compared to its earlier incarnation — an unfortunate but unavoidable cost.

Once all of the necessary integrations are complete,
the SNR \( 3.10 \) and the Fisher matrix \( 3.12 \) can be calculated for each of the two effective detectors of LISA using the noise \( S_n(f) \) (see Sec. 3.14). Some previous work \( 3, 7 \) investigated parameter estimation using the signal from only one synthesized detector; we will always assume that both are operational. It would be interesting to see how measurement degradation due to only having a single operating detector can be ameliorated by including precession effects.

At this stage, the necessary integrals are performed using Curtis-Clenshaw quadrature, which depends on the decomposition of the integrand into Chebyshev polynomials \( 60 \). This method keeps the code reasonably fast even with the addition of the Runge-Kutta routine. At each step of the integration, the integrator uses the values that were calculated using that Runge-Kutta routine to evaluate the waveform and/or its appropriate derivatives. The derivatives are calculated using

\[
\frac{df}{d\theta} \approx \frac{f(\theta + \Delta \theta) - f(\theta - \Delta \theta)}{2 \Delta \theta}.
\]  

(4.4)

For all parameters, we use \( \Delta \theta = 10^{-5} \theta \). We invert the Fisher matrix using LU decomposition to produce the covariance matrix \( 60 \). In “poor” cases, (e.g., high mass binaries at large redshift) the Fisher matrix can be nearly singular, with a large condition number\(^{10} \). In such a case, the covariance matrix produced by the code may not be the true inverse of the Fisher matrix (and may not even be positive definite). This problem is largely ameliorated by representing our numerical data in long double format — this improves (relative to type double) matrix inverses in many “bad” cases but leaves all other cases essentially unchanged.

It is worth noting that the “bad” cases are typically ones in which the binary executes very few orbits over the course of the measurement. We are confident in our results for all cases in which the number of measured orbits, \( N_{\text{orb}} = \Phi_{\text{orb}}/2\pi \), is greater than \( \sim 10 - 20 \). When the number of orbits is small (and the condition number is concomitantly high), the errors are so large that they are basically meaningless. In such a case, measurement would not determine the system’s characteristics in any meaningful sense.

**B. Black hole masses and spins**

Representative examples of our results are shown in Figures 2 and 3. These histograms show the spread of errors in \( M \) and \( \mu \) for a sample of \( 10^4 \) binaries at \( z = 1 \) with rest frame masses \( m_1 = 10^6 M_\odot \) and \( m_2 = 3 \times 10^5 M_\odot \). Each figure compares the results of the new code to those of the original code of \( 6 \), which neglects precession. (This code has been updated to reflect up-to-date models for LISA noise; some minor coding errors have also been corrected.) Clearly, including spin precession leads to a significant improvement in the measurement of these mass parameters. The reduced mass \( \mu \) in particular is improved. This is because the time variation of \( \beta \) and \( \sigma \) breaks a near degeneracy between those terms and \( \mu \) in the post-Newtonian phase \( 2.31 \). The masses also control the precession rate, as seen in (2.9) – (2.11). (Recall that in those equations, \( S_i = \chi_i m_i^2 \).) This means that they now influence the polarization amplitude and polarization phase; they do not influence those quantities when precession is neglected. These precession-induced influences on the waveform make it possible to determine the masses even more accurately than before.

![Distribution of errors in chirp mass \( M \) for \( 10^4 \) binaries with \( m_1 = 10^6 M_\odot \) and \( m_2 = 3 \times 10^5 M_\odot \) at \( z = 1 \). The dashed line is the precession-free calculation; the solid line includes precession. Precession reduces the measurement error by about an order of magnitude.](image)

As discussed earlier, we have found the masses \( m_1 \) and \( m_2 \) to be more useful parameters than \( M \) and \( \mu \) when precession is included. Figure 4 shows the error in measurements of the individual masses for our example system. While these masses are measured quite accurately, they are not measured as accurately as \( M \) and \( \mu \). This reflects the fact that even though the individual masses play a role in the precession, the other parts of the waveform depend explicitly on the combinations \( M \) and \( \mu \). Notice also that the smaller mass is typically determined a bit better than the larger one, though the difference is not large.

Precession makes it possible to determine the spins of the binary’s members. Figure 5 shows the error in measurements of the two dimensionless spin parameters \( \chi_1 \) and \( \chi_2 \). We see that \( \chi \) is generally determined very well: Taking a typical spin parameter to be about 0.5 (recall

\(^{10}\) The “condition number” is the ratio of the largest eigenvalue of a matrix to the smallest. A rule of thumb is that matrix inversion breaks down when the logarithm of the condition number of a matrix exceeds the number of digits of accuracy in the matrix elements (see, e.g., discussion in 60).
hole spin scales as mass squared (appears to be a simple consequence of the fact that black holes are better determined than that of the smaller hole. This dimensionless spin parameter of the larger hole tends to 

Next, we examine how well spin is measured as a function of spin magnitude. Figure 6 shows the error in spin has more of an impact on the waveform.

we randomly choose \( \chi \) between 0 and 1, the bulk of this distribution corresponds to errors of a bit less than a percent. For this entirely random distribution of \( \chi \), the dimensionless spin parameter of the larger hole tends to be better determined than that of the smaller hole. This appears to be a simple consequence of the fact that black hole spin scales as mass squared \( (S_i = \chi_i m_i^2) \), and larger spin has more of an impact on the waveform.

Next, we examine how well spin is measured as a function of spin magnitude. Figure 4 shows the error in \( \chi_1 \) for the same system as in Fig. 3, except that we set \( \chi_1 = \chi_2 = 0.9 \) (solid line) and \( \chi_1 = \chi_2 = 0.1 \) (dashed line), rather than randomly distributing their values. This allows us to more accurately assess how well spin is determined as a function of its value, as well as to more accurately determine the percent error we expect in these measurements. For \( \chi_1 = \chi_2 = 0.1 \), the error is almost 10%, while for \( \chi_1 = \chi_2 = 0.9 \), the error is closer to 0.1%. This is a considerable difference and is easily ascribed to the fact that rapid spin has a much stronger impact on the waveform.

Table I shows the median errors in intrinsic parameters for different mass ratios at \( z = 1 \). We continue to include the errors in \( M \) and \( \mu \) for comparison with the precession-free case, but only in binaries of unequal mass where the Gaussian approximation is well defined. Examining the table, we see some interesting features. The errors in general are worse for higher mass binaries, which spend less time in the LISA band. At \( m_1 = 10^5 M_\odot \) and \( m_2 = 3 \times 10^5 M_\odot \), the mass errors jump to nearly 10 percent, compared to tenths of a percent at the next lower mass combination. In addition, the spin determination becomes very unreliable. Mass ratio also has an important effect on the results. Taking into account the general trend caused by total mass, we see that unequal mass ratios generally produce better results. This is good news for eventual measurements of astrophysical systems, since merger tree calculations show that binaries are most likely to have mass ratios of about 10. To understand the mass ratio dependence, we again turn to the precession equations. For unequal mass ratios, the geodetic spin-orbit and spin-spin terms will cause the two spins to precess at different rates, creating richer features in the signal than for equal mass ratios. This illustrates the importance of effects beyond the “simple precession” of. We
C. Sky position and distance to source

We now focus on extrinsic parameters, the sky position and the luminosity distance to the source. We find that the determination of these parameters is likewise improved when precession physics is taken into account, though not as strongly as for intrinsic parameters. This might be expected, since precession is an intrinsic effect local to the binary and has no direct dependence on these extrinsic parameters. Precession’s impact on the extrinsic parameters is somewhat more indirect — it largely improves their determination by reducing the (otherwise quite strong) correlation between sky position and the orbital angular momentum direction \( \mathbf{L} \) and between these angles and the source’s luminosity distance.

In our analysis, a binary’s position on the sky is characterized by the two parameters \( \theta = (\mu_N \equiv \cos \bar{\theta}_N, \bar{\phi}_N) \). Consider the subspace containing just these two parameters. The Fisher matrix is a \( 2 \times 2 \) matrix which characterizes the probability density for the true parameters given a measured signal [see (3.11)]:

\[
p(\theta|s) \propto \exp\left(-\frac{1}{2} \Gamma_{\alpha\beta} \delta\theta^\alpha \delta\theta^\beta\right),
\]

To accurately describe the error ellipse on the sky, we need to manipulate the right hand side of (4.5) in several ways. First, we change coordinates from \( \mu_N \) to \( \bar{\theta}_N \), while leaving \( \bar{\phi}_N \) alone; schematically, we can represent this as \( \theta^a \rightarrow \bar{\theta}^a \). This transformation gives

\[
\delta\bar{\theta}_N = (d\bar{\theta}_N/d\mu_N)\delta\mu_N = -\delta\mu_N / \sin \bar{\theta}_N.
\]

Next, we need to redefine what we mean by “error” in order to make the results more relevant to observations. To do so, we define the “proper” angular errors

\[
\delta\bar{\theta}_N^p = \delta\bar{\theta}_N \quad \text{and} \quad \delta\bar{\phi}_N^p = \sin \bar{\theta}_N \delta\bar{\phi}_N.
\]

The proper angular errors are just the normal coordinate errors rescaled by the metric of the sphere to correctly account for the proper size of a segment \( \delta\bar{\theta}_N \) at different \( \bar{\theta}_N \). Substitut-
TABLE I: Median errors in intrinsic quantities for $10^4$ binaries of various masses at $z = 1$, including comparisons with the “no precession” case where possible. We have omitted the errors in chirp mass and reduced mass for equal mass binaries because that parameterization of the waveform fails the Gaussian approximation at those points.

<table>
<thead>
<tr>
<th>$m_1$ (M$_\odot$)</th>
<th>$m_2$ (M$_\odot$)</th>
<th>$\Delta m_1/m_1$</th>
<th>$\Delta m_2/m_2$</th>
<th>$\Delta \chi_1$</th>
<th>$\Delta \chi_2$</th>
<th>$\Delta M/M$ (no precession)</th>
<th>$\Delta M/M$ (precession)</th>
<th>$\Delta \mu/\mu$ (no precession)</th>
<th>$\Delta \mu/\mu$ (precession)</th>
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<td>$10^5$</td>
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<td>0.000629</td>
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<td>$10^7$</td>
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TABLE II: Median errors in intrinsic quantities for $10^4$ binaries of various masses at $z = 3$.

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<th>$\Delta m_2/m_2$</th>
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<th>$\Delta \chi_2$</th>
<th>$\Delta M/M$ (no precession)</th>
<th>$\Delta M/M$ (precession)</th>
<th>$\Delta \mu/\mu$ (no precession)</th>
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</table>

Finally, we diagonalize the Fisher matrix by rotating our parameterization, $\theta'^i \rightarrow \theta^a$, such that the probability

$$
\rho(\theta|s) \propto \exp \left[-\left(\frac{\delta \theta^b}{\sigma^b} \right)^2 + \left(\frac{\delta \theta^c}{\sigma^c} \right)^2 \right].
$$

In these coordinates, the covariance matrix is

$$
\Sigma_{ab} = \begin{bmatrix}
\sigma^2 & 0 \\
0 & \sigma^2
\end{bmatrix}.
$$

Following Cutler [2], we define the error ellipse such that the probability that the source lies outside the error ellipse is $e^{-1}$. The semiaxes of the error ellipse are given
by \( \sqrt{2(\sigma_{1,2}^P)^2} \). These quantities follow from the eigenvalues of the covariance matrix \((4.12)\); since eigenvalues are invariant under rotation, we can calculate them before performing the rotation. In terms of our original covariance matrix, the major axis \(2a\) and minor axis \(2b\) of the ellipse are given by

\[
2 \left[ \csc^2 \theta_N \Sigma_{\mu N} \bar{\phi} N + \sin^2 \theta_N \Sigma_{\phi N} \bar{\phi} N \pm \sqrt{\left(\csc^2 \theta_N \Sigma_{\mu N} \bar{\phi} N - \sin^2 \theta_N \Sigma_{\phi N} \bar{\phi} N \right)^2 + 4(\Sigma_{\mu N} \bar{\phi} N)^2} \right]^{1/2}
\]

(taking the plus and minus for major and minor axes respectively. We also find the area of the error ellipse:

\[
\Delta \Omega_N = \pi a b = 2 \pi \sqrt{\Sigma_{\mu N} \bar{\phi} N \Sigma_{\phi N} \bar{\phi} N - (\Sigma_{\mu N} \bar{\phi} N)^2} \quad (4.13)
\]

Many previous analyses have reported the ellipse’s area \(\Delta \Omega_N\) or \(\sqrt{\Delta \Omega_N}\), the side of a square of equivalent area, as the sky position error \([8, 10]\). Information about the ellipse’s shape, crucial input to coordinating GW observations with telescopes, is not included in such a measure. By examining both \(2a\) and \(2b\), this information is restored. Figure 7 shows the major axis of the error ellipse \(2a\) for both the original code, with no precession, and the code including precession effects. Figure 8 shows the same for the minor axis \(2b\). (Note that these figures cannot tell us which major axis is associated with which minor axis; that information is lost in the construction of the histograms.)

Compared to the code which does include precession physics, the peak of the major axis distribution is reduced by about an order of magnitude; the median of this distribution is reduced by about a factor of 2. For the minor axis, both the peak and median are reduced by about a factor of 2. The minor axis distribution also shows a long tail of very small errors. In those cases, the improvement is less than for the masses. For comparison purposes, we include

\[
\begin{array}{|c|c|c|c|c|c|c|c|c|}
\hline
m_1 (M_\odot) & m_2 (M_\odot) & \Delta m_1/m_1 & \Delta m_2/m_2 & \Delta \chi_1 & \Delta \chi_2 & \Delta M/M (no precession) & \Delta M/M (precession) & \Delta \mu/\mu (no precession) & \Delta \mu/\mu (precession) \\
\hline
10^7 & 10^7 & 0.00646 & 0.00645 & 0.0348 & 0.0340 & — & — & — & — \\
3 \times 10^7 & 10^7 & 0.00735 & 0.00597 & 0.0151 & 0.0312 & 0.00107 & 8.06 \times 10^{-3} & 0.178 & 0.00264 \\
3 \times 10^7 & 3 \times 10^7 & 0.0109 & 0.0109 & 0.0542 & 0.0527 & — & — & — & — \\
10^7 & 10^7 & 0.00955 & 0.00421 & 0.00701 & 0.0384 & 0.00336 & 0.000184 & 0.313 & 0.00330 \\
10^7 & 3 \times 10^7 & 0.0133 & 0.0107 & 0.0183 & 0.0387 & 0.00458 & 0.000309 & 0.491 & 0.00516 \\
10^7 & 10^7 & 0.0373 & 0.0374 & 0.181 & 0.184 & — & — & — & — \\
3 \times 10^7 & 3 \times 10^7 & 0.0141 & 0.0105 & 0.0130 & 0.0961 & 0.0122 & 0.00135 & 0.652 & 0.00830 \\
3 \times 10^7 & 10^7 & 0.0517 & 0.0426 & 0.0515 & 0.105 & 0.0113 & 0.00374 & 0.635 & 0.0193 \\
3 \times 10^7 & 3 \times 10^7 & 0.422 & 0.419 & 2.60 & 2.60 & — & — & — & — \\
10^7 & 10^7 & 0.260 & 0.315 & 0.387 & 3.41 & 0.0416 & 0.137 & 0.713 & 0.273 \\
10^7 & 3 \times 10^7 & 15.5 & 11.5 & 10.1 & 16.4 & 0.120 & 2.75 & 0.797 & 5.97 \\
10^7 & 10^7 & 57400 & 57100 & 262000 & 261000 & — & — & — & — \\
\hline
\end{array}
\]

TABLE III: Median errors in intrinsic quantities for \(10^4\) binaries of various masses at \(z = 5\). The results for the highest masses are meaningless — the parameters are completely undetermined.
results that neglect spin precession. Binaries with the best determined parameters at this redshift have total mass several $\times 10^5 M_\odot \lesssim M_{\text{tot}} \lesssim$ several $\times 10^6 M_\odot$ — smaller binaries are not quite determined so well due to the weakness of their signal, while larger ones are not determined so well because they radiate fewer cycles in band. We also see again that unequal mass binaries give better results than equal mass binaries due to the impact of mass ratio on precession effects. Overall, we find that the major axis of the error ellipse is on the order of 10 arcminutes for low mass and several tens of arcminutes for higher mass, while the minor axis is a factor of 2 smaller. This represents an improvement by a factor $\sim 2 - 6$ over the “no precession” case. The distance errors are on the order of $0.1 - 0.4\%$ for most masses, a factor of $\sim 2 - 7$ improvement.

Tables V and VI show the same results for higher redshift. We see the same trends as at $z = 1$, but with some degradation in numerical value. The sky position errors reach a degree or more in the major axis and several tens of arcminutes up to a degree or two in the minor axis. The distance errors are on the order of one to several percent for most masses. At the highest masses, we again see that these parameters are essentially undetermined and that precession makes things worse by requiring extra parameters to be fit.

V. SUMMARY AND CONCLUSIONS

The general relativistic precession of black holes in binary systems can have a strong influence on the binary’s dynamics and thus upon the GWs that it generates. It has been known for some time that it will be necessary to take these dynamics into account in order to detect these black holes in noisy detector data; clearly, taking these dynamics into account will be just as (if not more) important for the complementary problem of determining the parameters which characterize a detected system. Vecchio first demonstrated that by taking into account precession physics, quite a few near-degeneracies among binary source parameters can be broken, making our estimates for how accurately they can be determined more optimistic. This analysis largely confirms and extends Vecchio’s pioneering work. By taking the equations of motion to higher order to include spin-spin couplings, and by surveying measurement accuracy as a function of mass ratio, we have found that the improvement noted by Vecchio holds rather broadly. The degeneracy breaking due to precession physics is a rather robust phenomenon.

Two conclusions from this work are particularly important with regard to the astrophysical reach of future LISA measurements. The first is that modeling spin precession physics makes it possible to determine the magnitudes of the spins of the black holes which constitute the binary. If the spins are rapid, they can be measured quite accurately (as good as 0.1% accuracy for high spin, low redshift systems) due to the strong modulation imposed on the signal by their interaction. Coupled with the fact
<table>
<thead>
<tr>
<th>$m_1 (M_\odot)$</th>
<th>$m_2 (M_\odot)$</th>
<th>$2a \text{ (arcminutes)}$ (no precession)</th>
<th>$2b \text{ (arcminutes)}$ (no precession)</th>
<th>$\Delta D_L/D_L$ (no precession)</th>
<th>$\Delta D_L/D_L$ (precession)</th>
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<tr>
<td>$10^6$</td>
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<td>179</td>
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**TABLE IV:** Median errors in extrinsic quantities for $10^4$ binaries of various masses at $z = 1$, including comparisons with the “no precession” case. Note that the given major axis and minor axis are the medians for each data set and do not correspond to the same binary. However, they still represent an average sky position error ellipse in the following sense: $\sqrt{\Delta a \Delta b}$, calculated using the median values of $2a$ and $2b$, differs in most cases by less than 10 percent from the median value of $\sqrt{\Delta\Omega}$ calculated from the covariance matrix and (4.13).

<table>
<thead>
<tr>
<th>$m_1 (M_\odot)$</th>
<th>$m_2 (M_\odot)$</th>
<th>$2a \text{ (arcminutes)}$ (no precession)</th>
<th>$2b \text{ (arcminutes)}$ (no precession)</th>
<th>$\Delta D_L/D_L$ (no precession)</th>
<th>$\Delta D_L/D_L$ (precession)</th>
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<td>2130</td>
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</table>

**TABLE V:** Median errors in extrinsic quantities for $10^4$ binaries of various masses at $z = 3$.  

that the black hole masses can likewise be measured with good precision, this suggests that LISA will be a valuable tool for tracking the evolution of both mass and spin over cosmic time. Such observations could provide a direct window into the growth of cosmological structures. Measuring spin may also make it possible to indirectly test the black hole area theorem. The requirement that black hole area can only grow implies a consistency between the initial and final masses and spins. By measuring the initial masses and spins through the inspiral, and the mass and spin of the merged remnant hole through the ringdown waves, we can check this consistency relation in a manner analogous to the mass loss test proposed in \[38\]. We intend to investigate whether this test is feasible in future work.

Second, we confirm Vecchio’s result that precession breaks degeneracies between the angles which determine a binary’s orientation and its position on the sky, improving the accuracy with which sky position can be fixed using GWs alone. At low redshift ($z \sim 1$), we find that sources can be localized to within an ellipse whose major axis is typically $\sim 10 - \sim 10$ arcminutes across and whose minor axis is typically a factor $\sim 2$ smaller. This is small enough that searching the GW “pixel” for an electromagnetic counterpart to the merger event should not be too arduous a task \[39\]. For mergers at higher
redshift, the waves weaken and the source is not so well localized. The field which would need to be searched for sources at $z \sim 3 - 5$ is typically about a degree to a few degrees across in the long axis and tens of arcminutes to a degree or two in the short direction — a rather more difficult challenge, but not hopeless. We intend to more thoroughly investigate the nature of localization with spin precession, including how the pixel evolves with observation time up to final merger, in future work.

As mentioned in the Introduction, our analysis makes many assumptions and approximations which are likely to affect our results; a goal of future work will be to lift these approximations. One major concern is the Gaussian approximation we have taken to the likelihood function. As already discussed, this approximation is known to be good when the SNR is “large” \cite{3}; however, it is not apparent what “large” really means, particularly given that we are fitting for 15 parameters. Lifting this simplifying approximation can be done by simply computing the likelihood function \cite{3}, directly and examining how well parameters are thereby determined. In the context of GW measurements, Markov Chain-Monte Carlo techniques have been investigated and found to be very useful \cite{4, 5, 6}; the application of these techniques to LISA measurement problems is now being rather actively investigated \cite{7, 8, 9}.

Because we have used the Gaussian approximation (among other simplifications taken in this analysis), we cannot claim that this analysis gives a definitive statement about the accuracy with which LISA could measure binary black hole source parameters. However, it is certainly indicative of the accuracy which we expect LISA to achieve. In particular, we are confident that the trends we have seen as parameters are varied (e.g., masses, redshift, spin magnitude) are robust. Most importantly, it is very clear that the influence of spin-induced precession upon the measured waveform allows parameters to be measured to greater accuracy than before.

### Acknowledgments

This work benefited greatly from discussions with and feedback from members of the LISA International Science Team, especially Neil Cornish, E. Sterl Phinney, and Thomas Prince. We also thank Bence Kocsis for very useful discussions about parameterizing the sky position error ellipse, Shane Larson for providing code to process his noise files, and Jeremy Schnittman for helpful comments about the precession equations. We are especially grateful to Emanuele Berti for discussions regarding integrators, which led to a drastic improvement in the ability of our code to perform large Monte Carlo surveys, as well as many insightful comments about this manuscript. This work was supported by NASA Grant NAGW-12906 and NSF Grant PHY-0449884. SAH gratefully acknowledges the support of the MIT Class of 1956 Career Development fund.

### References
