Towards communication in a curved spacetime geometry

Qasem Exirifard, Eric Culf & Ebrahim Karimi

The current race in quantum communication - endeavours to establish a global quantum network - must account for special and general relativistic effects. The well-studied general relativistic effects include Shapiro time-delay, gravitational lensing, and frame dragging which all are due to how a mass distribution alters geodesics. Here, we report how the curvature of spacetime geometry affects the propagation of information carriers along an arbitrary geodesic. An explicit expression for the distortion onto the carrier wavefunction in terms of the Riemann curvature is obtained. Furthermore, we investigate this distortion for anti de Sitter and Schwarzschild geometries. For instance, the spacetime curvature causes a 0.10 radian phase-shift for communication between Earth and the International Space Station on a monochromatic laser beam and quadrupole astigmatism; can cause a 12.2% cross-talk between structured modes traversing through the solar system. Our finding shows that this gravitational distortion is significant, and it needs to be either pre- or post-corrected at the sender or receiver to retrieve the information.
Photons, electromagnetic waves, are widely used in classical and quantum communication since they do not possess electric charge or rest mass. However, a photon’s traits, e.g. group and phase velocity, wavelength, linear and optical angular momentum, are modified inside or during propagation through a linear or a nonlinear medium. These traits are governed by Maxwell’s equations, which are the relativistic quantum field theory of the \( U(1) \) gauge connection. Understanding how these optical properties are altered upon propagation is a key element for any optical communication network. In optical communication, the sender and the receiver, namely Alice and Bob, use one or several internal photonic degrees of freedom, such as wavelength, polarisation, transverse mode or time-bins, to share information, including a ciphertext and the secret key to decrypt the ciphertext. The propagation, e.g. through fibre, air or underwater channels, causes these photonic degrees of freedom to be altered, and thus causes undesired errors to the shared information. Therefore, the alteration to those degrees of freedom for any communication channel needs to be considered and well examined. Sharing information with a longer range or with several internal photonic degrees of freedom, such as wave-packet propagating along a geodesic. The alteration operators depend on the geodesic and components of the Riemann tensor on the geodesic. The alteration is calculated for examples including the spacetime geometry around the Earth and the Sun.

### Results and discussion

We start by considering a relativistic massless scalar field \( \psi = \psi(x^\mu) \) that propagates in a curved spacetime geometry \( x^\mu = (t, x^i) \) with an arbitrary Riemann curvature tensor \( R_{abcd} \). The units are chosen such that the speed of light in vacuum and the reduced Planck constant are equal to one, i.e. \( c = 1 \) and \( \hbar = 1 \). Let us consider a localised wavefunction (information carrier) whose size is small compared to the curvature of the spacetime geometry. At the leading order, therefore, the carrier can be treated as a massless point-like particle that travels along a null geodesic \( y \), see Fig. 1a. We choose the local Fermi coordinates\(^{20}\) along the geodesics in order to compute the quantum relativistic corrections. The metric’s components in the Fermi coordinates can be expanded in terms of the components of the Riemann tensor \( R_{\mu
u} \) and its covariant derivatives evaluated on the geodesic, see Fig. 1b,c. The expansion of the metric up to quadratic order in the transverse coordinates of the geodesic of a massless particle \( y \) is given by\(^{21}\),

\[
\begin{align*}
\text{d}^2s &= 2\text{d}x^\pm \text{d}x^- + \delta_{\mu\nu} \text{d}x^\mu \text{d}x^\nu - \left[ R_{\mu\nu\rho\sigma} x^\rho (\text{d}x^\sigma)^2 \right] \\
&\quad + \frac{4}{3} R_{+\mu\nu} x^\mu (\text{d}x^\nu)^2 + \frac{1}{3} R^{\alpha\beta\gamma\delta} x^{\alpha} (\text{d}x^{\beta})(\text{d}x^{\gamma})(\text{d}x^{\delta}) + \ldots, \\
\end{align*}
\]

where \( x^\pm = (x^3 \pm t)/\sqrt{2} \) represent the Dirac light-cone coordinates\(^{22}\) in the Fermi coordinates \( x^\mu \) is always tangent to

![Fig. 1 Schematics of communication in a general curved spacetime geometry and proper chosen coordinates.](image)

A two parties, namely Alice and Bob, communicate in a general curved spacetime geometry. Alice encodes her message in a sequence of information carriers and sends them to Bob. The information traverses through the spacetime over a geodesic \( y \). The physical traits of information are distorted by the curvature of the spacetime, causing errors to the field wavepacket. The shown null geodesic \( y \). Close to \( y \), at the local coordinates, the metric is approximately pseudo-Euclidean. The associated Fermi coordinates, the null-geodesic-path is mapped to \( x^+ \).
the null geodesic Fig. 1b), $\delta_{ab}$ is the Kronecker delta, $\dot{y}$ is the tangent of the null geodesic and $a, b \in \{1, 2\}$, and $(x^a) = (x^-, x^+)$. The curvature components are evaluated on $y$. The tree-level action of a massless scalar field $\psi$ in a general curved spacetime geometry is given by $S[\psi] = \frac{1}{2} \int d^4x \sqrt{-g} \partial_\mu \psi \partial^\mu \psi$, where $g$ is the determinant of the metric $g_{\mu\nu}$ and $\mu, \nu \in \{\pm, \pm\}$. Since the tree-level action is quadratic in terms of $\psi$, it is quantum mechanically exact, which can be verified by looking at its generating function, i.e. $Z[J] = \int \mathcal{D}\psi e^{-S[\psi]} \int d^4x \sqrt{-g} \partial_\mu \psi \partial^\mu \psi$ represents the integration over field configurations and $J$ is the source field. Its exact effective action, as defined by the Legendre transformation of $\ln Z$, coincides with the tree-level action, i.e. $\Gamma[\psi] = S[\psi]$. For a general action, $\phi$, resembles a “classical” field whose action is given by $\Gamma[\phi]$, while $\Gamma[\psi]$ encapsulates all the quantum loop corrections. The exact effective action includes both the classical and quantum effects. The classical effects are those that can be reproduced by motion of a point-like particle along the geodesic; the rest are quantum. The effective action of a free photon propagating in curved spacetime geometry coincides with the tree-level action, therefore, we omit the subscript $\epsilon$.

The massless scalar (quantum) field $\psi$ obeys the (covariant) wave equation, $\Box \psi = -(g)^{\mu\nu} \partial_{\mu} \psi \partial_{\nu} \psi = 0$. In the Fermi coordinates, $g_{\mu\nu}$ can be viewed as a perturbation to the Minkowski metric, inducing expansion series for the inverse and determinant of the metric: $g^{\mu\nu} = h^{\mu\nu} + \epsilon \partial^{\mu} \partial^{\nu} + O(\epsilon^2)$ and $\ln(\sqrt{-\det g}) = \epsilon \partial^{\mu} g^{\mu\nu} + O(\epsilon^2)$. $\epsilon$ is the systematic perturbation parameter introduced to keep track of the perturbation series, which means that all the components of the Riemann tensor in Eq. (1) are multiplied with $\epsilon$, and $\epsilon$ is treated as an infinitesimal parameter. At the end of the computation, we set $\epsilon = 1$. This technique helps us to systematically perform perturbations for small curvatures. Utilising the perturbation gives,

$$\Box \psi = \Box^{(0)} \psi + \epsilon (\partial_\mu (g^{\mu\nu} \partial_\nu \psi) + \partial_\mu \partial_\nu \partial_\nu \partial_\mu \psi) + O(\epsilon^2),$$

where $\Box^{(0)} \psi = 0$, and the perturbed term to the wave-function, $\psi^{(1)}$, yields,

$$\Box^{(0)} \psi^{(1)} = -\partial_\mu (g^{\mu\nu} \partial_\nu \psi^{(0)}) + \partial_\mu \partial_\nu \partial_\nu \partial_\mu \psi^{(0)}.$$

We assume the Fourier expansion in terms of $x^-$ variable, $\psi^{(0)} = \int d\omega \mathcal{F}^{(0)}(\omega, x^+) e^{i \omega x^-}$, where $\mathcal{F}^{(0)}$ satisfies the paraxial equation, $(2i\omega \partial_\omega + V^2) f^{(0)} = 0$. This implies that the solutions, given by the paraxial approximation in optics$^{23,24}$, are exact. The paraxial equation is isomorphic to the Schrödinger equation, and its solutions (the transverse and longitudinal parts) can be expressed in the form of Laguerre-Gauss (LG) modes (with an azimuthally symmetric intensity profile) or Hermite-Gaussian (HG) wavepackets$^{25}$. We consider a wavepacket wherein the field is slowly varying, and assume that the metric does not significantly change inside the wavepacket. Therefore, all derivatives of $\partial_\mu \psi^{(0)}$, except $\partial \psi^{(0)}$, are negligible, and the leading term in the right hand side of Eq. (3) is $\partial_\mu (g^{\mu\nu} \partial_\nu \psi)$, where $\phi g^{\mu\nu} = -\frac{1}{2} g_{\mu\nu} = R_{\mu\nu}^a x^a x^{\mu}$. Therefore, Eq. (3) reduces to,

$$\Box^{(0)} \psi^{(1)} = -R_{\mu\nu} \partial_\mu x^\nu \partial_\mu \psi^{(0)}.$$ (4)

Here, $\psi^{(1)} = \int d\omega \mathcal{F}^{(1)}(\omega, x^+) e^{i \omega x^-}$ with $\mathcal{F}^{(1)}$ being the correction to the structure function for the frequency of $\omega$. We have found the solutions to Eq. (4)—see the Supplementary Note 1 for more detail on the derivation. The solution is,

$$\psi(x^a) = \int d\omega \sum_{p, m, n} c_{p, m, n}(\omega) e^{i \omega x^-} (1 + \epsilon (O^{(0)} + O^{(a)} + O^{(b)})) f^{(0)}(\omega).$$

where $c_{p, m, n}(\omega)$ are defined based on the initial and boundary conditions. The operators, $O^{(a)}$, $O^{(b)}$, and $O^{(0)}$ encodes how the curvature of the spacetime geometry distorts the wavepacket. They are given by,

$$O^{(a)} = -\frac{i \omega}{2} g_{ab} x^b + \frac{1}{\omega} \partial_\omega g_{ab} \partial_\omega \partial_\omega \psi,$$

$$O^{(b)} = \int_{\Lambda} d\omega U^{(b)} g_{ab} x^b + \frac{1}{\omega} \partial_\omega U^{(b)} \partial_\omega \partial_\omega \psi,$$

$$O^{(0)} = x^a g_{ab} + g_{ab} x^b + 2 \partial_\omega g_{ab} \partial_\omega \partial_\omega \psi,$$

where $g_{ab}$, $\tilde{g}_{ab}$ are integrals of the components of the Riemann tensor $R_{\mu\nu}^a$ evaluated on the geodesic—see Supplementary Note 1:

$$G^{ab} = \int_0^{\tau} d\tau R_{+ab}^+, \quad \tilde{G}^{ab} = \int_0^{\tau} d\tau \tilde{R}_{ab}^+, \quad \tilde{G}^{ab} = \int_0^{\tau} d\tau \tilde{\tilde{G}}^{ab},$$ (6)

where $\tau$ is the affine parameter on the geodesic. The distortions provided by Eq. (5) is the solution to the exact quantum effective action and cannot be reproduced by motion of a point-like particle along a geodesic. They do not exist in flat spacetime geometry, so they manifest a set of quantum effects in curved spacetime geometry. A similar approach can be used to find the wavefunction of a massive scalar particle. The physical degrees of (U(1)) gauge fields get corrected as if they were scalar fields, see Supplementary Note 2 and Supplementary Note 3. Supplementary Note 4 presents the operators in the Hilbert space that corresponds to these corrections. The following subsections show how these operators distort the physical information encoded in wavepackets traveling along couple of examples of null geodesics in the Solar system and around the Earth.

We now study the distortion operators in a couple of spacetime geometries, including de Sitter and Schwarzschild spacetime geometries. We first study the de Sitter and anti de Sitter space-times because their symmetry allows one to immediately write down the components of the Riemann tensor in Fermi coordinates evaluated on the geodesic.

**de Sitter and anti de Sitter spacetime geometries.** The de Sitter and anti de Sitter space-times are maximally symmetric and the Riemann tensor at any given event in the spacetime in any coordinates, including the Fermi coordinates, is given by $R_{\mu\nu;\rho\sigma} = \Lambda g_{\mu\rho} g_{\nu\sigma} - g_{\mu\sigma} g_{\nu\rho}$. The value of $\Lambda$ determines different geometries: $\Lambda > 0$ represents the de Sitter space geometry; $\Lambda < 0$ represents the anti de Sitter space geometry; and $\Lambda = 0$ is the Minkowski spacetime geometry. $R_{\mu\nu} = \Lambda g_{\mu\nu}$ is the only non-zero component for $R_{\mu\nu}^a$ evaluated on the geodesic, and thus the correction operators, Eq. (5) are,

$$O^{(a)} = 0, \quad \mathcal{Q}^{(b)} = \frac{1}{\omega} \left(1 - \frac{2 \omega}{a^2 (x^-)^2}\right) \Lambda x^+, \quad \mathcal{Q}^{(0)} = \Lambda x^+ x^-.$$

(7)

Let us consider a Gaussian wavepacket with normal distribution for $\omega$ around $\omega_0$ with the width of $\sigma$, i.e. $\psi_{Alice} = f^{(0)}(x^+, x^-) e^{i \omega x^-} e^{-\omega^2/2 \sigma^2}$. The validity of the perturbative solution demands that $|\Lambda| \ll \omega_0^2$, $|\Lambda| \ll \sigma^2$ and $\sigma \ll \omega_0$. The wavepacket after the propagation is $\psi_{Bob} \simeq (1 + \epsilon \mathcal{Q}^{(a)} + \epsilon \mathcal{Q}^{(0)}) f^{(0)}(x^+) e^{-\omega^2/2 \sigma^2}$. $\mathcal{Q}^{(0)}$ and $\mathcal{Q}^{(b)}$ change the wavepacket amplitude and phase, respectively. The maximum of $|\mathcal{Q}^{(0)} \psi_{Alice}|$ occurs at $x^- = \pm 1/\sigma$. Requiring it to be smaller than 1 yields $T < \tau_a$ where $\tau_a = (\sqrt{\sigma}/|\Lambda|)$. $\tau_a$ is the maximum time that the wavepacket feels the curvature of the spacetime geometry and keeps its amplitude intact. $\mathcal{Q}^{(a)}$ alters the phase of the wavepacket. The maximum of $|\mathcal{Q}^{(a)} \psi_{Alice}|$ occurs at $x^- = \pm \sqrt{2}/\sigma$. Requiring it to be smaller than 1 results in $T < \tau_p$ where $\tau_p = (|\sigma|/\sqrt{2}|\omega_0|)$.  


\( \tau_p \) is the maximum time that a wavepacket can feel the curvature of the spacetime geometry and keep its phase intact. For \( T > \tau_p \), information about phase is lost at perturbation. This may point to a "gravitational decoherence", and its possible consequence on anti de Sitter/conformal field theory correspondence demands attention. \( \tau_A \) represents the amount of time interaction with the curvature that the wavepacket can keep its amplitude intact. We observe that \( \tau_p < \tau_A \). So the phase changes sooner than the change in the amplitude.

**Schwarzschild spacetime geometry.** We choose the standard spherical coordinates \( r, \theta, \varphi \) where geodesics are extrema of

\[
 \mathcal{L} = -\left(1 - \frac{m}{r}\right)\dot{r}^2 + \left(1 - \frac{m}{r}\right)^{-1} \dot{\theta}^2 + \sin^2 \theta \sin^2 \varphi \left(1 - \frac{1}{\tau_\rho} \right)^2 + \right(1 - \frac{1}{\tau_\rho} \right)^2 \phi + \sin^2 \theta \dot{\varphi}^2 ,
\]

and \( m = 2GM \), is the Schwarzschild radius, \( M \) is the mass of the blackhole and \( G \) is the gravitational constant—the units are such that \( m = 1 \). The components of the Riemann tensor in the Fermi coordinates adapted to a general null geodesic of Schwarzschild spacetime geometry are derived in the Supplementary Note 5. Figure 2a shows several null geodesics that go very close to a blackhole. We first consider that the wavepacket propagates along the radial direction, Fig. 2b, where the only non-zero components of the Riemann tensor is \( R_{++} = -1/r^2 \). This is the same component that appeared in the de Sitter spacetime geometry. The radial geodesic has \( l = 0 \), and its correction operators are \( \mathcal{O}_{00}^\gamma = 0 \), \( \mathcal{Q}_{00}^\gamma = -\frac{i}{2} \left(1 - \frac{\omega x - \omega y}{2} \right) \left( \frac{1}{\tau_\rho} - \frac{1}{\tau_\rho} \right) \), and \( \mathcal{Q}_{00}^\gamma = -\frac{i}{2} \left( \frac{1}{\tau_\rho} - \frac{1}{\tau_\rho} \right) x^- \), where Alice is located at \( r_\alpha \). These correction terms do not contain derivatives of the spatial transverse coordinates. Thus, the Riemann tensor does not affect the spatial transverse profile of the wavepacket. This is due to the symmetry, as the radial geodesic inherits the static and spherical symmetry of the background. Figure 3a shows the amplitude and phase of a Gaussian time-bin signal. Figure 3b depicts the alternations in the amplitude and the phase:

\[
 \delta \mathcal{A} = -\frac{\varepsilon x^-}{2} \left( \frac{1}{r_\alpha} - \frac{1}{r_\beta} \right) e^{\frac{ix^-}{2}} + O(x^-^2) \]

\[
 \delta \chi = \frac{i\omega_0}{4\sigma^2} \left( \frac{1}{r_\alpha} - \frac{1}{r_\beta} \right) (\sigma x^-)^2 e^{-\frac{x^-}{2}} + O(x^-^2) ,
\]

where \( r_\beta \) is the location of Bob. The maximum alteration to the amplitude and phase of the Gaussian wavepacket occur at \( x^- = \pm 1/\sigma \) and \( x^- = \pm \sqrt{2}/\sigma \), respectively. For a Gaussian wavepacket that propagates radially close to the Earth, the maximum alteration to the amplitude and phase are, respectively, \( |\delta \mathcal{A}_{\text{max}}| = \frac{1}{2\sigma} \left( \frac{1}{\tau_\rho} - \frac{1}{\tau_\rho} \right) \) and \( \delta \chi_{\text{max}} = \frac{\omega_0}{4\sigma^2} \left( \frac{1}{\tau_\rho} - \frac{1}{\tau_\rho} \right) \), where \( m_{\text{Earth}} \) is the Schwarzschild radius of Earth.

For \( \nu_0 = 456 \) THz and \( \Delta \nu = 1 \) kHz, \( \delta \chi_{\text{max}} = 0.10 \) rad. Supplementary Note 8 provides further details on choosing these values.

As a final example, we examine the weak regime of gravity when the beam possesses well-defined transverse modes—see Fig. 2a. The wavepacket carrying a well-defined transverse mode traverses through space and reaches to the minimum distance of \( l \) to the central mass—here, we assume \( l \) is large. We now consider a specific wavepacket, a Hermite-Gauss transverse mode \( f_{0,0,1,0}(x^- , x^+ , x^1 , x^2) = e^{i\omega x^-} x^2 e^{-\frac{x^-}{2}} \)—Hermite-Gauss modes are used to extend the communication alphabet beyond bits, i.e. 0 and 1. The longitudinal and frequency distributions are assumed to be Gaussian. When \( \sigma \) is large, the dominant correction operator is calculated to be:

\[
 O = -\frac{9.5i\omega^2}{z_R^2} (x^2 - y^2) ,
\]

where \( x = \sqrt{2} x_1 / w(x^+ ) \) and \( y = \sqrt{2} x_2 / w(x^+ ) \) are dimensionless coordinates, \( w(x^+) = w_0 \sqrt{1 + \left( \frac{x^{-}}{z_R} \right)^2} \) is the beam radius, \( z_R = \)

![Fig. 2 Propagation of wavepacket and associated null geodesics in the Schwarzschild spacetime geometry.](image-url)
The dimensionless coordinates are used for these plots. For instance, it causes up to 12.2% crosstalk between HG0,3 and HG0,1 modes. The crosstalk between mode HG0,3 and HG0,1 is mostly due to the beam radius at Alice’s position—see Supplementary Note 6 for more details. This operator contains coordinate parameters \( x \) and \( y \), and thus, alters both the amplitude and phase of the transverse modes upon propagation. The correction for the solar system, when Alice and Bob are at the mean Earth-Sun distance from the Sun and the wavepacket passes at \( l = 2R_o \), and for \( z_R = 2.8 \text{ km} \times a^2 = 1.34 \times 10^9 \text{ m} \), remains perturbative and is given by \( \epsilon \psi^{(1)}(x^\mu) = i 0.10 \ (x^2 - y^2) \psi_{\text{Alice}}(x^\mu) |_{x^\mu = x^\mu} \).

The amplitude and phase of the mode \( \psi^{(0)}(x^\mu) \) and the correction \( \epsilon \psi^{(1)}(x^\mu) \) for a few Hermite-Gaussian modes are shown in Fig. 4. As seen, these alterations on the modes are considerable. For instance, it causes up to 12.2% crosstalk between HG0,3 and HG0,1 modes.

The incidence of cross-talk is significant and it alters both the amplitude and phase of the modes. Moreover, whenever \( \epsilon \approx 1 \), the higher-order terms of correction need to be considered. For instance, for \( z_R < 2.8 \text{ km} \times a^2 \), the correction becomes larger than 1 and we need to take into account higher \( \epsilon \) terms. Taking into account all the corrections is tantamount to knowing the Riemann tensor in whole of the spacetime geometry, a piece of knowledge which is not attainable.

Thus, we tend to argue that once the perturbation breaks, in addition to known well-studied gravitational decoherence, a decoherence occurs. We observe that, in addition to the known decoherence of a bipartite entangled system when each particle traverses through a different gravitational field gradient, a coherent beam decoheres when different segments of the spatial spread of the wave experience different tidal gravitational field gradients. The phenomenon we are reporting also occurs for geodesics passing very close to the event horizon—see Supplementary Note 7.

Finally, it is noteworthy that photon pairs \( |\Psi_{\text{entangled}}\rangle \) in entangled spatial, frequency or temporal modes, would be affected by the curved spacetime geometry whenever they are shared between two parties, namely Alice and Bob. The final state of the entangled photon, indeed, is given by applying the non-local operators \( \mathcal{U} = (1 + \epsilon O + U_{\text{out}} + Q_{\text{out}}) \) onto the entangled states, \( \langle U_A \otimes U_B \rangle |\Psi_{\text{entangled}}\rangle \) here, \( U_A \) and \( U_B \) are associated with the correction operators at Alice and Bob’s places, respectively.

**Conclusion**

We have presented how the curvature of the spacetime geometry affects the propagation of an arbitrary wavepacket along a general geodesic in a general curved spacetime geometry. The effect is beyond classical general relativity, residing in the same category as Hawking radiation. A set of linear operators are presented that encode the effect of the curvature. The corrections to the information carrier wavepacket are investigated in cases of de Sitter and Schwarzschild spacetime geometries. It has been shown that the corrections accumulate overtime and distort...
the wavepacket. The gravitational distortion, therefore, needs to be accounted for in quantum communication performed over long distances in a curved spacetime geometry.

Data availability
The authors declare that the data supporting the findings of this study are available within the paper and its supplementary information file.

Received: 22 February 2021; Accepted: 5 July 2021; Published online: 30 July 2021

References
1. Ursin, R. et al. Entanglement-based quantum communication over 144 km. New J. Phys. 3, 1–14 (2001).
2. Yin, J. et al. Satellite-based entanglement distribution over 1200 kilometers. Science 356, 1140–1144 (2017).
3. Yin, J. et al. Satellite-to-ground entanglement-based quantum key distribution. Phys. Rev. Lett. 119, 200501 (2017).
4. Wenerowsky, S. et al. Passively stable distribution of polarisation entanglement over 192 km of deployed optical fibre. Quantum Inf. & Comp. 6, 5 (2020).
5. Hufnagel, F. et al. Investigation of underwater quantum channels in a 30 meter flume tank using structured photons. N. J. Phys. 22, 093074 (2020).
6. Einstein, A. Über den einfluß der schwerkraft auf die ausbreitung des lichtes. Ann. der Phys. 340, 898–908 (1911).
7. Pound, R. V. & Rebka, G. A. Apparent weight of photons. Phys. Rev. Lett. 4, 337–341 (1960).
8. Chou, C.-W., Hume, D. B., Rosenband, T. & Wineland, D. J. Optical clocks and relativity. Science 329, 1630–1633 (2010).
9. Ashby, N. Relativity in the global positioning system. Living Rev. Relativ. 6, 1 (2003).
10. Shapiro, I. et al. Fourth test of general relativity: preliminary results. Phys. Rev. Lett. 20, 1265–1269 (1968).
11. Perlick, V. Gravitational lensing from a spacetime perspective. Living Rev. Relativ. 7, 9 (2004).
12. Everitt, C. et al. Gravity probe B: final results of a space experiment to test general relativity. Phys. Rev. Lett. 106, 221101 (2011).
13. Will, C. G. The confrontation between general relativity and experiment. Living Rev. Relativ. 17, 4 (2011).
14. Liu, T., Wu, S. & Cao, S. The influence of the Earth’s curved spacetime on Gaussian quantum coherence. Laser Phys. Lett. 16, 095201 (2019).
15. Kish, S. P. & Ralph, T. C. Quantum metrology in the Kerr metric. Phys. Rev. D. 99, 124015 (2019).
16. Pierini, R. Effects of gravity on continuous-variable quantum key distribution. Phys. Rev. D. 98, 125007 (2018).
17. Bruschi, D. E., Ralph, T., Fuentes, I., Jennewein, T. & Rzazewski, M. Spacetime effects on satellite-based quantum communications. Phys. Rev. D. 90, 045041 (2014).
18. Bruschi, D. E., Datta, A., Ursin, R., Ralph, T. C. & Fuentes, I. Quantum estimation of the Schwarzschild spacetime parameters of the Earth. Phys. Rev. D. 90, 124001 (2014).
19. Jonsson, R. H., Aruquipa, D. Q., Casals, M., Kempf, A. & Martin-Martinez, E. Communication through quantum fields near a black hole. Phys. Rev. D. 101, 125005 (2020).
20. Ferrini, E. Sopra i fenomeni che avvengono in vinicinanza di una linea oraria. Atti R. Accad. Lincei Rend. Cl. Sci. Fis. Mat. Nat. 31, 21–51 (1922).
21. Blu, M., Frank, D. & Weiss, S. Fermi coordinates and Penrose limits. Classical Quantum Gravity 23, 3993–4010 (2006).
22. Dirac, P. A. M. Forms of relativistic dynamics. Rev. Mod. Phys. 21, 392–399 (1949).
23. Segnan, A. E. Lasers university science books. Mill. Val., CA 37, 462–466 (1986).
24. Białynicki-Birula, I. & Białynicka-Birula, Z. Beams of electromagnetic radiation carrying angular momentum: the riemann–silberstein vector and the classical–quantum correspondence. Opt. Commun. 264, 342–351 (2006).
25. Bloch, K. Y., Bloch, Y. P., Savélev, S. & Nori, F. Semicliscal dynamics of electron wavepacket states with phase vortices. Phys. Rev. Lett. 99, 190404 (2007).
26. Wang, Y. M. The Large N limit of superconformal field theories and supergravity. Int. J. Theor. Phys. 38, 1113–1133 (1999).
27. Krenn, M. et al. Communication with spatially modulated light through turbulent air across Vienna. N. J. Phys. 16, 113028 (2014).
28. Pang, B. H., Chen, Y. & Khalili, F. Y. Universal decoherence under gravity: a perspective through the equivalence principle. Phys. Rev. Lett. 117, 090401 (2016).
29. Stefanov, V., Suitsou, I. & Mogilevtsev, D. Gravitational dephasing in spontaneous emission of atomic ensembles in timed Dicke states. Phys. Rev. D. 101, 044042 (2020).
30. Bassi, A., Großardt, A. & Ulbricht, H. Gravitational decoherence. Classical Quantum Gravity 34, 193002 (2017).
31. Penrose, R. On gravity’s role in quantum state reduction. Gen. Relativ. Gravit. 28, 581–600 (1996).
32. Joshi, S. K. et al. [Space QUEST topical Team], Space QUEST mission proposal: experimentally testing decoherence due to gravity. N. J. Phys. 20, 063016 (2018).
33. Ralph, T. C., Milburn, G. J. & Downes, T. Quantum connectivity of space-time and gravitationally induced decorrelation of entanglement. Phys. Rev. A 79, 022121 (2009).
34. Hawking, S. W. Black hole explosions? Nature 248, 30–31 (1974).