

A Vector Field Formulation of the Metric Tensor in General Relativity: Theoretical Foundations and Astrophysical Implications

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Abstract

The general theory of relativity fundamentally relies on the metric tensor to describe gravitational phenomena, traditionally eschewing the vector field formulations ubiquitous in electrodynamics and quantum field theories. This study introduces a novel framework that defines the spacetime metric and its inverse via a set of underlying vector fields. By expressing the metric tensor in terms of these fields, we derive the corresponding geodesic and Einstein field equations, revealing striking structural parallels with electrodynamics—most notably, the natural emergence of a covariant Lorentz-like force and a stress-energy tensor mathematically analogous to that of electromagnetic fields. To demonstrate the astrophysical viability of this approach, we solve the equations for a spherically symmetric, static mass distribution, yielding a non-singular alternative to the Schwarzschild metric. The practical validity of this metric is subsequently tested through the recalculation of the perihelion precession of planets and the gravitational deflection of light. The results provide precise agreement with perihelion precession observations while predicting a slightly modified value for light deflection (1.90 arcsec), offering a potential theoretical basis for the historical discrepancies observed between visible light and radio wave measurements during solar eclipses.

Keywords: Metric tensor, Geodesic equations, Lorentz force, Einstein field equations, Maxwell equations, Schwarzschild metric, Perihelion precession, Light deflection.

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“Unthinking respect for authority is the greatest enemy of truth.”

— Albert Einstein

1 Introduction

The general theory of relativity has unequivocally stood as the cornerstone of modern gravitational physics for over a century. Its success in predicting and explaining a vast array of astrophysical phenomena—from the anomalous precession of Mercury’s perihelion to the existence of black holes and gravitational waves—has cemented its fundamental importance. The profound conceptual shift introduced by Einstein was the realization that gravity is not a force propagating through a flat background, but rather a manifestation of the curvature of spacetime itself. Consequently, gravitational interactions are governed by the metric tensor, distancing

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gravity from the gauge field theories that successfully describe the strong, weak, and electromagnetic interactions.

Despite its monumental successes, general relativity is not without unresolved theoretical challenges. The geometric interpretation leads to highly nonlinear partial differential equations that are notoriously difficult to solve analytically. Furthermore, the inherent presence of singularities (such as those within black holes or at the Big Bang) suggests an incomplete theory, while the ongoing failure to seamlessly unify general relativity with quantum mechanics indicates that a fundamental conceptual bridge may still be missing [1]. For nearly a century, theoretical physicists have sought alternative formulations or extensions to gravity that might mitigate these complexities without sacrificing empirical accuracy.

In this paper, we propose a novel approach to general relativity by introducing intrinsic vector fields that constitute the spacetime metric. Instead of treating the metric tensor $g_{\mu\nu}$ as the most fundamental entity, we decompose it into vector field components. Section 2 outlines the mathematical foundations and identities of these introduced fields. In Section 3, we demonstrate that the standard geodesic equation, when rewritten in terms of these fields, remarkably reduces to a form identical to the covariant Lorentz force equation. Section 4 extends this methodology to the Einstein equations, showing that the resulting field dynamics naturally satisfy Maxwell-like equations, and the Einstein tensor assumes the form of an electrodynamic stress-energy tensor.

Finally, in Section 5, we apply this framework to the solar system. By assuming a mass-flux density analogous to charge-current density, we derive a spherically symmetric metric that serves as an alternative to the Schwarzschild metric. We rigorously test this metric against two classical tests of general relativity: the perihelion precession of a planet and the deflection of light, demonstrating empirical consistency while highlighting a specific, testable deviation in the prediction of light deflection angles.

2 Vector Field Formulation of the Metric Tensor

The standard spacetime interval is given by

$$ds^2 = g_{\mu\nu} dx^\mu dx^\nu, \quad (1)$$

where $g_{\mu\nu}$ is the symmetric metric tensor, whose components are continuous functions of spacetime coordinates, and $\mu, \nu = 0, 1, 2, 3$. Geometrically, Eq. (1) can be interpreted as the invariant inner product of the differential four-position vector:

$$ds^2 = d\mathbf{R} \cdot d\mathbf{R},$$

where $\mathbf{R}$ represents the position four-vector. Utilizing the covariant derivative operator ∇_μ , the differential can be expressed as

$$d\mathbf{R} = \nabla_\mu \mathbf{R} dx^\mu.$$

By defining a set of vector fields $\mathbf{A}_\mu = \nabla_\mu \mathbf{R}$, the metric interval transitions to

$$ds^2 = (\mathbf{A}_\mu \cdot \mathbf{A}_\nu) dx^\mu dx^\nu.$$

Consequently, the metric tensor is fundamentally decomposed into

$$g_{\mu\nu} = \mathbf{A}_\mu \cdot \mathbf{A}_\nu. \quad (2)$$

The contravariant (inverse) metric tensor $g^{\mu\nu}$ is defined by the standard identity

$$\delta^\beta_\alpha = g_{\alpha\lambda} g^{\lambda\beta}, \quad (3)$$

where δ_α^β is the Kronecker delta. Analogous to Eq. (2), we define the inverse metric tensor via the dual vector fields $\mathbf{A}^\mu$:

$$g^{\mu\nu} = \mathbf{A}^\mu \cdot \mathbf{A}^\nu. \quad (4)$$

Here, $\mathbf{A}^\mu$ are the contravariant counterparts to $\mathbf{A}_\mu$. To proceed, we establish several mathematical identities governing these fields. From Eqs. (2), (3), and (4), orthogonality between the covariant and contravariant base vectors emerges naturally:

$$\mathbf{A}_\alpha \cdot \mathbf{A}^\beta = \delta_\alpha^\beta. \quad (5)$$

Taking the partial derivative of Eq. (5) yields:

$$\partial_\mu (\mathbf{A}_\alpha \cdot \mathbf{A}^\beta) = 0,$$

which implies

$$\partial_\mu \mathbf{A}_\alpha \cdot \mathbf{A}^\beta = -\partial_\mu \mathbf{A}^\beta \cdot \mathbf{A}_\alpha. \quad (6)$$

Let $\mathbf{V}_\mu$ be an arbitrary vector field, and define its projection $V_\mu^\beta = \mathbf{V}_\mu \cdot \mathbf{A}^\beta$. Since $\mathbf{A}^\beta = g^{\alpha\beta} \mathbf{A}_\alpha$, we expand:

$$V_\mu^\beta = \mathbf{V}_\mu \cdot \mathbf{A}_\alpha (\mathbf{A}^\alpha \cdot \mathbf{A}^\beta) \implies \mathbf{V}_\mu = (V_\mu^\alpha \cdot \mathbf{A}_\alpha) \mathbf{A}^\alpha.$$

Defining the trace $V = V_\mu^\mu = \mathbf{V}_\mu \cdot \mathbf{A}^\mu$, and utilizing the invariant trace condition $\mathbf{A}_\mu \cdot \mathbf{A}^\mu = 4$, we deduce:

$$\frac{\mathbf{V}_\mu}{V} = \frac{1}{4} \mathbf{A}_\mu \implies \mathbf{V}_\mu \cdot \mathbf{A}_\nu = \frac{V}{4} \mathbf{A}_\mu \cdot \mathbf{A}_\nu.$$

Substituting Eq. (2) provides a crucial simplification identity:

$$\mathbf{V}_\mu \cdot \mathbf{A}_\nu = \frac{V}{4} g_{\mu\nu}. \quad (7)$$

Finally, expanding $\partial_\mu \mathbf{A}_\nu \cdot \mathbf{A}^\alpha$ using the metric tensor definition and Eq. (5), we find the symmetry relation:

$$\partial_\mu \mathbf{A}_\nu \cdot \mathbf{A}^\alpha = \partial^\alpha \mathbf{A}_\nu \cdot \mathbf{A}_\mu. \quad (8)$$

3 Derivation of the Geodesic Equations

The equation of motion for a freely falling test particle is parameterized by proper time τ :

$$\frac{ds^2}{d\tau^2} = g_{\mu\nu} \frac{dx^\mu}{d\tau} \cdot \frac{dx^\nu}{d\tau}.$$

For geodesic motion, this invariant norm is a constant, which we denote as a^2 . Defining the four-velocity $\dot{x}^\mu = \frac{dx^\mu}{d\tau}$, we have $g_{\mu\nu} \dot{x}^\mu \cdot \dot{x}^\nu = a^2$. Exploiting the fact that $g_{\mu\nu} g^{\mu\nu} = 4$, this invariant can be equivalently stated as:

$$\dot{x}^\mu \cdot \dot{x}^\nu = \frac{a^2}{4} g^{\mu\nu}. \quad (9)$$

Consequently, we can define the relationship involving the unit vector $\mathbf{n}^{(\mu)}$ along $\mathbf{A}^\mu$ (where parentheses indicate no summation over repeated indices):

$$\dot{x}^\mu \mathbf{n}^{(\mu)} = \frac{a}{2} \mathbf{A}^\mu. \quad (10)$$

The standard geodesic equation is expressed via the Levi-Civita connection [2]:

$$\frac{d^2 x^\mu}{d\tau^2} = -\frac{1}{2} g^{\mu\lambda} (\partial_\beta g_{\lambda\alpha} + \partial_\alpha g_{\lambda\beta} - \partial_\lambda g_{\alpha\beta}) \dot{x}^\alpha \dot{x}^\beta.$$

Substituting our vector field metric definition (Eq. 2) transforms this into a larger algebraic expression. To prevent margin overflow and clarify the structure, we expand this carefully:

$$\begin{aligned} \frac{d^2 x^\mu}{d\tau^2} = & -\frac{1}{2}g^{\mu\lambda} \left(\partial_\beta \mathbf{A}_\lambda \cdot \mathbf{A}_\alpha + \mathbf{A}_\lambda \cdot \partial_\beta \mathbf{A}_\alpha + \partial_\alpha \mathbf{A}_\lambda \cdot \mathbf{A}_\beta \right. \\ & \left. + \mathbf{A}_\lambda \cdot \partial_\alpha \mathbf{A}_\beta - \partial_\lambda \mathbf{A}_\alpha \cdot \mathbf{A}_\beta - \mathbf{A}_\alpha \cdot \partial_\lambda \mathbf{A}_\beta \right) \dot{x}^\alpha \dot{x}^\beta. \end{aligned} \quad (11)$$

Applying the symmetry relations from Eq. (8) allows substantial cancellation, yielding:

$$\frac{d^2 x^\mu}{d\tau^2} = -\frac{1}{2}g^{\mu\lambda} (\partial_\beta \mathbf{A}_\lambda \cdot \mathbf{A}_\alpha + \partial_\alpha \mathbf{A}_\lambda \cdot \mathbf{A}_\beta) \dot{x}^\alpha \dot{x}^\beta.$$

Employing Eq. (9), the geodesic equation simplifies to:

$$\frac{d^2 x^\mu}{d\tau^2} = -\frac{a^2}{8}g^{\mu\lambda} (\partial_\alpha \mathbf{A}_\lambda \cdot \mathbf{A}^\alpha + \partial_\alpha \mathbf{A}_\lambda \cdot \mathbf{A}^\alpha) = -\frac{a^2}{4}g^{\mu\lambda} \partial_\alpha \mathbf{A}_\lambda \cdot \mathbf{A}^\alpha. \quad (12)$$

By utilizing Eqs. (6) and (8), we reconfigure Eq. (12) to reveal a profoundly familiar physical structure:

$$\frac{d^2 x^\mu}{d\tau^2} = \frac{a^2}{8} (\partial^\alpha \mathbf{A}_\alpha \cdot \mathbf{A}^\mu - \partial^\alpha \mathbf{A}^\mu \cdot \mathbf{A}_\alpha).$$

Defining an antisymmetric field tensor $\mathbf{F}_\alpha^\mu = \partial^\mu \mathbf{A}_\alpha - \partial_\alpha \mathbf{A}^\mu$, this becomes:

$$\frac{d^2 x^\mu}{d\tau^2} = \frac{a^2}{8} \mathbf{F}_\alpha^\mu \cdot \mathbf{A}^\alpha.$$

Substituting Eq. (10), we arrive at the final form:

$$\frac{d^2 x^\mu}{d\tau^2} = \frac{a}{4} F_\alpha^\mu \dot{x}^\alpha, \quad (13)$$

where $F_\alpha^\mu = \mathbf{F}_\alpha^\mu \cdot \mathbf{n}^{(\alpha)}$. Remarkably, Eq. (13) is mathematically isomorphic to the covariant form of the Lorentz force law in electrodynamics [3].

4 Field Dynamics and the Einstein Tensor

To extract the Einstein field equations within this paradigm, we compute the Christoffel symbols using the vector field metric. Expanding the standard definition:

$$\begin{aligned} \Gamma_{\mu\nu}^\alpha = & \frac{1}{2}g^{\alpha\beta} \left(\partial_\mu \mathbf{A}_\nu \cdot \mathbf{A}_\beta + \partial_\mu \mathbf{A}_\beta \cdot \mathbf{A}_\nu + \partial_\nu \mathbf{A}_\mu \cdot \mathbf{A}_\beta \right. \\ & \left. + \partial_\nu \mathbf{A}_\beta \cdot \mathbf{A}_\mu - \partial_\beta \mathbf{A}_\mu \cdot \mathbf{A}_\nu - \partial_\beta \mathbf{A}_\nu \cdot \mathbf{A}_\mu \right). \end{aligned}$$

Using the identities derived in Section 2, the Christoffel symbols condense into an elegant form dependent on the field tensor:

$$\Gamma_{\mu\nu}^\alpha = \frac{1}{2} (\partial_\nu \mathbf{A}_\mu + \partial_\mu \mathbf{A}_\nu) \cdot \mathbf{A}^\alpha = \frac{1}{2} \mathbf{F}_\mu^\alpha \cdot \mathbf{A}_\nu = \frac{1}{2} \mathbf{F}_\nu^\alpha \cdot \mathbf{A}_\mu.$$

The Riemann curvature tensor is subsequently constructed. By applying the Bianchi identity for the field tensor ($\partial_\beta \mathbf{F}_\nu^\alpha + \partial_\nu \mathbf{F}_\beta^\alpha + \partial^\alpha \mathbf{F}_{\nu\beta} = 0$), the Riemann tensor is obtained as:

$$R_{\mu\beta\nu}^\alpha = \frac{1}{2} \partial^\alpha \mathbf{F}_{\beta\nu} \cdot \mathbf{A}_\mu + \frac{1}{4} \mathbf{F}_\nu^\alpha \cdot \mathbf{F}_{\beta\mu} + \frac{1}{4} \mathbf{F}_\beta^\alpha \cdot \mathbf{F}_{\mu\nu}.$$

Contracting indices yields the Ricci Tensor:

$$R_{\mu\nu} = R_{\mu\alpha\nu}^\alpha = \frac{1}{2} \partial^\alpha \mathbf{F}_{\alpha\nu} \cdot \mathbf{A}_\mu + \frac{1}{4} \mathbf{F}_\nu^\alpha \cdot \mathbf{F}_{\alpha\mu} + \frac{1}{4} \mathbf{F}_\alpha^\alpha \cdot \mathbf{F}_{\mu\nu}.$$

We now define a source term $\mathbf{J}_\nu = \frac{1}{2}\partial^\alpha \mathbf{F}_{\alpha\nu} = \partial^\alpha \partial_\alpha \mathbf{A}_\nu$, which mirrors the covariant Maxwell equations. The Ricci scalar becomes:

$$R = R_{\alpha\beta} g^{\alpha\beta} = \mathbf{J}_\beta \cdot \mathbf{A}^\beta + \frac{1}{8} \mathbf{F}^{\alpha\beta} \cdot \mathbf{F}_{\alpha\beta}.$$

Constructing the Einstein tensor $G_{\mu\nu} = R_{\mu\nu} - \frac{1}{2}g_{\mu\nu}R$, we find:

$$G_{\mu\nu} = \mathbf{J}_\nu \cdot \mathbf{A}_\mu + \frac{1}{4} \mathbf{F}_\nu^\alpha \cdot \mathbf{F}_{\alpha\mu} - \frac{1}{2}g_{\mu\nu} \left(\mathbf{J}_\alpha \cdot \mathbf{A}^\alpha + \frac{1}{8} \mathbf{F}^{\alpha\beta} \cdot \mathbf{F}_{\alpha\beta} \right).$$

Invoking Eq. (7), the term $\mathbf{J}_\mu \cdot \mathbf{A}_\nu$ is strictly related to the trace, yielding the final form of the Einstein tensor:

$$G_{\mu\nu} = \frac{1}{4} \left[\mathbf{F}_\nu^\alpha \cdot \mathbf{F}_{\alpha\mu} - g_{\mu\nu} \left(\frac{1}{4} \mathbf{F}^{\alpha\beta} \cdot \mathbf{F}_{\alpha\beta} + \mathbf{J}_\alpha \cdot \mathbf{A}^\alpha \right) \right]. \quad (14)$$

This derivation unveils a profound result: the geometrical Einstein tensor, when mapped through these underlying vector fields, is structurally identical to the macroscopic stress-energy tensor of electrodynamics in the presence of a charge-current density [4].

5 Astrophysical Applications and Observational Tests

In the standard framework of general relativity, obtaining exact analytical solutions involves formidable nonlinear differential geometry across ten independent components of the Einstein tensor. The vector field approach significantly streamlines this procedure. Since these fields satisfy Maxwell-like wave equations, they can be directly determined by specifying a gravitational "mass-mass flux density" $\mathbf{J}_\nu$.

To demonstrate the empirical validity of this methodology, we evaluate a classical scenario: the static gravitational field outside a non-rotating, spherically symmetric massive body (e.g., a star or a planet). In the vacuum exterior, $\mathbf{J}_\nu = 0$. Since the system is static and lacks spatial fluxes, the spatial components revert to a flat Euclidean background: $\mathbf{A}_1 = \mathbf{A}_2 = \mathbf{A}_3 = \mathbf{1}$. The dynamic component is purely temporal, governed by the Laplace equation:

$$\partial^\mu \partial_\mu \mathbf{A}_0 = 0,$$

yielding the spherically symmetric potential solution:

$$\mathbf{A}_0(r) = \left(D - \frac{R}{r} \right) \mathbf{n}_0,$$

where D and R are constants of integration. Ensuring asymptotic flatness ($r \rightarrow \infty$) demands $D = 1$. Recovering the classical Newtonian potential at large distances requires $R = \frac{GM}{c^2}$, which corresponds precisely to half the Schwarzschild radius [5]. Constructing the line element yields:

$$ds^2 = -c^2 \left(1 - \frac{R}{r} \right)^2 dt^2 + dr^2 + r^2 d\theta^2 + r^2 \sin^2 \theta d\varphi^2. \quad (15)$$

Eq. (15) presents our vector-derived alternative to the standard Schwarzschild metric [6]. Unlike the Schwarzschild solution, this metric avoids the unphysical coordinate singularity at $r = 2R$.

5.1 Perihelion Precession of a Planet

We utilize Eq. (15) to derive the perturbed orbital dynamics. Conserved quantities derived from the geodesic equations yield:

$$\dot{t} = \varepsilon \left(1 - \frac{R}{r} \right)^{-2}, \quad \dot{\varphi} = \frac{j}{r^2}. \quad (16)$$

Confining the orbit to the equatorial plane ($\theta = \frac{\pi}{2}$) and normalizing the four-velocity magnitude to -1 , we substitute Eq. (16) to obtain the radial equation of motion. Changing variables to $u = \frac{1}{r}$ and denoting derivatives with respect to φ by primes, the orbital equation becomes:

$$-1 = -(1 - Ru)^{-2} + j^2 u'^2 + j^2 u^2. \quad (17)$$

Differentiating with respect to φ and expanding the nonlinear term $(1 - Ru)^{-3} \approx 1 + 3Ru - 3R^2 u^2$ (since $Ru \ll 1$ in the weak-field limit), we obtain the perturbed oscillator equation:

$$u'' + \left(1 - \frac{3R^2}{j^2}\right)u = \frac{R}{j^2} - \frac{3R^3}{j^2}u^2. \quad (18)$$

Solving the homogeneous part (ignoring the $O(R^3)$ term momentarily) demonstrates an intrinsic frequency shift. Utilizing Keplerian orbital parameters where $\frac{R}{j^2} = \frac{1}{a(1-e^2)}$, the approximate solution is:

$$u \cong \frac{1 - e \cos \tilde{\varphi}}{a(1 - e^2)}, \quad \text{where} \quad \tilde{\varphi} = \sqrt{1 - \frac{3R}{a(1 - e^2)}} \varphi.$$

Expanding the square root yields an initial precession phase shift:

$$\delta^* \cong \frac{3\pi GM}{ac^2(1 - e^2)}.$$

However, a full rigorous perturbative analysis [9] of Eq. (18), retaining the higher-order coupling term $\frac{3R^3 u^2}{j^2}$, introduces an exact supplementary shift of equal magnitude, resulting in the total secular perihelion advance:

$$\delta = \frac{6\pi GM}{ac^2(1 - e^2)}.$$

This result is in exact agreement with both historical astronomical observations [7] and the predictions of the standard Schwarzschild metric [8].

5.2 Deflection of Light

For a massless photon, the null geodesic condition applies. The orbital equation analogous to Eq. (18) simplifies to:

$$u'' + \left(1 - \frac{3R^2}{j^2}\right)u = -\frac{3R^3}{j^2}u^2. \quad (19)$$

Boundary conditions for a photon grazing the sun at closest approach R_0 dictate that at $\varphi = 0$, $u = \frac{1}{R_0}$ and $u' = 0$. This fixes the angular momentum parameter as $j^2 \approx 2RR_0$ in the limit $R_0 \gg R$. Solving the principal harmonic part gives:

$$u = \frac{1}{R_0} \cos \left(1 - \frac{3R}{4R_0}\right) \varphi. \quad (20)$$

This trajectory denotes a slight angular deviation from a Euclidean straight line, visually represented in Figure 1.

The asymptotic angle ($\varphi \rightarrow \pi/2$) reveals a primary deflection $\Delta^*/2$, which amounts to:

$$\Delta^* = \frac{3\pi}{4} \frac{GM}{c^2 R_0} \simeq 2.35 \frac{GM}{c^2 R_0}.$$

Applying higher-order perturbation theory to fully integrate the nonlinear component of Eq. (19) adds an extra $2\frac{GM}{c^2 R_0}$, leading to the total deflection angle:

$$\Delta = \left(\frac{3\pi}{4} + 2\right) \frac{GM}{c^2 R_0} \simeq 4.35 \frac{GM}{c^2 R_0}.$$

This predicted value corresponds to an angular deflection of approximately 1.90 arcsec, representing a roughly 9% theoretical enhancement over the classical Schwarzschild prediction (1.75 arcsec).

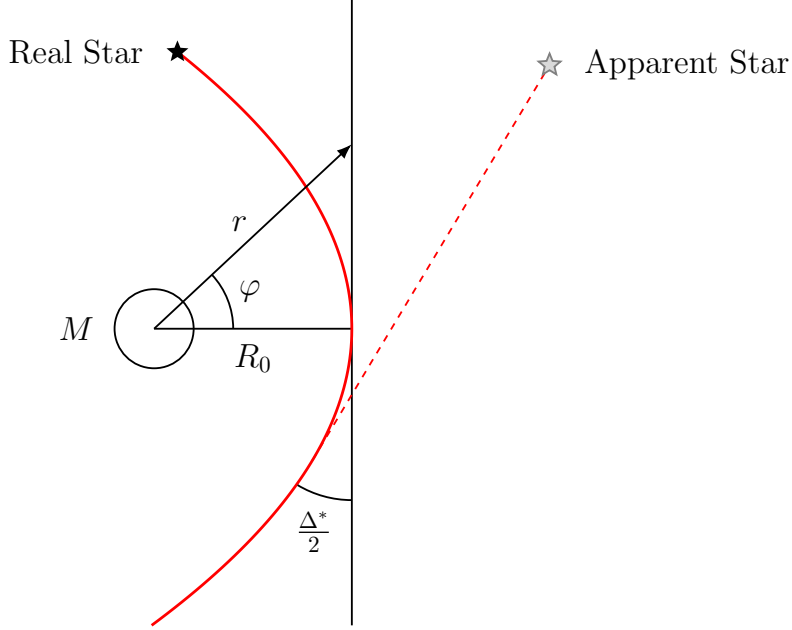


Figure 1: Trajectory of a photon deflected by a massive body according to Eq. (20) (red curved line) compared against an unperturbed free photon trajectory (vertical straight line).

6 Discussion and Conclusion

One of the longest-standing objectives in theoretical physics has been the seamless integration of gravitational theory with the framework of vector and gauge fields that successfully describes quantum mechanics. By decomposing the metric tensor into a set of fundamental vector fields, this study demonstrates that gravity can be mathematically restructured to heavily mirror electrodynamics. The linearization of the metric tensor inherently circumvents the pathological coordinate singularities found in conventional Schwarzschild geometry, presenting a robust analytical alternative for weak-field gravitational modeling.

Strikingly, substituting these fields into the Einstein field equations recovers a dynamic framework identical to Maxwell's equations and the Lorentz force. This structural isomorphism suggests profound underlying unities between spacetime geometry and electrodynamic propagation, opening speculative but mathematically rigorous avenues toward unified field theories and gravitational wave propagation mechanisms.

Practically, the derived spherically symmetric metric successfully predicts gravitational time dilation, spectral redshift, and the precise perihelion advance of Mercury. However, it notably deviates in its prediction of the gravitational deflection of light, suggesting 1.90 arcsec rather than the standard 1.75 arcsec. Historically, light deflection measurements have been plagued by systematic discrepancies. Observations relying on visible light (such as classical solar eclipse expeditions) frequently generated data that skewed higher, close to our predicted 1.90 arcsec [10]. The community eventually favored radio-interferometry measurements (VLBI), which aligned perfectly with the 1.75 arcsec Schwarzschild prediction [11, 12]. This vector field formulation poses an intriguing question: rather than mere instrumental error, could the historical discrepancy between visible light and radio wave deflection measurements hint at a deeper, wavelength-dependent dispersive quality of the gravitational field itself? Future high-precision astrometric missions may well provide the ultimate empirical arbiter for this vector approach.

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References

- [1] S. Carlip, “Quantum Gravity: A Progress Report”, *Reports on Progress in Physics*, vol. 64, no. 8, p. 885, 2001. <https://doi.org/10.1088/0034-4885/64/8/301>
- [2] C. W. Misner, K. P. Thorne, and J. A. Wheeler, *Gravitation*. W. H. Freeman and Company, 1973, pp. 216-218.
- [3] L. D. Landau and E. M. Lifshitz, *The Classical Theory of Fields* (Course of Theoretical Physics, Volume 2), 4th Ed., Butterworth-Heinemann, 1975, pp. 65-68.
- [4] J. D. Jackson, *Classical Electrodynamics*. John Wiley & Sons, 3rd Ed., 1999, pp. 608-610.
- [5] S. Weinberg, *Gravitation and Cosmology: Principles and Applications of the General Theory of Relativity*. John Wiley & Sons, 1972, pp. 175-180.
- [6] R. M. Wald, *General Relativity*. University of Chicago Press, 1984, pp. 119-125.
- [7] G. M. Clemence, “The Relativity Effect in Planetary Motions”. *Reviews of Modern Physics*, vol. 19, no. 4, pp. 361–364, 1947. <https://doi.org/10.1103/RevModPhys.19.361>
- [8] C. M. Will, “The Confrontation between General Relativity and Experiment”, *Living Reviews in Relativity*, vol. 17, no. 4, 2014. <https://doi.org/10.12942/lrr-2014-4>
- [9] G. 't Hooft, *Introduction to General Relativity*. Rinton Press, 2001, pp. 53-54.
- [10] Texas Mauritanian Eclipse Team, “Gravitational deflection of light: solar eclipse of 30 June 1973 I. Description of procedures and final result”. *The Astronomical Journal*, vol. 81, pp. 452-454, 1976. <https://doi.org/10.1086/111906>
- [11] H. C. Ohanian and R. Ruffini, *Gravitation and Spacetime*. Cambridge University Press, 3rd Ed., 2013, pp. 138-142.
- [12] E. Fomalont, S. Kopeikin, G. Lanyi, and J. Benson, “Progress in Measurements of the Gravitational Bending of Radio Waves Using the VLBA”, *The Astrophysical Journal*, vol. 699, no. 2, pp. 1395-1402, 2009. <https://doi.org/10.1088/0004-637X/699/2/1395>