

The Dirac point electron in zero-gravity Kerr–Newman spacetime

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Abstract

Dirac’s wave equation for a point electron in the electromagnetic Kerr–Newman spacetime is studied in a zero-gravity limit; here, “zero-gravity” means $G \rightarrow 0$, where G is Newton’s constant of universal gravitation. The zero- G limit studied here eliminates the troublesome Cauchy horizon of the Kerr–Newman spacetime and also its physically problematic acausal region of closed timelike loops. While the gravitational features of the Kerr–Newman manifold vanish as well when $G \rightarrow 0$, the limit retains the nontrivial topology associated with its ring singularity of radius $|a|$, and its electromagnetic Appell–Sommerfeld fields with charge Q and ADM magnetic dipole moment Qa . Moreover, one readily obtains (Q, I) -generalizations of this electromagnetic zero- G Kerr–Newman (zGKN) spacetime with same charge Q but with magnetic dipole moments $I\pi a^2$, where I is the equivalent electric current supported on a ring of radius $|a|$ in \mathbb{R}^3 which produces the same magnetic dipole moment in the far field multipole expansion. Dirac’s equation for a point electron in such spacetimes is commented on, too. The following results are obtained: First, by adapting an argument of Winklmeier–Yamada for the Dirac equation of a point electron in the outer region of the Kerr–Newman black hole spacetime, combined with a perturbation argument, it is shown that the formal Dirac Hamiltonian on the static spacelike slices of the maximal analytically extended zGKN spacetime is essentially self-adjoint. Next, with the help of an operator which anti-commutes with the Dirac Hamiltonian of a point electron in any (Q, I) -generalization of the zGKN spacetime, it is shown that the spectrum of any self-adjoint extension is symmetric about zero for any charge and magnetic dipole moment. Then, exploiting the Chandrasekhar–Page–Toop separation-of-variables technique for Dirac’s equation on a general Kerr–Newman spacetime it is shown that the Dirac operator on the zGKN spacetime has a continuous spectrum with a gap about zero that contains a point spectrum. The point spectrum is associated with time-periodic normalizable solutions, representing bound states of Dirac’s point electron in the electromagnetic field of the ring singularity of the zGKN spacetime.

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1 Introduction

In 1976 Chandrasekhar [11, 10], Page [25], and Toop [36] showed that Dirac’s equation for a point electron in the Kerr–Newman spacetime separates essentially completely¹ in oblate spheroidal coordinates. Although this remarkable discovery enabled detailed mathematical studies of the behavior of a Dirac electron in a charged, rotating black hole spacetime [20, 5, 14, 15, 16, 17, 3, 39, 40, 4] (see also [12] for neutral rotating black holes), there are perplexing conceptual issues which await clarification. Beside those that hark back to the enigmatic meaning of Dirac’s equation in Minkowski spacetime, see Thaller [34], serious new issues arise because of the physically somewhat questionable character of the Kerr–Newman solution, unveiled by Carter [8].

Namely, the maximal analytical extension of the stationary axisymmetric Kerr–Newman spacetime has a very strong curvature singularity on a timelike² cylindrical surface whose cross-section with constant- t hypersurfaces is a circle; here, t is a coordinate pertinent to the *asymptotically (at spacelike ∞) timelike* Killing field that encodes the stationarity of the “outer regions” of the Kerr–Newman spacetime. This circle is commonly referred to as *the “ring” singularity*. The region near the ring is especially pathological since it harbors closed timelike loops.³ Carter [8] also showed that the maximal analytically extended Kerr–Newman manifold is “cross-linked through the ring.”⁴ This non-trivial topology survives the vanishing-charge⁵ limit of the Kerr–Newman manifold, which yields the maximal analytic extension [6] of Kerr’s solution [21] to Einstein’s vacuum equations ($\mathbf{R}_{\mu\nu} = 0$), cf. [19]. Carter furthermore showed that this topology also survives the vanishing-mass⁶ limit of the Kerr manifold, which yields an otherwise flat vacuum spacetime consisting of two static spacetime ends which are cross-linked through the ring. This vanishing-mass limit of the Kerr manifold coincides with the vanishing-mass limit of the family of static vacuum spacetimes discovered, and completely described, a few years earlier by Zipoy [41]. Since Zipoy seems to have been the first to discover this non-trivial topology in exact spacetime solutions to Einstein’s vacuum equations, we henceforth will refer to it as the *Zipoy topology*.

In the black-hole sector of their parameter space the Kerr–Newman spacetimes also have a Cauchy horizon and an event horizon. From the “safe perspective of an observer at spatial infinity” the ring singularity, the acausal region, and the Cauchy horizon are invisible, being “hidden” behind the event horizon, and no exotic or even objectionable physics would ever seem to happen: a Dirac spinor wavefunction initially supported outside the event horizon will either keep spreading within the outer region or eventually (as $t \rightarrow \infty$) accumulate (in parts or wholly) at the event horizon, see [16, 17]. However, mathematical (and theoretical) physicists may also want to study the spinor wave function in other coordinates designed to “follow it across the event horizon,”⁷ yet

¹In contrast to the familiar separation-of-variables results for, say, the Laplacian in a rectangular box or a cylinder or a sphere, Chandrasekhar, Page, and Toop obtained a system of ODEs for functions of only one variable each, but each of these ODEs is coupled to another one through some shared parameters.

²In a limiting sense of course, since the metric is singular on this surface.

³The timelike ring singularity of the Kerr–Newman manifold is itself the limit of closed timelike loops, for which reason it is not possible to interpret this singular source of the stationary and axisymmetric Kerr–Newman electromagnetic fields outside of the event horizon as a “rotating charged ring.”

⁴The complement of a wedding ring in ordinary three-dimensional Euclidean space is topologically non-trivial, too, but “looping through the ring once brings you back to where you began;” in a spacelike slice of the maximal analytically extended Kerr–Newman spacetime “you need to loop through the ring twice to get back to square one.”

⁵By “charge” we mean the usual charge parameter in the solution.

⁶By “mass” we mean the usual mass parameter in the solution.

⁷When the analogous study was carried out by Oppenheimer and Snyder [24] for classical gaseous matter undergoing gravitational collapse it leveled the ground for building our modern understanding of the physics of gravitational collapse, involving the formation of black holes and their singularities. Poetically speaking their work revealed that there is more physics in general relativity than meets the (distant observers) eye.

it is neither clear how to continue in a “physically correct” manner beyond the Cauchy horizon⁸ nor what to make of the acausal region of closed timelike loops, nor how to correctly handle the timelike singularity. Moreover, if one inquires into the “physics beyond the event horizon,” one has the option of allowing the support of the initial spinor wave function to be spread over both asymptotically flat ends, or some other parts of the maximal analytically extended Kerr–Newman spacetime. It is not so clear which options (if any) are physically reasonable and which ones are science fiction, although astrophysicists can argue for the “physical black hole spacetime,” i.e. the part of the maximal analytically extended Kerr–Newman spacetime which is the asymptotic limit of the topologically simple spacetime of a charged, rotating star collapsing into a black hole [30].

The horizons are absent in the hyper-extremal parameter regime. Even though the absence of the Cauchy horizon is a welcome simplification, this regime is rarely studied because the absence of the event horizon renders the singularity “naked,” and the (weak) cosmic censorship hypothesis, according to which “nature abhors naked singularities” [26, 27], has (unfortunately) discouraged physicists from investigating spacetimes with naked singularities. Yet once a piece of a Dirac spinor wave function has crossed the event horizon of a Kerr–Newman black hole it is no longer shielded from possible harm done by the spacetime singularity — viz., inside the event horizon the singularity *is* naked —, and so one may as well study the effects of naked singularities directly. Be that as it may, the hyper-extremal regime retains the closed timelike loops which according to the standard interpretation of general relativity turn the entire manifold into a *causally vicious set*, something that many physicists (including the authors) would regard as physically suspicious.

A strategy to rid the Kerr–Newman manifold from its Cauchy horizon, and all its other acausal aspects, is to take a *zero-gravity limit*, which means taking $G \rightarrow 0$, where G is Newton’s constant of universal gravitation. This would be quite uninteresting if the zero- G limit of the Kerr–Newman manifold would simply yield a Minkowski spacetime decorated with the electric field of a point charge and the magnetic field of a point dipole, as one might be tempted to guess from the asymptotically flat ends of the Kerr–Newman spacetime. However, as shown by one of us in the accompanying paper [32], the zero-gravity limit of the maximal analytically extended Kerr–Newman spacetime yields a static, flat, yet *two-leafed, cross-linked* spacetime which is decorated with Appell–Sommerfeld [2, 29] electromagnetic fields⁹ whose sources appear to be certain “finite charge and current distributions” supported by the one-dimensional ring.¹⁰ Although the gravitational (viz.: curvature) aspects of the Kerr–Newman manifold, its event horizon included, vanish in this limit too, one does retain the topological, the singular, and all the electromagnetic aspects of the spacetime. Studies of the Dirac equation for a point electron in this zero- G Kerr–Newman (zGKN) spacetime will therefore illuminate the role of the topological and electromagnetic aspects of the Kerr–Newman manifold in regard to the relativistic quantum mechanics of the electron.

In this paper we study the Dirac equation for a point electron in static, electromagnetic, flat spacetimes with Zipoy topology which include the zGKN spacetimes as special case, but which in general can sport any Sommerfeld fields one wants. We will consider Sommerfeld fields which differ

⁸If instead of the Cauchy problem one studies t -periodic solutions, then one can continue across the Cauchy and the event horizons using a weak matching procedure [14, 15]. However, Finster et al. [14, 15] found that no t -periodic solutions exist which are normalized over a constant- t slice of the “physical black hole spacetime;” see main text.

⁹The zero- G limit of the electromagnetic Kerr–Newman fields yields fields originally discovered by Appell [2], who obtained them from the Coulomb potential of a point charge by a complex translation of the charge’s position. Appell noticed that the fields change sign when looping once through the ring, but did not conclude — apparently — that they live naturally on a topologically nontrivial space. Sommerfeld [29] seems to have been first to introduce “branched Riemann spaces,” three-dimensional analogues of topologically non-trivial Riemann surfaces, to which we will refer as *Sommerfeld spaces*, and to construct electromagnetic fields (harmonic functions) on them, which in general we will call *Sommerfeld fields*. Eventually Evans [13] and his students [23], [1] laid their rigorous foundations.

¹⁰Strictly speaking, the ring singularity is not part of the manifold; it’s rather a ring “defect.”

from the Appell–Sommerfeld fields only in a single number, the ratio $I\pi a^2/Qa$ of their magnetic dipole moment to the magnetic dipole moment of the Appell–Sommerfeld fields of same charge Q ; here, $|a|$ is the radius of the ring singularity. By constructing an operator that anti-commutes with the pertinent Dirac Hamiltonian we show that the spectrum of any of its self-adjoint extensions is symmetric about zero; this result holds for arbitrary (Q, I) . All other results are obtained for Dirac’s electron in the $zGKN$ spacetime ($I\pi a^2/Q|a| = 1$): by adapting an argument of Winklmeier–Yamada for the Dirac equation of a point electron in the outer region of the Kerr–Newman black hole spacetime, we show that our formal Dirac Hamiltonian is essentially self-adjoint on a complete spacelike slice of the maximal analytically extended, static $zGKN$ spacetime. Then we exploit the Chandrasekhar–Page–Toop separation-of-variables theorem for Dirac’s equation on a general Kerr–Newman spacetime, and the Prüfer transform, to show that the self-adjoint Dirac operator has a continuous spectrum with a gap about zero that contains a pure point spectrum associated with time-periodic L^2 spinor fields, representing bound states of Dirac’s point electron in the electromagnetic field of the ring singularity of the $zGKN$ spacetime.

In the next section we formulate our main results about the Dirac equation for a point electron in the $zGKN$ spacetime; one result is valid also for a Dirac electron in static, flat spacetimes having Zipoy topology featuring electromagnetic Sommerfeld fields of arbitrary $I\pi a/Q$ -ratio. In sections 3, 4, 5, 6, 7, and 8 we prove our main theorems about the spectrum. In section 9 we conclude with a list of interesting questions left unanswered by this work.

2 Formulation of the main results

We begin by formulating the Dirac equation for a point electron in electromagnetic, static, flat spacetimes with Zipoy topology which generalize zero- G Kerr–Newman spacetimes to zero- G Kerr spacetimes equipped with Sommerfeld fields of arbitrary $I\pi a/Q$ -ratio. We then state our main theorems about the spectrum of the pertinent Dirac operators.

2.1 Dirac’s equation for a point electron on zero- G Kerr spacetimes equipped with electromagnetic Sommerfeld fields of arbitrary $I\pi a/Q$ -ratio

2.1.1 Zero- G Kerr spacetimes

Our limit $G \rightarrow 0$ of the maximal analytic extension of the well-known *Kerr* family of stationary, axisymmetric spacetime solutions of Einstein’s vacuum equations yields a one-parameter family of static, flat, but topologically nontrivial spacetimes¹¹ $(\mathcal{M}, \mathbf{g})$ which consist of two “cross-linked leafs.” Explicitly, let $\mathcal{C} \equiv \{(t, r, \theta, \varphi) : t \in \mathbb{R}, r \in \mathbb{R}, \theta \in [0, \pi], \varphi \in [0, 2\pi)\}$ denote a rectangular “four-dimensional cylinder,” and let $\mathcal{S} \equiv \{(t, r, \theta, \varphi) : t \in \mathbb{R}, r = 0, \theta = \pi/2, \varphi \in [0, 2\pi)\} \subset \mathcal{C}$ denote a rectangular “two-dimensional slab” in \mathcal{C} . Then $\mathcal{C} \setminus \mathcal{S}$ is a covering chart of oblate spheroidal (Boyer–Lindquist, or BL) coordinates¹² for this spacetime, with line element¹³

$$ds_{\mathbf{g}}^2 = dt^2 - (r^2 + a^2) \sin^2 \theta d\varphi^2 - \frac{r^2 + a^2 \cos^2 \theta}{r^2 + a^2} (dr^2 + (r^2 + a^2) d\theta^2); \quad (1)$$

here, $a^2 > 0$ is the only parameter of these spacetimes, and we have set the speed of light $c = 1$.

¹¹We note that Zipoy [41] found a large class of static, axisymmetric, flat but topologically nontrivial solutions to the Einstein vacuum equations which in their zero- G limit coincides with the zero- G Kerr spacetime family.

¹²It is unfortunate that the notation (t, r, θ, φ) for the spacetime coordinates is somewhat overused in the relativity literature, for instance meaning also Schwarzschild coordinates on the outer region of that spacetime, but of course also just standard spherical coordinates of Minkowski spacetime. All standard non-spherical (t, r, θ, φ) coordinate systems reduce to standard spherical coordinates of Minkowski spacetime near “ $r = \infty$,” which is perhaps part of the original rationale for choosing this notation. Note that in BL coordinates r takes both positive and negative values.

¹³Note our sign convention of $(+, -, -, -)$ for the metric.

The static, axisymmetric character of these zero- G Kerr spacetimes is manifest in (1). Also, since $r \in \mathbb{R}$ occurs strictly quadratically in (1), it is clear that the manifold consists of two “conjoined identical twins.” To exhibit their flatness, and in the process also the topological nontrivial juncture, we introduce cylindrical coordinates (t, ϱ, z, φ) on Minkowski spacetime $\mathbb{R}^{1,3}$, with the same (t, φ) as in Boyer–Lindquist coordinates, and with the cylindrical coordinates (ϱ, z) related to the elliptical coordinates (r, θ) by

$$\varrho = \sqrt{r^2 + a^2} \sin \theta, \quad z = r \cos \theta. \quad (2)$$

In cylindrical (t, ϱ, z, φ) coordinates the metric takes the familiar form for flat Minkowski spacetime

$$ds_{\mathbf{g}}^2 = dt^2 - d\varrho^2 - \varrho^2 d\varphi^2 - dz^2, \quad (3)$$

except that the map (2) makes it plain that the chart $\mathcal{C} \setminus \mathcal{S}$ will be mapped into *two* copies of Minkowski spacetime which are “doubly conjoined,” in a smooth yet crossing manner, at the interior of the set \mathcal{S} . The metric \mathbf{g} given by the line element (1) has a singularity at \mathcal{S} , which is the singularity *of the spacetime*, not *in the spacetime*. The set \mathcal{S} is the boundary of a timelike open solid cylinder; the cross section at any instant t of this cylindrical surface is a translate of the *ring* $\mathcal{R}_0 \equiv \{(t, r, \theta, \varphi) : t = 0, r = 0, \theta = \pi/2, \varphi \in [0, 2\pi)\}$, for which reason one speaks of a “ring singularity.” The points on the ring are *conical singularities* for the metric, meaning that the limit as the radius goes to zero of the ratio of the circumference to radius of a small circle centered at a point of the ring and lying in a meridional plane $\varphi = \text{const.}$ is 4π instead of 2π . See [32] for details.

The key topological features of this manifold can easily be visualized. Namely, although the fixed timelike planes $\{(t, r, \theta, \varphi) : t = t_0, \varphi = \varphi_0\}$ cannot be embedded into \mathbb{R}^3 , each such plane can be immersed in it, the immersion consisting of two Euclidean half planes, stacked up upon each other, then cut along a line segment of length $|a|$ orthogonal to the planes’ boundaries “with scissors,” then smoothly “cross-glued” at the cut such that the “upper” and “lower” sheets are cross-linked like an \times along the cut, while remaining like \parallel beyond the cut; the singular endpoint of the \times -line is not part of the two-sheeted manifold. Shown in Fig.1 are the ring singularity and the part $\{r \in (-1, 1), \theta \in (0, \pi)\}$ of a constant-azimuth section (slightly curved, for the purpose of visualization) of the two-sheeted static spacelike slice of the zero- G Kerr spacetime.

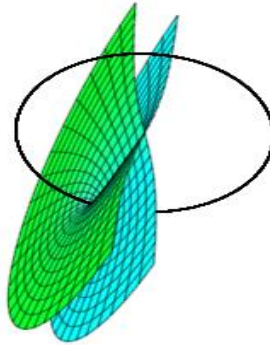


Figure 1: An illustration of the Zipoy topology.

We can view \mathcal{M} as a bundle over the base manifold $\mathbb{R}^{1,3} \setminus \mathcal{S}$ (mildly abusing notation), with the projection map $\Pi : \mathcal{M} \rightarrow \mathbb{R}^{1,3} \setminus \mathcal{S}$ being $\Pi(t, r, \theta, \varphi) = (t, \varrho, z, \varphi)$. The fiber over a point in the base consists of two points, degenerating into one point at each “ring”

$$\mathcal{R}_{t_0} = \{(t, r, \theta, \varphi) \mid t = t_0, r = 0, \theta = \pi/2, 0 \leq \varphi \leq 2\pi\}. \quad (4)$$

The pullback of the Minkowski metric η under Π endows \mathcal{M} with a flat Lorentzian metric $\mathbf{g} = \Pi^*\eta$, whose line element is given in (1).

2.1.2 Zero-G Kerr-Newman spacetimes

The spacetime $(\mathcal{M}, \mathbf{g})$ introduced above can be decorated with any static electromagnetic Sommerfeld field $\mathbf{F} = d\mathbf{A}$, satisfying the flat space Maxwell equations locally but respecting the topologically nontrivial character of the spacetime. The zero- G limit of the maximal analytical extension of the *Kerr–Newman* family¹⁴ of stationary axisymmetric solutions to the Einstein–Maxwell equations, written in BL coordinates, yields precisely the zero- G Kerr (zGK) spacetime decorated with a particular electromagnetic Sommerfeld field, the *Appell–Sommerfeld field*, whose four-potential one-form reads

$$\mathbf{A} = -\frac{r}{r^2 + a^2 \cos^2 \theta} (Qdt - Qa \sin^2 \theta d\varphi). \quad (5)$$

Here, Q is the total charge “seen from infinity” in one of the two sheets, defined by computing the electric outward flux through a spherical surface surrounding the ring singularity in that sheet. The field \mathbf{F} is singular on the same ring \mathcal{R}_t as is the metric, while for r very large positive its electric and magnetic components approach, respectively, the asymptotics of an “electric monopole field in \mathbb{R}^3 of a charge Q ” and a “magnetic dipole field in \mathbb{R}^3 of dipole moment Qa ,” for $-r$ very large positive, in the other sheet, these electric monopole and magnetic dipole fields correspond to a charge $-Q$ and magnetic dipole moment $-Qa$.

Remark 2.1. It may be tempting to speculate whether the magnetic dipole moment Qa can be interpreted as due to a “gyrating charged ring,” with a the angular momentum per unit mass “of the singularity,” as has been attempted for Kerr–Newman spacetimes. Moreover, since Kerr–Newman spacetimes have a *gyromagnetic ratio* $Q/M =: g_{\text{KN}}Q/2M$ (in units with $c = 1$), amounting to a g -factor $g_{\text{KN}} = 2$ (see [8]), and since the KN parameters (M, Q, a) are independent of G , and so is the KN gyromagnetic ratio, one could be tempted to assign the zGKN spacetime the same gyromagnetic ratio of Q/M and g -factor of 2. However, since M does not show in the zGKN metric, such an assignment would be reasonable only if there were no other way to construct zGKN than taking the zero- G limit of KN. Yet this is not the case: as already pointed out in footnote 11, the underlying spacetime manifold of zGKN can be obtained as zero- G limit of either, the *stationary* family of Kerr spacetimes — having both an ADM mass M_{K} and ADM angular momentum $J = M_{\text{K}}a$ —, or a *static* family of Zipoy spacetimes — having an ADM mass M_{Z} but no ADM angular momentum; note also that $M_{\text{Z}} \neq M_{\text{K}}$ in general. This (say) zGZ spacetime can now be equipped with an arbitrary Sommerfeld field, in particular: the Appell–Sommerfeld field of zGKN, without being logical compelled to interpret its magnetic moment Qa as being due to a “gyrating ring of charge Q ” with angular momentum per unit mass a ; this will become even more clear in the next subsection. So it’s better to refrain from assigning zGKN any spacetime g -factor. For a careful analysis of the ring sources of the electromagnetic zGKN fields, see [32].

2.1.3 Generalizations of zGKN spacetimes to arbitrary charge and current.

Since the electric and magnetic components of Maxwell’s vacuum equations decouple in the zero- G limit, to decorate the zero- G Kerr spacetime with a generalization of the electromagnetic Appell–Sommerfeld field $\mathbf{F} = d\mathbf{A}$ having electric charge Q and current I , all that needs to be done is to

¹⁴We recall that for fixed $G > 0$ (and speed of light c) the Kerr–Newman family is a three-parameter family of electrovac spacetimes, the parameters being ADM mass (energy) $M(> 0)$, ADM angular momentum $J = Ma \in \mathbb{R}$, and total charge $Q \in \mathbb{R}$, all defined in a single asymptotic end. Note that in units where $c = 1$ the angular momentum per unit mass a has physical dimension of length; equivalently, of time.

replace the magnetic dipole moment Qa in formula (5) by $I\pi a^2$, thus

$$\mathbf{A} = -\frac{r}{r^2 + a^2 \cos^2 \theta} (Qdt - I\pi a^2 \sin^2 \theta d\varphi). \quad (6)$$

Again, electric charge Q and magnetic dipole moment $I\pi a^2$ are as “seen” from spacelike infinity in one sheet; viewing from spacelike infinity in the other sheet one “sees” $-Q$ and $-I\pi a^2$.

Remark 2.2. The Sommerfeld field (6) makes it plain that the zGKN spacetimes are but a special one-parameter subfamily in a two-parameter family of qualitatively similar spacetimes with arbitrary charge Q and magnetic moment $I\pi a^2$ (given a). The ease with which this result was accomplished stands in stark contrast to the difficulties in generalizing the Kerr–Newman family to electromagnetic spacetimes with a magnetic moment different from Qa .

2.1.4 The Dirac equation on electromagnetic spacetimes: Cartan’s frame method

In arbitrary coordinates (x^μ) (with $c = 1$ and $\hbar = 1$), the Dirac equation for a spin-1/2 electron of empirical rest mass m and charge $-e < 0$ interacting (through minimal coupling) with an electromagnetic field $\mathbf{F} = d\mathbf{A}$ in a spacetime $(\mathcal{M}, \mathbf{g})$ reads

$$\tilde{\gamma}^\mu (-i\nabla_\mu + eA_\mu)\Psi + m\Psi = 0; \quad (7)$$

here ∇ is the covariant derivative (on spinors) associated to the spacetime metric \mathbf{g} , and $(\tilde{\gamma}^\mu)_{\mu=0}^3$ are Dirac matrices associated to this metric, i.e. satisfying

$$\tilde{\gamma}^\mu \tilde{\gamma}^\nu + \tilde{\gamma}^\nu \tilde{\gamma}^\mu = 2\mathbf{g}^{\mu\nu} \mathbf{1}_{4 \times 4}, \quad (8)$$

while the A_μ are the pertinent components of the electromagnetic potential, $\mathbf{A} = A_\mu dx^\mu$.

Using Cartan’s frame method (see [7] and refs. therein) one can express the above covariant derivative on spinors in terms of standard derivatives:

$$\tilde{\gamma}^\mu \nabla_\mu = \gamma^\mu \mathbf{e}_\mu + \frac{1}{4} \Omega_{\mu\nu\lambda} \gamma^\lambda \gamma^\mu \gamma^\nu \quad (9)$$

Here $\{\mathbf{e}_\mu\}_{\mu=0}^3$ is a *Cartan frame*, i.e. an orthonormal frame of vectors spanning the tangent space at each point of the spacetime manifold. We thus have

$$(\mathbf{e}_\mu)^\nu (\mathbf{e}_\lambda)^\kappa \mathbf{g}_{\nu\kappa} = \eta_{\mu\lambda}. \quad (10)$$

On the one hand, it follows that

$$\tilde{\gamma}^\mu = (\mathbf{e}_\nu)^\mu \gamma^\nu, \quad (11)$$

where the γ^ν are Dirac gamma matrices for the Minkowski space, satisfying $\gamma^\nu \gamma^\mu + \gamma^\mu \gamma^\nu = 2\eta^{\mu\nu} \mathbf{1}_{4 \times 4}$. On the other hand, let $\{\boldsymbol{\omega}^\mu\}_{\mu=0}^3$ denote the *dual* frame to $\{\mathbf{e}_\mu\}$, i.e. the orthonormal basis for the cotangent space at each point of the manifold that is dual to the basis for the tangent space:

$$\boldsymbol{\omega}_\mu (\mathbf{e}^\nu) = \mathbf{e}^\nu (\boldsymbol{\omega}_\mu) = \delta_\mu^\nu. \quad (12)$$

Then the $\Omega_{\mu\nu\lambda}$ are by definition the *Ricci rotation coefficients* of the frame $\{\boldsymbol{\omega}^\mu\}_{\mu=0}^3$, defined in the following way: Let the one-forms Ω_ν^μ satisfy

$$d\boldsymbol{\omega}^\mu + \Omega_\nu^\mu \wedge \boldsymbol{\omega}^\nu = 0. \quad (13)$$

This does not uniquely define the Ω_ν^μ . However, there exists a unique set of such 1-forms satisfying the extra condition

$$\Omega_{\mu\nu} = -\Omega_{\nu\mu}, \quad (14)$$

where the first index is lowered by the Minkowski metric: $\Omega_{\mu\nu} := \eta_{\mu\lambda}\Omega_\nu^\lambda$. Since $\{\omega^\mu\}$ forms a basis for the space of 1-forms, we then have $\Omega_{\mu\nu} = \Omega_{\mu\nu\lambda}\omega^\lambda$, which defines the rotation coefficients $\Omega_{\mu\nu\lambda}$.

The Dirac equation (7) on a spacetime $(\mathcal{M}, \mathbf{g})$ with an electromagnetic 4-potential \mathbf{A} can thus be written in the following form:

$$\gamma^\mu \left(\mathbf{e}_\mu + \Gamma_\mu + ie\tilde{A}_\mu \right) \Psi + im\Psi = 0; \quad (15)$$

here, the Γ_μ are connection coefficients,

$$\Gamma_\mu := \frac{1}{4}\Omega_{\nu\lambda\mu}\gamma^\nu\gamma^\lambda = \frac{1}{8}\Omega_{\nu\lambda\mu}[\gamma^\nu, \gamma^\lambda], \quad (16)$$

and the \tilde{A}_μ are the components of the potential \mathbf{A} in the ω^μ basis, i.e. $\mathbf{A} = \tilde{A}_\mu\omega^\mu$, or,

$$\tilde{A}_\mu := (\mathbf{e}_\mu)^\nu A_\nu. \quad (17)$$

2.1.5 Frame formulation of the Dirac equation on zero- G Kerr spacetimes featuring electromagnetic Appell–Sommerfeld fields with arbitrary $I\pi a/Q$ -ratio

As explained in section 2.1.1, the single chart $\mathcal{C} \setminus \mathcal{S}$ of oblate spheroidal coordinates (t, r, θ, φ) covers the whole zero- G Kerr spacetime $(\mathcal{M}, \mathbf{g})$, and in section 2.1.2 we saw that in these coordinates the electromagnetic Appell–Sommerfeld one-form \mathbf{A} is everywhere on $(\mathcal{M}, \mathbf{g})$ given by the simple formula (6). It is therefore only natural that one would like to write Dirac’s equation (7) in these coordinates as well, in the hope of achieving at least some partial separation of variables.¹⁵

However, unlike Cartesian coordinates (x^μ) in Minkowski spacetime, oblate spheroidal coordinate derivatives do not give rise to an orthonormal basis for the tangent space at each point of a zero- G Kerr spacetime. Thus, to bring (7) into the Cartan form (15) using oblate spheroidal coordinates, one also needs to construct a suitable Cartan frame. Following Chandrasekhar [11, 10], Page [25], Toop [36] (see also Carter-McLenaghan [9]), we introduce a special orthonormal frame $\{\mathbf{e}_\mu\}_{\mu=0}^3$ on the tangent bundle $T\mathcal{M}$ which is adapted to the oblate spheroidal coordinates in order for the Dirac equation to take a comparatively simple form.

We begin by introducing a Cartan (co-)frame $\{\omega^\mu\}_{\mu=0}^3$ for the cotangent bundle¹⁶:

$$\omega^0 := \frac{\Delta}{|\rho|}(dt - a \sin^2 \theta d\varphi), \quad \omega^1 := |\rho|d\theta, \quad \omega^2 := \frac{\sin \theta}{|\rho|}(-adt + \Delta^2 d\varphi), \quad \omega^3 := \frac{|\rho|}{\Delta}dr, \quad (18)$$

with the conventional abbreviations

$$\Delta := \sqrt{r^2 + a^2}, \quad \rho := r + ia \cos \theta. \quad (19)$$

Let us denote the oblate spheroidal coordinates (t, r, θ, φ) collectively by (y^ν) . Let $\mathbf{g}_{\mu\nu}$ denote the coefficients of the spacetime metric (1) in oblate spheroidal coordinates, i.e. $\mathbf{g}_{\mu\nu} = \mathbf{g}\left(\frac{\partial}{\partial y^\mu}, \frac{\partial}{\partial y^\nu}\right)$. One easily checks that written in the $\{\omega^\mu\}$ frame, the spacetime line element is

$$ds_{\mathbf{g}}^2 = \mathbf{g}_{\mu\nu}dy^\mu dy^\nu = \boldsymbol{\eta}_{\alpha\beta}\omega^\alpha\omega^\beta, \quad (20)$$

¹⁵The idea of using special frames adapted to a coordinate system in order to separate spinorial wave equations in those coordinates goes back to Kinnersley [22] and Teukolsky [33].

¹⁶This particular frame is called a *canonical symmetric tetrad* in [9].

where

$$\boldsymbol{\eta} = \text{diag}(1, -1, -1, -1) \quad (21)$$

is the Minkowski metric in rectangular coordinates. This shows that the frame $\{\boldsymbol{\omega}^\mu\}_{\mu=0}^3$ is indeed orthonormal. With respect to this frame the electromagnetic Sommerfeld potential (6) becomes $\mathbf{A} = \tilde{A}_\mu \boldsymbol{\omega}^\mu$, with

$$\tilde{A}_0 = -Q \frac{r}{|\rho|\Delta} - (Q - I\pi a) \frac{a^2 r \sin^2 \theta}{\Delta |\rho|^3}, \quad \tilde{A}_1 = 0, \quad \tilde{A}_2 = -(Q - I\pi a) \frac{ar \sin \theta}{|\rho|^3}, \quad \tilde{A}_3 = 0. \quad (22)$$

Remark 2.3. We observe that for $Q = I\pi a$, all but one of the quantities \tilde{A}_μ vanish, and the non-vanishing one, \tilde{A}_0 , reduces to $-Qr/|\rho|\Delta$.

Next, let the frame of vector fields $\{\mathbf{e}_\mu\}$ be the *dual* frame to $\{\boldsymbol{\omega}^\mu\}$. Thus $\{\mathbf{e}_\mu\}$ yields an orthonormal basis for the tangent space at each point in the manifold:

$$\mathbf{e}_0 = \frac{\Delta}{|\rho|} \partial_t + \frac{a}{\Delta |\rho|} \partial_\varphi, \quad \mathbf{e}_1 = \frac{1}{|\rho|} \partial_\theta, \quad \mathbf{e}_2 = \frac{a \sin \theta}{|\rho|} \partial_t + \frac{1}{|\rho| \sin \theta} \partial_\varphi, \quad \mathbf{e}_3 = \frac{\Delta}{|\rho|} \partial_r. \quad (23)$$

Next, the anti-symmetric matrix $(\Omega_{\mu\nu}) = (\boldsymbol{\eta}_{\mu\lambda} \Omega_\nu^\lambda)$ is computed to be

$$(\Omega_{\mu\nu}) = \begin{pmatrix} 0 & -C\boldsymbol{\omega}^0 - D\boldsymbol{\omega}^2 & D\boldsymbol{\omega}^1 - B\boldsymbol{\omega}^3 & -A\boldsymbol{\omega}^0 - B\boldsymbol{\omega}^2 \\ & 0 & D\boldsymbol{\omega}^0 + F\boldsymbol{\omega}^2 & -E\boldsymbol{\omega}^1 - C\boldsymbol{\omega}^3 \\ & (\text{asym}) & 0 & -B\boldsymbol{\omega}^0 - E\boldsymbol{\omega}^2 \\ & & & 0 \end{pmatrix}, \quad (24)$$

with

$$A := \frac{a^2 r \sin^2 \theta}{\Delta |\rho|^3}, \quad B := \frac{ar \sin \theta}{|\rho|^3}, \quad C := \frac{a^2 \sin \theta \cos \theta}{|\rho|^3}, \quad D := \frac{a \cos \theta \Delta}{|\rho|^3}, \quad E := \frac{r \Delta}{|\rho|^3}, \quad F := \frac{\Delta^2 \cos \theta}{|\rho|^3 \sin \theta}. \quad (25)$$

With respect to this frame on a zero- G Kerr spacetime the covariant derivative part of the Dirac operator (7) can be expressed with the help of the operator

$$\mathfrak{D} := \tilde{\gamma}^\mu \nabla_\mu = \begin{pmatrix} 0 & \mathfrak{l}' + \mathfrak{m}' \\ \mathfrak{l} + \mathfrak{m} & 0 \end{pmatrix}, \quad (26)$$

where

$$\mathfrak{l} := \frac{1}{|\rho|} \begin{pmatrix} D_+ & L_- \\ L_+ & D_- \end{pmatrix} \quad (27)$$

and

$$\mathfrak{l}' := \frac{1}{|\rho|} \begin{pmatrix} D_- & -L_- \\ -L_+ & D_+ \end{pmatrix}, \quad (28)$$

with

$$D_\pm := \pm \Delta \partial_r + \left(\Delta \partial_t + \frac{a}{\Delta} \partial_\varphi \right), \quad L_\pm := \partial_\theta \pm i (a \sin \theta \partial_t + \csc \theta \partial_\varphi), \quad (29)$$

while

$$\begin{aligned} \mathfrak{m} &:= \frac{1}{2} [(-2C + F + iB)\sigma_1 + (-A + 2E + iD)\sigma_3] \\ &= \frac{1}{2|\rho|} \begin{pmatrix} \frac{r}{\Delta} + \frac{\Delta}{\bar{\rho}} & \cot \theta + \frac{ia \sin \theta}{\bar{\rho}} \\ \cot \theta + \frac{ia \sin \theta}{\bar{\rho}} & -\frac{r}{\Delta} - \frac{\Delta}{\bar{\rho}} \end{pmatrix} \end{aligned} \quad (30)$$

and

$$\mathbf{m}' := \frac{1}{2}[(2C - F + iB)\sigma_1 + (A - 2E + iD)\sigma_3] = -\mathbf{m}^*, \quad (31)$$

where the σ_k are Pauli matrices:

$$\sigma_1 = \begin{pmatrix} 0 & 1 \\ 1 & 0 \end{pmatrix}, \quad \sigma_2 = \begin{pmatrix} 0 & -i \\ i & 0 \end{pmatrix}, \quad \sigma_3 = \begin{pmatrix} 1 & 0 \\ 0 & -1 \end{pmatrix}. \quad (32)$$

We note that the principal part of $|\rho|\mathfrak{D}$ has an additive separation property:

$$|\rho| \begin{pmatrix} 0 & l' \\ l & 0 \end{pmatrix} = \left[\gamma^3 \Delta \partial_r + \gamma^0 \left((\Delta \partial_t + \frac{a}{\Delta} \partial_\varphi) \right) \right] + \left[\gamma^1 \partial_\theta + \gamma^2 (a \sin \theta \partial_t + \csc \theta \partial_\varphi) \right], \quad (33)$$

where the coefficients of the two square-bracketed operators are functions of only r , respectively only θ . Moreover, it is possible to transform away the lower order term in \mathfrak{D} , so that exact separation can be achieved for $|\rho|\mathfrak{D}$. Namely, let

$$\chi(r, \theta) := \frac{1}{2} \log(\Delta \bar{\rho} \sin \theta). \quad (34)$$

It is easy to see that

$$\mathbf{m} = l\chi, \quad \mathbf{m}' = l'\bar{\chi}. \quad (35)$$

Let us therefore define the diagonal matrix

$$\mathfrak{D} := \text{diag}(e^{-\chi}, e^{-\chi}, e^{-\bar{\chi}}, e^{-\bar{\chi}}) \quad (36)$$

and a new bispinor $\hat{\Psi}$ related to the original Ψ by

$$\Psi = \mathfrak{D} \hat{\Psi}. \quad (37)$$

Denoting the upper and lower components of a bispinor Ψ by ψ_1 and ψ_2 respectively, it then follows that

$$(l + \mathbf{m})\psi_1 = (l + \mathbf{m})(e^{-\chi} \hat{\psi}_1) = e^{-\chi} [l - l\chi + \mathbf{m}] \hat{\psi}_1 = e^{-\chi} l \hat{\psi}_1, \quad (38)$$

and similarly

$$(l' + \mathbf{m}')\psi_2 = e^{-\bar{\chi}} l' \hat{\psi}_2. \quad (39)$$

We now put it all together. We set

$$\mathfrak{R} := \text{diag}(\rho, \rho, \bar{\rho}, \bar{\rho}) \quad (40)$$

and note that $|\rho|\mathfrak{D}^{-*}\mathfrak{D} = \mathfrak{R}$ while $\mathfrak{D}^{-*}\gamma^\mu\mathfrak{D} = \gamma^\mu$. Thus, setting $\Psi = \mathfrak{D}\hat{\Psi}$ in (7) and left-multiplying the equation by the diagonal matrix $\mathfrak{D}' := |\rho|\mathfrak{D}^{-*}$ we conclude that $\hat{\Psi}$ solves a new Dirac equation

$$\left(|\rho|\gamma^\mu (\mathbf{e}_\mu + ie\tilde{A}_\mu) + im\mathfrak{R} \right) \hat{\Psi} = 0. \quad (41)$$

Finally, let us compute the Hamiltonian form of (41). Let matrices M^μ be defined by

$$|\rho|\gamma^\mu \mathbf{e}_\mu = M^\mu \partial_\mu. \quad (42)$$

Thus in particular

$$M^0 = \Delta \gamma^0 + a \sin \theta \gamma^2. \quad (43)$$

We may thus rewrite (41) as

$$M^0 \partial_t \hat{\Psi} = - \left(M^k \partial_k + ie|\rho|\gamma^\mu \tilde{A}_\mu + im\mathfrak{R} \right) \hat{\Psi}, \quad (44)$$

so that, defining

$$\hat{H} := -i(M^0)^{-1} \left(M^k \partial_k + ie|\rho|\gamma^\mu \tilde{A}_\mu + im\mathfrak{R} \right), \quad (45)$$

we can now rewrite the Dirac equation (41) in Hamiltonian form:

$$i\partial_t \hat{\Psi} = \hat{H} \hat{\Psi}. \quad (46)$$

Remark 2.4. We note that for $Q \neq I\pi a$ the quantity $|\rho|\gamma^\mu \tilde{A}_\mu$ is a function of both r and θ , and unlike the other terms in the Dirac equation (41) it does not separate into a sum of two terms each depending only on one of these variables. It follows that the Dirac equation will not be exactly separable on spacetimes with a magnetic moment different from Qa .

Even for $Q = I\pi a$ when $|\rho|\gamma^\mu \tilde{A}_\mu$ reduces to $|\rho|\gamma^0 \tilde{A}_0 = -(Qr/\Delta)\gamma^0$, which is a function of only r , the separation of variables Ansatz does not yield a system of ordinary differential equations which can be solved one at a time, unlike the situation for the familiar Dirac equation for the spectrum of hydrogen in Minkowski spacetime.

2.1.6 A Hilbert space for \hat{H}

In order to decide what is the correct inner product to use for the space of bispinor fields defined on the zGKN spacetime, we pause to consider the action for the original Dirac equation (7), which should be obtainable from this equation upon left-multiplying it by the conjugate bispinor $\bar{\Psi}$, defined as

$$\bar{\Psi} := \Psi^\dagger \gamma^0, \quad (47)$$

and integrating the result on the spacetime. Thus, using oblate spheroidal coordinates,

$$\mathcal{S}[\Psi] = \int dt \int_{\Sigma_t} \Psi^\dagger \gamma^0 [\tilde{\gamma}^\mu \nabla_\mu \Psi + \dots] d\mu_{\Sigma_t}, \quad (48)$$

where $d\mu_{\Sigma_t} = |\rho|^2 \sin\theta d\theta d\varphi dr$ is the volume element of Σ_t , the spacelike $t = \text{constant}$ slice of zGKN. It follows that the natural inner product for bispinors on Σ_t needs to be

$$\langle \Psi, \Phi \rangle = \int_{\Sigma} \Psi^\dagger \gamma^0 \tilde{\gamma}^0 \Phi d\mu_{\Sigma} = \int_0^{2\pi} \int_0^\pi \int_{-\infty}^\infty \Psi^\dagger M \Phi |\rho|^2 \sin\theta d\theta d\varphi dr, \quad (49)$$

with

$$M := \gamma^0 \tilde{\gamma}^0 = \gamma^0 \mathbf{e}_\nu^0 \gamma^\nu = \frac{\Delta}{|\rho|} \alpha^0 + \frac{a \sin\theta}{|\rho|} \alpha^2. \quad (50)$$

Here, α^2 is the second one of the three Dirac alpha matrices in the Weyl (spinor) representation, viz.

$$\alpha^k = \gamma^0 \gamma^k = \begin{pmatrix} \sigma_k & \\ & -\sigma_k \end{pmatrix}, \quad k = 1, 2, 3; \quad (51)$$

for notational convenience, we have also set

$$\alpha^0 = \begin{pmatrix} \mathbf{1}_{2 \times 2} & \\ & \mathbf{1}_{2 \times 2} \end{pmatrix} \quad (52)$$

for the 4×4 identity matrix.

Now, let $\Psi = \mathfrak{D}\hat{\Psi}$ and $\Phi = \mathfrak{D}\hat{\Phi}$, with \mathfrak{D} as in (36). Then we have

$$\langle \Psi, \Phi \rangle = \int_0^{2\pi} \int_0^\pi \int_{-\infty}^\infty \hat{\Psi}^\dagger \hat{M} \hat{\Phi} d\theta d\varphi dr, \quad (53)$$

where

$$\hat{M} := \alpha^0 + \frac{a \sin \theta}{\Delta} \alpha^2. \quad (54)$$

The eigenvalues of \hat{M} are $\lambda_\pm = 1 \pm \frac{a \sin \theta}{\Delta}$, both of which are positive everywhere on this space with Zipoy topology. (Note that $\lambda_- \rightarrow 0$ on the ring, which is not part of the space time but at its boundary.) We may thus take the above as the definition of a positive definite inner product given by the matrix \hat{M} for bispinors defined on the rectangular cylinder¹⁷ $\mathcal{Z} := \mathbb{R} \times [0, \pi] \times [0, 2\pi]$ with its natural measure:

$$\langle \hat{\Psi}, \hat{\Phi} \rangle_{\hat{M}} := \int_{\mathcal{Z}} \hat{\Psi}^\dagger \hat{M} \hat{\Phi} d\theta d\varphi dr. \quad (55)$$

The corresponding Hilbert space¹⁸ is denoted by \mathbf{H} , thus

$$\mathbf{H} := \left\{ \hat{\Psi} : \mathcal{Z} \rightarrow \mathbb{C}^4 \mid \|\hat{\Psi}\|_{\hat{M}}^2 := \langle \hat{\Psi}, \hat{\Psi} \rangle_{\hat{M}} < \infty \right\}. \quad (56)$$

After these preparations we are now ready to state our main results.

2.2 Statement of the Main Theorems

Our results about the symmetry of the spectrum are valid for the Dirac Hamiltonian on a static spacelike slice of the zero- G Kerr spacetime decorated with Sommerfeld fields of arbitrary charge Q and current I . The essential self-adjointness, and location of essential and point spectra, are stated only for the Dirac Hamiltonian on a static spacelike slice of the z GKN spacetime; however, we conjecture that these results are and hold for the more general Hamiltonian.

In the ensuing four sections we will prove the following Theorems about \hat{H} .

2.2.1 Symmetry of the spectrum of the Dirac Hamiltonians

We shall find an operator which commutes with any self-adjoint extension of the formal Dirac operator \hat{H} on \mathbf{H} , with the help of which we prove:

THEOREM 2.5. *Let any self-adjoint extension of the formal Dirac operator \hat{H} on \mathbf{H} be denoted by the same letter. Suppose $E \in \text{spec } \hat{H}$. Then $-E \in \text{spec } \hat{H}$.*

Note that the above result holds for *any* self-adjoint extension of \hat{H} , whatever Q and I are.

2.2.2 Essential self-adjointness of the Dirac Hamiltonian on z GKN

By adapting an argument of Winklmeier–Yamada, we shall prove:

THEOREM 2.6. *For $Q = I\pi a$, i.e., for z GKN, the operator \hat{H} is e.s.a. on \mathbf{H} .*

¹⁷Note that \mathcal{Z} is the $t = \text{const.}$ section of \mathcal{C} .

¹⁸Note that \mathbf{H} is *not equivalent* to standard $L^2(\mathcal{Z})$ whose inner product has the identity matrix in place of \hat{M} .

2.2.3 The continuous spectrum of the Dirac Hamiltonians on zGKN

THEOREM 2.7. *For $Q = I\pi a$ the continuous spectrum of \hat{H} on \mathbb{H} is $\mathbb{R} \setminus (-m, m)$.*

2.2.4 The point spectrum of the Dirac Hamiltonian on zGKN

With the help of the Chandrasekhar–Page–Toop formalism to separate variables, and the Prüfer transform, we will be able to control the point spectrum for the zGKN Dirac Hamiltonian:

THEOREM 2.8. *For $Q = I\pi a$ the point spectrum of \hat{H} on \mathbb{H} is nonempty and located in $(-m, m)$; the end points are not included.*

This completes the formulation of our main results. We next turn to their proofs. The proofs of our main theorems are distributed over four sections corresponding to the various aspects of the spectrum, i.e. symmetry, essential self-adjointness, continuous spectrum, and point spectrum.

3 Proof of Theorem 2.5 (Symmetry of the energy spectrum)

Suppose $E \in \mathbb{R}$ is an eigenvalue of \hat{H} . Then there exists $\hat{\Psi} \in \mathbb{H}$ such that

$$\hat{H}\hat{\Psi} = E\hat{\Psi}. \quad (57)$$

Suppose one can find a bounded linear, or conjugate-linear, operator $\hat{C} : \mathbb{H} \rightarrow \mathbb{H}$ that anti-commutes with \hat{H} , i.e.

$$[\hat{C}, \hat{H}]_+ = \hat{C}\hat{H} + \hat{H}\hat{C} = 0. \quad (58)$$

It is then easy to see that $-E$ must also be an eigenvalue of \hat{H} , since

$$\hat{H}\hat{C}\hat{\Psi} = -\hat{C}\hat{H}\hat{\Psi} = -\hat{C}E\hat{\Psi} = -E\hat{C}\hat{\Psi}. \quad (59)$$

This argument can be extended to show the symmetry of other parts of the spectrum. (See e.g. Glazman [18], p. 205.)

Let $\hat{K} : \mathbb{H} \rightarrow \mathbb{H}$ denote the complex conjugation operator $\hat{K}\hat{\Psi}(x) = \hat{\Psi}^*(x)$, and let $\hat{S} : \mathbb{H} \rightarrow \mathbb{H}$ denote the operator $(\hat{S}\hat{\Psi})(x) = \hat{\Psi}(\varsigma(x))$ where $\varsigma : \mathcal{M} \rightarrow \mathcal{M}$ is the sheet swapping map,

$$\varsigma(r, \theta, \varphi) = (-r, \pi - \theta, \varphi). \quad (60)$$

We claim that the operator $\hat{C} : \mathbb{H} \rightarrow \mathbb{H}$ given by $\hat{C} := \gamma^0 \hat{K} \hat{S}$, viz.

$$(\hat{C}\hat{\Psi})(x) = \gamma^0 \hat{\Psi}^*(\varsigma(x)), \quad (61)$$

anti-commutes with \hat{H} .

To prove the claim, first note that $\gamma^0 = \beta$ anti-commutes with all three α^k matrices. Recall that

$$\hat{H}(x) = \hat{M}^{-1} \mathfrak{H} \quad (62)$$

and

$$\mathfrak{H} := -i\alpha^3 \partial_r + \frac{1}{\Delta} \left(-i\alpha^1 \partial_\theta - i\alpha^2 \csc \theta \partial_\varphi \right) - \frac{ia}{\Delta^2} \alpha^0 \partial_\varphi + \frac{m}{\Delta} \gamma^0 \mathfrak{R} + \frac{e|\rho|}{\Delta} \left(\tilde{A}^0(x) \alpha^0 + \tilde{A}^2(x) \alpha^2 \right). \quad (63)$$

Now

$$\hat{M}^{-1} = \frac{\Delta^2}{|\rho|^2} \left(\alpha^0 - \frac{a \sin \theta}{\Delta} \alpha^2 \right). \quad (64)$$

Thus, keeping in mind that $\overline{\alpha^2} = -\alpha^2$, we find that

$$\hat{C}\hat{M}^{-1} = \gamma^0 \overline{\hat{M}^{-1} \circ \zeta} \hat{K} \hat{S} = \frac{\Delta^2}{|\rho|^2} \gamma^0 \left(\alpha^0 + \frac{a \sin \theta}{\Delta} \alpha^2 \right) \hat{K} \hat{S} = \frac{\Delta^2}{|\rho|^2} \left(\alpha^0 - \frac{a \sin \theta}{\Delta} \alpha^2 \right) \gamma^0 \hat{K} \hat{S} = \hat{M}^{-1} \hat{C}. \quad (65)$$

So we only need to check that \hat{C} anti-commutes with \mathfrak{H} .

It is enough to check that each term in \mathfrak{H} goes through an odd number of sign changes (either one or three) as the three operators \hat{K} , \hat{S} , and multiplication by γ^0 , filter through that term. Recalling that the potential \mathbf{A} is anti-symmetric with respect to sheet swap: $\tilde{A}_\mu \circ \zeta = -\tilde{A}_\mu$, this becomes obvious for most terms in \mathfrak{H} . Only the term involving \mathfrak{R} requires some care. We first check that $\mathfrak{R} \circ \zeta = -\mathfrak{R}$ and that $\gamma^0 \mathfrak{R} = \mathfrak{R} \gamma^0$. Then

$$\hat{C} \gamma^0 \mathfrak{R} = \gamma^0 \gamma^0 \overline{\mathfrak{R} \circ \zeta} \hat{K} \hat{S} = -\gamma^0 \gamma^0 \mathfrak{R} \hat{K} \hat{S} = -\gamma^0 \mathfrak{R} \gamma^0 \hat{K} \hat{S} = -\gamma^0 \mathfrak{R} \hat{C}, \quad (66)$$

establishing the anti-commutation property. The proof of Theorem 2.5 is complete.

For the proof the remaining statements in this paper we rely on the fact that the Dirac equation of a point electron in zGKN separates into four (coupled) ordinary differential equations, each of which depends on only one of the four oblate spheroidal coordinates, with the coupling being effected through shared parameters in the equations. This is carried out in the next section before we resume with proving our claims.

4 Proof of Theorem 2.6 (Essential self-adjointness ($Q = I\pi a$))

We now show that the Dirac Hamiltonian \hat{H} is essentially self-adjoint on \mathbf{H} ; recall that \mathbf{H} is equipped with the inner product (55).

We observe that $M^0 = \Delta \gamma^0 \hat{M}$, so we may rewrite (45) as

$$\hat{H} = \hat{M}^{-1} \gamma^0 \left(\frac{-i}{\Delta} M^k \partial_k + e \frac{|\rho|}{\Delta} \gamma^\mu \tilde{A}_\mu + \frac{m}{\Delta} \mathfrak{R} \right) = \hat{M}^{-1} \mathfrak{H}, \quad (67)$$

where

$$\mathfrak{H} := \mathfrak{M} + \frac{1}{\Delta} \mathfrak{R} + \mathfrak{P} + \mathfrak{Q}, \quad (68)$$

with

$$\mathfrak{M} := -i\alpha^3 \partial_r \quad (69)$$

$$\mathfrak{R} := -i\alpha^1 \partial_\theta - i\alpha^2 \csc \theta \partial_\varphi \quad (70)$$

$$\mathfrak{P} := -i \frac{a}{\Delta^2} \alpha^0 \partial_\varphi + \frac{m}{\Delta} \gamma^0 \mathfrak{R} \quad (71)$$

$$\mathfrak{Q} := e \frac{|\rho|}{\Delta} \gamma^0 \gamma^\mu \tilde{A}_\mu. \quad (72)$$

Thus,

$$\langle \hat{\Psi}, \hat{H} \hat{\Phi} \rangle_{\hat{M}} = \int_{\mathcal{Z}} \hat{\Psi} \mathfrak{H} \hat{\Phi} d\theta d\varphi dr. \quad (73)$$

Evidently, \mathfrak{H} is Hermitian symmetric on the Hilbert space $L^2(\mathcal{Z}; \mathbb{C}^4)$ with its natural inner product

$$(\hat{\Phi}, \hat{\Psi}) = \int_{\mathcal{Z}} \hat{\Phi}^\dagger \hat{\Psi} d\theta d\varphi dr. \quad (74)$$

It is furthermore easy to see that \hat{H} is e.s.a. on \mathbf{H} if and only if \mathfrak{H} is e.s.a. on $L^2(\mathcal{Z}; \mathbb{C}^4)$.

We shall first prove that \mathfrak{H} is e.s.a. on $L^2(\mathcal{Z}; \mathbb{C}^4)$ when $Q = I\pi a$, i.e. for a Dirac point electron in zGKN. Then we will extend this result to small $(Q - I\pi a)e$ using a Hardy-type estimate and perturbation theory.

THEOREM 4.1. *For $Q = I\pi a$ the operator \mathfrak{H} is e.s.a. on $L^2(\mathcal{Z}; \mathbb{C}^4)$.*

Proof. Let us write

$$\mathfrak{H} = \mathfrak{H}^0 + \mathfrak{Q}. \quad (75)$$

Here, \mathfrak{H}^0 is the free Hamiltonian. We will first show that \mathfrak{H}^0 is essentially self-adjoint; this proof is an easy adaptation of the method first employed by Winklmeier and Yamada [40]. We will then conclude essential self-adjointness of \mathfrak{H} by using a perturbation argument.

To this end, let us consider the decomposition with respect to the azimuthal angle φ of the Hilbert space $L^2(\mathcal{Z}; \mathbb{C}^4)$ into partial wave subspaces $L^2([0, \pi] \times \mathbb{R}, d\theta dr)$:

$$L^2(\mathcal{Z}; \mathbb{C}^4) = \bigoplus_{\kappa \in \mathbb{Z} + \frac{1}{2}} \left(L^2_{\kappa}([0, \pi] \times \mathbb{R}, d\theta dr) \right)^4 \quad (76)$$

corresponding to the expansion of a bispinor field $\hat{\Psi} \in L^2(\mathcal{Z}; \mathbb{C}^4)$ given by

$$\hat{\Psi}(r, \theta, \varphi) = \sum_{\kappa \in \mathbb{Z} + \frac{1}{2}} e^{i\kappa\varphi} \hat{\Psi}_{\kappa}(r, \theta). \quad (77)$$

The reason κ needs to be half of an odd integer is that a spinor needs to change sign upon a 2π rotation about any axis, and therefore we need to have

$$\lim_{\varphi \rightarrow 2\pi} \hat{\Psi} = - \lim_{\varphi \rightarrow 0} \hat{\Psi}. \quad (78)$$

Let $\mathfrak{H}_{\kappa}^0 := \mathfrak{H}^0|_{L^2_{\kappa}}$. Then $\mathfrak{H}_{\kappa}^0 = \mathfrak{S}_{\kappa} + \Delta^{-1}\mathfrak{T}_{\kappa} + \mathfrak{B}_{\kappa}$, with

$$\mathfrak{S}_{\kappa} = -i\alpha^3\partial_r, \quad \mathfrak{T}_{\kappa} = -i\alpha^1\partial_{\theta} + \alpha^2\kappa \csc\theta = \begin{pmatrix} \mathfrak{t}_{\kappa} & 0 \\ 0 & -\mathfrak{t}_{\kappa} \end{pmatrix}, \quad \mathfrak{B}_{\kappa} = \frac{a\kappa}{\Delta^2}\alpha^0 + \frac{m}{\Delta}\gamma^0\mathfrak{A}. \quad (79)$$

We note that \mathfrak{B}_{κ} is a symmetric *bounded* multiplication operator on L^2_{κ} ; in fact,

$$\|\mathfrak{B}_{\kappa}\|_{L^{\infty}} \leq |\kappa/a| + m, \quad (80)$$

so that the task of showing e.s.a.-ness of \mathfrak{H}_{κ}^0 reduces to showing e.s.a.-ness of $\mathfrak{H}'_{\kappa} := \mathfrak{S}_{\kappa} + \Delta^{-1}\mathfrak{T}_{\kappa}$.

Now \mathfrak{H}'_{κ} is block-diagonal:

$$\mathfrak{H}'_{\kappa} = \begin{pmatrix} \mathfrak{h}'_{\kappa} & 0 \\ 0 & -\mathfrak{h}'_{\kappa} \end{pmatrix}, \quad \mathfrak{h}'_{\kappa} := -i\sigma_3\partial_r + \Delta^{-1}\mathfrak{t}_{\kappa}, \quad \mathfrak{t}_{\kappa} := -i\sigma_1\partial_{\theta} + \sigma_2\kappa \csc\theta. \quad (81)$$

Thus it is enough to show \mathfrak{h}'_{κ} is e.s.a. We do so by showing that $\ker(\mathfrak{h}'_{\kappa} \pm i) = \{0\}$: Suppose $\hat{\psi}_{\kappa} \in (L^2_{\kappa}([0, \pi] \times \mathbb{R}, d\theta dr))^2$ satisfies

$$\mathfrak{h}'_{\kappa}\hat{\psi}_{\kappa} = \pm i\hat{\psi}_{\kappa}. \quad (82)$$

As observed in [40], it is possible to decompose (82) with respect to the eigenspaces of the operator

$$\mathfrak{a}_{\kappa} := W\mathfrak{t}_{\kappa}W^{-1} = -i\sigma_2\partial_{\theta} + \kappa \csc\theta\sigma_1, \quad (83)$$

where

$$W := \begin{pmatrix} 0 & 1 \\ i & 0 \end{pmatrix}. \quad (84)$$

The operator \mathfrak{a}_{κ} has pure point spectrum and a complete set of eigenfunctions. More precisely, one has the following result [38] (here quoted from [40]):

THEOREM 4.2. (Winklmeier, 2006) For all $\kappa \in \mathbb{Z} + \frac{1}{2}$ the operator \mathfrak{a}_κ with domain $(C_c^\infty((0, \pi)))^2$ is essentially self-adjoint in $(L^2((0, \pi), d\theta))^2$. Its closure (denoted again by \mathfrak{a}_κ) is compactly invertible and its spectrum consists of simple eigenvalues only, given by

$$\lambda_n^\kappa := \text{sgn}(n) \left(|\kappa| - \frac{1}{2} + |n| \right), \quad n \in \mathbb{Z}^* = \mathbb{Z} \setminus \{0\}, \quad (85)$$

with corresponding normalized eigenfunctions $\{g_n^\kappa\}_{n \in \mathbb{Z}^*}$ forming a complete orthonormal set in $(L^2((0, \pi), d\theta))^2$. Moreover,

$$\lambda_{-n}^\kappa = -\lambda_n^\kappa, \quad g_{-n}^\kappa = -\sigma_3 g_n^\kappa. \quad (86)$$

We can therefore write

$$\hat{\psi}_\kappa(r, \theta) = \sum_{n \in \mathbb{Z}^*} \xi_n(r) g_n^\kappa(\theta), \quad (87)$$

with functions $\xi_n \in L^2(\mathbb{R}, dr)$. Hence, performing a similarity transform on (82) with W and projecting on the spans of g_n^κ and g_{-n}^κ we obtain the following system (see [40] for details):

$$\begin{pmatrix} \frac{\lambda_n^\kappa}{\Delta} & i\partial_r \\ i\partial_r & -\frac{\lambda_{-n}^\kappa}{\Delta} \end{pmatrix} \begin{pmatrix} \xi_n \\ \xi_{-n} \end{pmatrix} = \pm i \begin{pmatrix} \xi_n \\ \xi_{-n} \end{pmatrix} \quad (88)$$

However the operator in the above eigenvalue problem $\mathcal{C}_\kappa = i\sigma_1 \partial_r + \frac{\lambda_n^\kappa}{\Delta} \sigma_3$ is clearly e.s.a., since $\frac{\lambda_n^\kappa}{\Delta}$ is bounded, hence $\xi_n = \xi_{-n} = 0$.

This completes the proof of essential self-adjointness of \mathfrak{H}^0 .

Consider now the term $\mathfrak{Q} = e \frac{|\rho|}{\Delta} \gamma^0 \gamma^\mu \tilde{A}_\mu$ coming from the electromagnetic potential. It can be rewritten as $\mathfrak{Q} = -eQ \mathfrak{V}_1 - e(Q - I\pi a) \mathfrak{V}_2$, where

$$\mathfrak{V}_1 := \frac{r}{\Delta^2} \alpha^0, \quad \mathfrak{V}_2 := \frac{ar \sin \theta}{\Delta |\rho|^2} \hat{M}. \quad (89)$$

The first term, \mathfrak{V}_1 , is clearly bounded, whereas the second one, \mathfrak{V}_2 , blows up on the ring. However, since by hypothesis we restrict ourselves to the case $Q = I\pi a$, the \mathfrak{V}_2 term is absent from \mathfrak{Q} , and essential self-adjointness of \mathfrak{H} follows easily from that of \mathfrak{H}^0 and the boundedness of $eQ \mathfrak{V}_1$.

The proof of Theorem 4.1 is complete. \square

5 Chandrasekhar–Page–Toop separation-of-variables ($Q = I\pi a$)

When $Q = I\pi a$ the Dirac equation (46) for the bispinor $\hat{\Psi}$ allows a clear separation also for the remaining r and θ derivatives (commonly referred to in the literature as “radial” and “angular” derivatives, even though r is not a radius and θ is not an angle, except at infinity). Thus, when $Q = I\pi a$ the Dirac equation (46) becomes

$$(\hat{R} + \hat{A}) \hat{\Psi} = 0, \quad (90)$$

where

$$\hat{R} := \begin{pmatrix} imr & 0 & D_- + ieQ \frac{r}{\Delta} & 0 \\ 0 & imr & 0 & D_+ + ieQ \frac{r}{\Delta} \\ D_+ + ieQ \frac{r}{\Delta} & 0 & imr & 0 \\ 0 & D_- + ieQ \frac{r}{\Delta} & 0 & imr \end{pmatrix}, \quad (91)$$

$$\hat{A} := \begin{pmatrix} -ma \cos \theta & 0 & 0 & -L_- \\ 0 & -ma \cos \theta & -L_+ & 0 \\ 0 & L_- & ma \cos \theta & 0 \\ L_+ & 0 & 0 & ma \cos \theta \end{pmatrix}, \quad (92)$$

where D_{\pm} and L_{\pm} have been given in (29). Once a solution $\hat{\Psi}$ to (90) is found, the bispinor $\Psi := \mathfrak{D}\hat{\Psi}$ solves the original Dirac equation (7).

5.0.5 The Chandrasekhar Ansatz

Assume now that a solution $\hat{\Psi}$ of (90) is of the form

$$\hat{\Psi} = e^{iEt - i\kappa\varphi} \begin{pmatrix} R_1 S_1 \\ R_2 S_2 \\ R_2 S_1 \\ R_1 S_2 \end{pmatrix}, \quad (93)$$

with R_k being complex-valued functions of r alone, and S_k real-valued functions of θ alone. Let

$$\vec{R} := \begin{pmatrix} R_1 \\ R_2 \end{pmatrix}, \quad \vec{S} := \begin{pmatrix} S_1 \\ S_2 \end{pmatrix}. \quad (94)$$

Plugging the Chandrasekhar Ansatz (93) into (90) one easily finds that there must be $\lambda \in \mathbb{C}$ such that

$$T_{rad}\vec{R} = E\vec{R}, \quad (95)$$

$$T_{ang}\vec{S} = \lambda\vec{S}, \quad (96)$$

where

$$T_{rad} := \begin{pmatrix} d_- & -m\frac{r}{\Delta} - i\frac{\lambda}{\Delta} \\ -m\frac{r}{\Delta} + i\frac{\lambda}{\Delta} & -d_+ \end{pmatrix} \quad (97)$$

$$T_{ang} := \begin{pmatrix} -ma \cos \theta & -l_- \\ l_+ & ma \cos \theta \end{pmatrix} \quad (98)$$

The operators d_{\pm} and l_{\pm} are now ordinary differential operators in r and θ respectively, with coefficients that depend on the unknown E , and parameters a , κ , and eQ :

$$d_{\pm} := i\frac{d}{dr} \pm \frac{-a\kappa + eQr}{\Delta^2} \quad (99)$$

$$l_{\pm} := \frac{d}{d\theta} \mp (aE \sin \theta - \kappa \csc \theta) \quad (100)$$

The angular operator T_{ang} in (96) is easily seen to be essentially self-adjoint on $(C_c^{\infty}((0, \pi), \sin \theta d\theta))^2$ and in fact is self-adjoint on its domain inside $(L^2((0, \pi), \sin \theta d\theta))^2$ (e.g. [31, 3]) with purely point spectrum $\lambda = \lambda_n(am, aE, \kappa)$, $n \in \mathbb{Z} \setminus 0$. Thus in particular $\lambda \in \mathbb{R}$. It then follows that the radial operator T_{rad} is also essentially self-adjoint on $(C_c^{\infty}(\mathbb{R}, dr))^2$ and in fact self-adjoint on its domain inside $(L^2(\mathbb{R}, dr))^2$.

Suppose $\vec{R} = (R_1, R_2)^T \in (L^2(\mathbb{R}))^2$ is a nontrivial solution to $T_{rad}\vec{R} = E\vec{R}$, with $E \in \mathbb{R}$. Then

$$\begin{aligned} \frac{dR_1}{dr} - i \left(E - \frac{a\kappa - eQr}{\Delta^2} \right) R_1 + \frac{1}{\Delta} (imr - \lambda) R_2 &= 0 \\ -\frac{dR_2}{dr} - i \left(E - \frac{a\kappa - eQr}{\Delta^2} \right) R_2 + \frac{1}{\Delta} (imr + \lambda) R_1 &= 0. \end{aligned}$$

Multiply the first equation by \bar{R}_1 and the second equation by \bar{R}_2 , add them and take the real part, to obtain

$$\frac{d}{dr} (|R_1|^2 - |R_2|^2) = 0. \quad (101)$$

Thus the difference of the moduli squared of R_1 and R_2 is constant, hence zero since they need to be integrable at infinity. I.e.,

$$|R_1| = |R_2| := R. \quad (102)$$

Let $R_j = Re^{i\Phi_j}$ for $j = 1, 2$. Multiply the first equation by \bar{R}_2 , multiply the complex conjugate of the second equation by R_1 , and add them to obtain

$$\frac{d}{dr} \left(\frac{R_1}{\bar{R}_2} \right) = 0. \quad (103)$$

Thus the ratio R_1/\bar{R}_2 , and hence the sum of the arguments $\Phi_1 + \Phi_2$ must be a constant, say δ . Thus $R_1 = \bar{R}_2 e^{i\delta}$. Since multiplication by a constant phase factor is a gauge transformation for Dirac bispinors, we can replace $\hat{\Psi}$ with $\hat{\Psi}' = e^{-i\delta/2} \hat{\Psi}$ without changing anything. The spinor thus obtained has the same form as (93), now with $R'_1 = \bar{R}'_2$. Thus without loss of generality we can assume $\delta = 0$ and $R_1 = \bar{R}_2$.

This motivates us to set

$$R_1 = \frac{1}{\sqrt{2}}(v - iu), \quad R_2 = \frac{1}{\sqrt{2}}(v + iu) \quad (104)$$

for real functions u and v . Consider the unitary matrix

$$U := \frac{1}{\sqrt{2}} \begin{pmatrix} -i & 1 \\ i & 1 \end{pmatrix}. \quad (105)$$

A change of basis using U brings the radial system (95) into the following standard (Hamiltonian) form

$$(H_{rad} - E) \begin{pmatrix} u \\ v \end{pmatrix} = \begin{pmatrix} 0 \\ 0 \end{pmatrix}, \quad (106)$$

where

$$H_{rad} := \begin{pmatrix} m \frac{r}{\Delta} + \frac{\gamma r + a\kappa}{\Delta^2} & -\partial_r + \frac{\lambda}{\Delta} \\ \partial_r + \frac{\lambda}{\Delta} & -m \frac{r}{\Delta} + \frac{\gamma r + a\kappa}{\Delta^2} \end{pmatrix}, \quad (107)$$

(cf. [35], eq (7.105)) with

$$\gamma := -eQ < 0. \quad (108)$$

6 Proof of Theorem 2.7 (Continuous spectrum of \hat{H} on zGKN)

Following Weidmann [37] we now prove the theorem about the continuous spectrum of \hat{H} . Recall the partial wave decomposition (76). Let \hat{H}_κ denote the restriction of \hat{H} to L^2_κ . From the Chandrasekhar separation (93) and equation (95) it follows that the spectrum of \hat{H}_κ coincides with that of T_{rad} , which coincides with that of H_{rad} since these last two are unitarily equivalent. Furthermore, the spectrum of \hat{H} equals the union of the spectra of \hat{H}_κ . Thus in order to prove the claim about the essential spectrum, it suffices to show that it holds for H_{rad} regardless of the values of κ and λ .

Since H_{rad} is a radial Dirac operator, one can then use results that are particular to one dimension. One such result is due to Weidmann [37]:

THEOREM. Let P and J be matrices such that $H_{rad} = J\partial_r + P$. Suppose P can be written as $P_1 + P_2$ in such a way that each component of P_1 is integrable in $[R, \infty)$ for some $R > 0$, P_2 is of bounded variation on $[R, \infty)$ and

$$\lim_{r \rightarrow \infty} P(r) = \begin{pmatrix} a & 0 \\ 0 & b \end{pmatrix}, \quad a > b.$$

Then every self-adjoint extension of h has a purely absolutely continuous spectrum in $(-\infty, a] \cup [b, \infty)$.

Using this result, our claim is established by noticing that the hypotheses on P are satisfied, and

$$\lim_{r \rightarrow \infty} P(r) = \begin{pmatrix} m & 0 \\ 0 & -m \end{pmatrix}$$

Proof of Theorem 2.7 is complete.

7 Proof of Theorem 2.8 (Point spectrum of \hat{H} on zGKN)

By the remarks at the beginning of Section 6, we are interested in the eigenvalues E and square-integrable eigenfunctions in $L^2(\mathbb{R}, dr)^2$ of the operator H_{rad} . One complication is that in our case the radial Hamiltonian H_{rad} depends on the unknown eigenvalues λ of the angular operator T_{ang} in (96), which in turn depend on the energy E . Since the angular operator is the same as the one on Kerr and Kerr-Newman spacetime studied in [31, 3], and since it is known that for a given value of E there is a largest negative eigenvalue $\lambda = \Lambda(E)$, our strategy is to show the existence, for a given value of $\lambda < 0$, of a smallest positive eigenvalue $E = \mathcal{E}(\lambda)$ for H_{rad} , and then set up an iteration that converges to a pair (E, λ) for which both the radial (95) and the angular (96) have L^2 solutions, thereby establishing the existence of a “positive ground state” for the full Dirac Hamiltonian; note that by the symmetry of the spectrum there also exists a “negative ground state.”

7.1 The Prüfer transform

Consider the equations (106) and (96) for unknowns (u, v) and (S_1, S_2) . Let us define new unknowns (R, Ω) and (S, Θ) via the Prüfer transform [28]

$$u = \sqrt{2}R \cos \frac{\Omega}{2}, \quad v = \sqrt{2}R \sin \frac{\Omega}{2}, \quad S_1 = S \cos \frac{\Theta}{2}, \quad S_2 = S \sin \frac{\Theta}{2}. \quad (109)$$

Thus

$$R = \frac{1}{2} \sqrt{u^2 + v^2}, \quad \Omega = 2 \tan^{-1} \frac{v}{u}, \quad S = \sqrt{S_1^2 + S_2^2}, \quad \Theta = 2 \tan^{-1} \frac{S_2}{S_1}. \quad (110)$$

As a result, $R_1 = -iRe^{i\Omega/2}$ and $R_2 = iRe^{-i\Omega/2}$. Hence $\hat{\Psi}$ can be re-expressed in terms of the Prüfer variables, thus

$$\hat{\Psi}(t, r, \theta, \varphi) = R(r)S(\theta)e^{i(Et - \kappa\varphi)} \begin{pmatrix} -i \cos(\Theta(\theta)/2)e^{+i\Omega(r)/2} \\ +i \sin(\Theta(\theta)/2)e^{-i\Omega(r)/2} \\ +i \cos(\Theta(\theta)/2)e^{-i\Omega(r)/2} \\ -i \sin(\Theta(\theta)/2)e^{+i\Omega(r)/2} \end{pmatrix}, \quad (111)$$

and we obtain the following equations for the new unknowns

$$\frac{d}{dr}\Omega = 2\frac{mr}{\Delta}\cos\Omega + 2\frac{\lambda}{\Delta}\sin\Omega + 2\frac{a\kappa + \gamma r}{\Delta^2} - 2E, \quad (112)$$

$$\frac{d}{dr}\ln R = \frac{mr}{\Delta}\sin\Omega - \frac{\lambda}{\Delta}\cos\Omega. \quad (113)$$

Similarly,

$$\frac{d}{d\theta}\Theta = -2ma\cos\theta\cos\Theta + 2\left(aE\sin\theta - \frac{\kappa}{\sin\theta}\right)\sin\Theta + 2\lambda, \quad (114)$$

$$\frac{d}{d\theta}\ln S = -ma\cos\theta\sin\Theta - \left(aE\sin\theta - \frac{\kappa}{\sin\theta}\right)\cos\Theta. \quad (115)$$

To simplify the analysis of these systems and reduce the number of parameters involved, we will henceforth set $m = 1$. Note that this is always possible by defining the constants $a' = ma$, $E' = E/m$, and a change of variable $r' = mr$.

7.2 The realm of L^2 solutions

We note that in both of the above systems, when a solution to the first equation is known, the second equation in the system can be solved by quadrature. Moreover, the requirement that R and S be L^2 functions of their argument determines what boundary values the solutions to the Ω and Θ equations should have. More precisely,

PROPOSITION 7.1. *Any bispinor $\hat{\Psi}$ of the form (111) constructed from solutions of (112), (113), (114), (115), with $\kappa \geq \frac{1}{2}$, $E > 0$ and $\lambda < 0$, belongs to the Hilbert space \mathbf{H} provided*

$$\lim_{r \rightarrow -\infty} \Omega(r) = -\pi + \cos^{-1}(E), \quad \lim_{r \rightarrow \infty} \Omega(r) = -\cos^{-1}(E), \quad (116)$$

and

$$\Theta(0) = 0, \quad \Theta(\pi) = -\pi. \quad (117)$$

Proof. It is straightforward to compute that for a $\hat{\Psi}$ of the form (111),

$$\begin{aligned} \|\hat{\Psi}\|_M^2 &= 2 \int_0^{2\pi} \int_0^\pi \int_{-\infty}^\infty R^2(r) S^2(\theta) \left(1 + \frac{a \sin \theta}{\Delta} \sin \Theta(\theta) \sin \Omega(r)\right) dr d\theta d\varphi \\ &= 4\pi \left[\int_{-\infty}^\infty R^2 dr \int_0^\pi S^2 d\theta + a \int_{-\infty}^\infty R^2 \sin \Omega \frac{dr}{\Delta} \int_0^\pi S^2 \sin \Theta \sin \theta d\theta \right] \\ &\leq 8\pi \|R\|_{L^2}^2 \|S\|_{L^2}^2, \end{aligned}$$

and thus $\hat{\Psi} \in \mathbf{H}$ provided $R \in L^2(\mathbb{R}, dr)$ and $S \in L^2((0, \pi), d\theta)$.

Now (114) can be written as a smooth dynamical system in the (θ, Θ) plane by introducing a new independent variable τ such that $\frac{d\theta}{d\tau} = \sin \theta$. Then, with dot representing differentiation in τ , we have,

$$\begin{cases} \dot{\theta} &= \sin \theta \\ \dot{\Theta} &= -2a \sin \theta \cos \theta \cos \Theta + 2aE \sin^2 \theta \sin \Theta - 2\kappa \sin \Theta + 2\lambda \sin \theta \end{cases} \quad (118)$$

Identifying the line $\Theta = \pi$ with $\Theta = -\pi$, this becomes a dynamical system on a closed finite cylinder $\mathcal{C}_1 = [0, \pi] \times \mathbb{S}^1$. The only equilibrium points of the flow are on the two circular boundaries: Two on the left boundary: $S_- = (0, 0)$, $N_- = (0, \pi)$; two on the right: $S_+ = (\pi, -\pi)$ and $N_+ = (\pi, 0)$.

The linearization of the flow at the equilibrium points reveals that S_- and S_+ are hyperbolic saddle points (with eigenvalues $\{1, -2\kappa\}$ and $\{-1, 2\kappa\}$ respectively), while N_- is a source node (with eigenvalues 1 and 2κ) and N_+ is a sink node (with eigenvalues -1 and -2κ). The α -limit set of the orbit of any point in the interior of the cylinder must necessarily be either S_- or N_- , and likewise its ω -limit set can only be either N_+ or S_+ . The only possible boundary values for $\Theta(\theta)$ are therefore 0 and $\pm\pi$ at each endpoint of the interval $0 \leq \theta \leq \pi$. The boundary values (117) correspond to a heteroclinic orbit connecting the two saddles S_- and S_+ . Suppose such a saddle connection exists. Since the eigendirection corresponding to the unstable manifold of S_- is $v_1 = (\kappa + \frac{1}{2}, \lambda - a)^T$ and the stable manifold of S_+ has the same eigendirection, it follows that for the said saddle connection, we have

$$\frac{d\Theta}{d\theta}\Big|_{\theta=0} = \frac{d\Theta}{d\theta}\Big|_{\theta=\pi} = \frac{\lambda - a}{\kappa + \frac{1}{2}} =: \delta < 0. \quad (119)$$

Thus $\Theta = \delta\theta + o(1)$ as $\theta \rightarrow 0$ and $\Theta = -\pi + \delta(\theta - \pi) + o(1)$ as $\theta \rightarrow \pi$. Consider now the S equation (115). By the above,

$$\frac{d}{d\theta} \ln S = \begin{cases} \frac{\kappa}{\theta} + o(1) & \text{as } \theta \rightarrow 0 \\ \frac{\kappa}{\theta - \pi} + o(1) & \text{as } \theta \rightarrow \pi \end{cases} \quad (120)$$

Integrating in θ we thereby conclude that $S \sim |\theta|^\kappa$ for θ small and $S \sim |\pi - \theta|^\kappa$ for θ near π . Therefore S is integrable on $(0, \pi)$ and indeed it belongs to $L^p((0, \pi), d\theta)$ for any $p \geq 1$.

Consider next the Ω equation (112). It can also be rewritten as a smooth dynamical system on a cylinder, in this case by setting $\tau := \frac{r}{a}$ as new independent variable, as well as introducing a new dependent variable

$$\xi := \tan^{-1} \frac{r}{a} = \tan^{-1} \tau \quad (121)$$

Then, with dot again representing differentiation in τ , (112) is equivalent to

$$\begin{cases} \dot{\xi} &= \cos^2 \xi \\ \dot{\Omega} &= 2a \sin \xi \cos \Omega + 2\lambda \cos \xi \sin \Omega + 2\gamma \sin \xi \cos \xi + 2\kappa \cos^2 \xi - 2aE \end{cases} \quad (122)$$

Once again, identifying $\Omega = -\pi$ with $\Omega = \pi$ turns this into a smooth flow on the closed finite cylinder $\mathcal{C}_2 := [-\frac{\pi}{2}, \frac{\pi}{2}] \times \mathbb{S}^1$. The only equilibrium points of the flow are on the two circular boundaries. For $E \in [0, 1)$ there are two equilibria on each: $S_- = (-\frac{\pi}{2}, -\pi + \cos^{-1} E)$ and $N_- = (-\frac{\pi}{2}, \pi - \cos^{-1} E)$ on the left boundary, and $S_+ = (\frac{\pi}{2}, -\cos^{-1} E)$ and $N_+ = (\frac{\pi}{2}, \cos^{-1} E)$ on the right boundary. S_\mp are non-hyperbolic (degenerate) saddle-nodes, with eigenvalues 0 and $\pm 2a\sqrt{1 - E^2}$, while N_- is a degenerate source-node and N_+ a degenerate sink-node¹⁹. The boundary values (116) correspond to a heteroclinic orbit connecting S_- and S_+ .

Suppose such a saddle connection exists, and consider the R equation (113). As $r \rightarrow \pm\infty$, we will then have

$$\frac{d}{dr} \ln R \sim -\text{sgn}(r)\sqrt{1 - E^2} \quad (123)$$

so that integrating in r we will obtain

$$R(r) \sim e^{-|r|\sqrt{1 - E^2}} \quad \text{as } r \rightarrow \pm\infty \quad (124)$$

which ensures that R is integrable at infinity. Since the right-hand-side of the R equation is smooth in Ω and r , and Ω itself is smooth, it follows that $R \in L^p(\mathbb{R}, dr)$ for all $p \geq 1$. \square

¹⁹For $E = 1$ each S, N pair coalesces into one degenerate equilibrium: $N_-^1 = (-\frac{\pi}{2}, \pm\pi)$ and $N_+^1 = (\frac{\pi}{2}, 0)$ with both eigenvalues being zero.

7.3 Existence of heteroclinic orbits connecting the two saddles

From the proof of Proposition 7.1 it is evident that in order to establish the existence of an eigenfunction for the Dirac Hamiltonian of a point electron in the zGKN spacetime, we need to show that there exists a pair (E, λ) such that both dynamical systems (118) and (122) have a *saddle-saddle connecting orbit* for those values of E and λ . We call this type of orbit a *saddle connector* for the corresponding flow. We pave the road for our proof by recalling some general facts of flow on a cylinder.

7.3.1 Flow on a finite cylinder

Let $\mathcal{C} := [x_-, x_+] \times \mathbb{S}^1$ be a finite cylinder. We denote its universal cover by $\bar{\mathcal{C}} := [x_-, x_+] \times \mathbb{R}$, with coordinates (x, y) , and fix a fundamental domain $\tilde{\mathcal{C}} := [x_-, x_+] \times [-\pi, \pi)$ in $\bar{\mathcal{C}}$. Consider the flow Φ_t on $\bar{\mathcal{C}}$ given by the dynamical system

$$\begin{cases} \dot{x} &= f(x) \\ \dot{y} &= g(x, y) \end{cases} \quad (125)$$

where the dot represents differentiation with respect to a formal “time” parameter τ , the functions f and g are smooth, and g is 2π -periodic in y : $g(x, y) = g(x, y + 2\pi)$. Let us moreover assume that f satisfies

$$f(x_-) = f(x_+) = 0, \quad f(x) > 0 \quad \forall x \in (x_-, x_+) \quad (126)$$

while g satisfies

$$g(x_-, y) = 0 \implies y \in \{n_-, s_-\}, \quad g(x_+, y) = 0 \implies y \in \{n_+, s_+\} \quad (127)$$

where $-\pi \leq s_- < n_- \leq \pi$ and $-\pi \leq s_+ < n_+ \leq \pi$. These assumptions imply that the following four distinct points in \mathcal{C} are equilibrium points for the flow:

$$N_{\pm} := (x_{\pm}, n_{\pm}), \quad S_{\pm} := (x_{\pm}, s_{\pm}). \quad (128)$$

We shall further assume that the flow does not have any non-wandering points other than the above four equilibria.

The following assumptions fix the character of the four equilibrium points:

$$f'(x_-) \geq 0, \quad f'(x_+) \leq 0, \quad f''(x_{\pm}) \neq 0 \quad (129)$$

(where by $f'(x_{\pm})$ we mean the left derivate at x_+ and the right derivative at x_-), and

$$D_y g(x_-, n_-) > 0, \quad D_y g(x_-, s_-) < 0, \quad D_y g(x_+, n_+) < 0, \quad D_y g(x_+, s_+) > 0, \quad (130)$$

where $D_y g$ is the y -derivative of $g(x, y)$. Thus N_- is a (source) node, N_+ a (sink) node, and S_{\pm} are saddle points. These will be hyperbolic if $f'(x_{\pm}) \neq 0$, and non-hyperbolic (degenerate) otherwise.

Later on, in order to have a well-defined notion of index for certain distinguished orbits on \mathcal{C} , we will also assume that the locations of the equilibria on the boundary of the cylinder are not arbitrary, but are subject to the single condition

$$s_- - n_- = n_+ - s_+ \pmod{2\pi} \quad (131)$$

(This is a condition on $g(x, y)$. Although we will not pursue this approach here, under this condition (125) can be viewed as a flow with two equilibrium points on a *2-torus*.)

For a point $p \in \mathcal{C}$, let $\mathcal{O}(p)$ denote the flow orbit through p . Since \mathcal{C} is compact, all orbits are complete, meaning they exist for all $s \in \mathbb{R}$, and since the flow is autonomous, two orbits are either disjoint or they coincide. The orbit of an equilibrium point consists of only one point, namely the equilibrium itself. The ω -limit of any other orbit in \mathcal{C} can be either N_+ or S_+ , and the α -limit likewise can only be either N_- or S_- . All these facts are easy consequences of the existence and uniqueness theorem for ODEs.

7.3.2 Connecting orbits and corridors

Given a flow on \mathcal{C} as in the above, there are two distinguished orbits in the interior of the cylinder: Let \mathcal{W}^- denote the unique orbit of the flow whose α -limit point is the saddle S_- , and let \mathcal{W}^+ denote the unique orbit whose ω -limit point is the saddle S_+ . In the hyperbolic case ($f'(x_{\pm}) \neq 0$) the uniqueness is immediate because \mathcal{W}^- is the unstable manifold of S_- and \mathcal{W}^+ is the stable manifold of S_+ . In the non-hyperbolic case $f'(x_{\pm}) = 0$ the orbits \mathcal{W}^{\pm} are *center* manifolds for the corresponding saddle-nodes S_{\pm} . Recall that center manifolds may be non-unique, but in our case the uniqueness is assured because the equilibrium points are on the boundary of the domain, so the relevant part of the center manifolds are on the “saddle side” of the equilibrium, and not on the “node side” (see Figure 2). If \mathcal{W}^+ and \mathcal{W}^- intersect, they must coincide, and the resulting

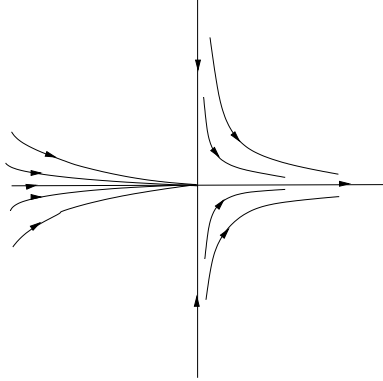


Figure 2: Flow near a saddle-node. The node part lies outside of the domain of concern.

orbit will connect the two saddle points, i.e. it will be the saddle connector we are after. Let us therefore assume that they are disjoint. The ω -limit of \mathcal{W}^- must then necessarily be N_+ , and the α -limit of \mathcal{W}^+ must be N_- .

On the other hand the assumptions we have made about the flow imply that there are also two orbits of the flow on the left boundary of the cylinder connecting N_- with S_- , call them $(N_-S_-)_{\pm}$, with $+$ denoting the counterclockwise one (when viewed from a point on the cylinder’s axis and to the left of the cylinder), and similarly two joining S_+ with N_+ , called $(S_+N_+)_{\pm}$. Consider therefore the following collection of six heteroclinic orbits

$$\mathcal{H} := \{(N_-S_-)_{\pm}, \mathcal{W}^{\pm}, (S_+N_+)_{\pm}\}. \quad (132)$$

The cylinder \mathcal{C} is divided into two invariant regions \mathcal{K}_1 and \mathcal{K}_2 , called *corridors*, by these orbits: $\mathcal{C} = \mathcal{K}_1 \cup \mathcal{H} \cup \mathcal{K}_2$. We would like to distinguish one of these two corridors. We do so as follows: Consider the lifting of the flow to the universal cover $\tilde{\mathcal{C}}$. Let \tilde{S}_- denote the unique copy of the node S_- that lies in the fundamental domain $\tilde{\mathcal{C}}$, and let $\tilde{\mathcal{W}}^-$ be the unique orbit in $\tilde{\mathcal{C}}$ whose α -limit point is \tilde{S}_- . The ω -limit point of this orbit is thus some copy of the node N_+ , call it \tilde{N}_+ , which

has coordinates $(x_+, n_+ - 2\pi k_+)$ for some $k_+ \in \mathbb{Z}$. Similarly, let \tilde{S}_+ denote the unique point in the preimage of S_+ under the covering map that lies in the fundamental domain $\tilde{\mathcal{C}}$ and let $\tilde{\mathcal{W}}^+$ denote the unique orbit whose ω -limit point is this \tilde{S}_+ . Let $\tilde{N}_- = (x_-, n_- + 2\pi k_-)$, $k_- \in \mathbb{Z}$ be the α -limit point of $\tilde{\mathcal{W}}^+$. By definition the corridor \mathcal{K}_1 is the open domain in $\tilde{\mathcal{C}}$ whose boundary contains the two orbits $\tilde{\mathcal{W}}^-$ and $\tilde{\mathcal{W}}^+$.

We note that in $\tilde{\mathcal{C}}$ only one of the two corridors will have *both* of these orbits on its boundary, so this is the distinguishing feature of \mathcal{K}_1 .

We orient the boundary of \mathcal{K}_1 (which is a closed simple curve) in such a way that the orientation induced on $\tilde{\mathcal{W}}^-$ coincides with the direction of the flow on that orbit.

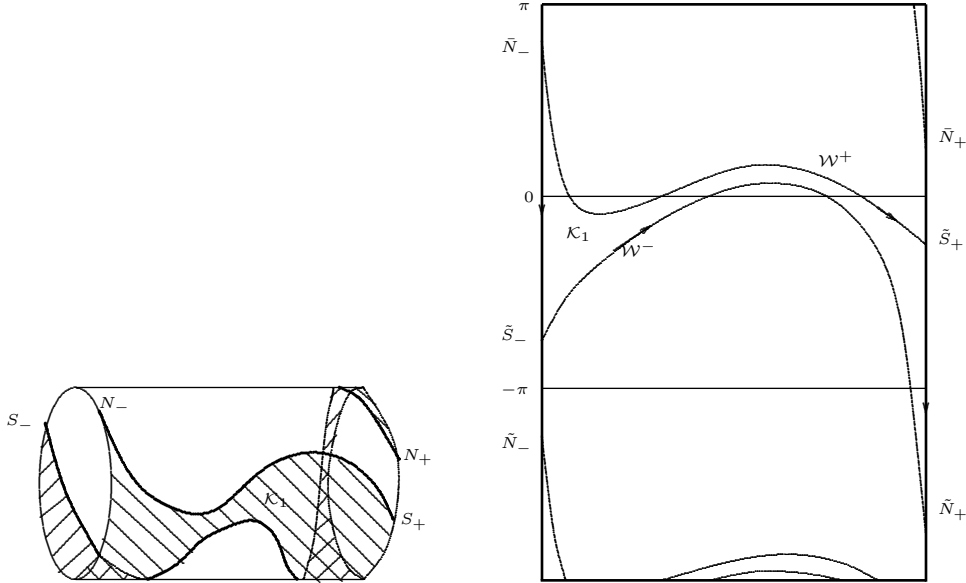


Figure 3: The corridor \mathcal{K}_1 in \mathcal{C} and in $\tilde{\mathcal{C}}$.

Furthermore, in the universal cover $\tilde{\mathcal{C}}$, using the well-orderedness of \mathbb{R} it is possible to speak of $\tilde{\mathcal{W}}^+$ as being situated “above” or “below” $\tilde{\mathcal{W}}^-$. It is evident that the boundary of the corridor \mathcal{K}_1 is oriented clockwise if $\tilde{\mathcal{W}}^+$ is below $\tilde{\mathcal{W}}^-$, and counterclockwise if it is the other way around. see Figure 3.

7.3.3 Parameter-dependent flows

Suppose that the function g in (125) depends smoothly on a parameter $\mu \in I$ where I is an open interval in \mathbb{R} (It’s enough for the μ -dependence to be C^1). Thus the flow is now

$$\begin{cases} \dot{x} = f(x) \\ \dot{y} = g_\mu(x, y) \end{cases} \quad (133)$$

By the implicit function theorem, the locations of the equilibria S_\pm, N_\pm also depend –in a C^1 fashion– on μ , so long as the non-degeneracy conditions (130) are satisfied.

We need to make certain assumptions about the μ -dependence of the flow regarding its monotonicity, and the topology of its nullclines:

Monotonicity We will only consider parameter-dependent flows for which the function $g_\mu(x, y)$ is *monotone non-increasing* in μ .

Assumption (M): For all $(x, y) \in \bar{\mathcal{C}}$, we have

$$\frac{\partial}{\partial \mu} g_\mu(x, y) \leq 0. \quad (134)$$

This assumption in particular implies a corresponding monotonicity for the orbits of the flow (133):

LEMMA 7.2. *Let $\mathcal{O}_\mu = \{(x(\tau), y_\mu(\tau))\}_{\tau \in \mathbb{R}}$ be a complete orbit of (133). Then*

$$\mu_1 < \mu_2 \implies y_{\mu_1}(\tau) \geq y_{\mu_2}(\tau) \quad \text{for all } \tau \in \mathbb{R}. \quad (135)$$

Proof. Easy! (will fill in later) □

Topology of nullclines Consider the subset of the cylinder \mathcal{C} defined by

$$\Gamma := \{(x, y) \in \mathcal{C} \mid g_\mu(x, y) = 0\}. \quad (136)$$

Thus Γ is the zero level-set of g_μ . Since g_μ is smooth, Γ is a curve (or collection of curves) in \mathcal{C} . These curves are referred to as the *y-nullclines* of the flow (125). They have the property that any orbit of the flow that crosses them must have a horizontal tangent at the crossing point. Moreover, Γ divides \mathcal{C} into two regions, thus $\mathcal{C} = \Gamma \cup \mathcal{N} \cup \mathcal{P}$, where

$$\mathcal{N} := \{(x, y) \in \mathcal{C} \mid g_\mu(x, y) < 0\}, \quad \mathcal{P} := \{(x, y) \in \mathcal{C} \mid g_\mu(x, y) > 0\}. \quad (137)$$

Thus the y coordinate of (the lift to the universal cover of) any orbit must decrease in \mathcal{N} and must increase in \mathcal{P} .

Evidently, Γ must include all the singular points of the flow: $S_\pm \in \Gamma$, $N_\pm \in \Gamma$.

It may happen that as μ crosses a critical value $\mu_c \in I$, the topology of the nullclines undergoes a dramatic change. This is indeed the case for the flows that we are studying in this paper. We introduce an assumption that amounts to having some control on this change in nullcline topology:

Assumption (A): There exists a $\mu_c \in I$ such that Γ , the zero level-set of g_{μ_c} has a *saddle point* at some interior point of \mathcal{C} . In particular, for $\mu < \mu_c$, the sets \mathcal{N} and \mathcal{P} are both connected, and Γ is the union of two disjoint curves $\Gamma = \Gamma_u \cup \Gamma_d$, with $N_-, S_+ \in \bar{\Gamma}_d$ and $S_-, N_+ \in \bar{\Gamma}_u$. On the other hand for $\mu > \mu_c$, \mathcal{N} is connected, while \mathcal{P} has two connected components \mathcal{P}_l and \mathcal{P}_r , each being a convex subset of \mathcal{C} . Moreover, Γ is the union of two disjoint curves $\Gamma = \Gamma_l \cup \Gamma_r$, with $N_-, S_- \in \bar{\Gamma}_l$ and $S_+, N_+ \in \bar{\Gamma}_r$; see Figure 4.

7.3.4 Winding number of orbits and corridors

Assuming (131), let

$$w_0 := \frac{s_+ - n_-}{2\pi} = \frac{n_+ - s_-}{2\pi}. \quad (138)$$

Given any orbit $\mathcal{O} = \{(x_o(\tau), y_o(\tau)) \mid \tau \in \mathbb{R}\}$ of the flow (125), the following quantity is well-defined:

$$w(\mathcal{O}) = w_0 - \frac{1}{2\pi} (y_o(\infty) - y_o(-\infty)) = w_0 - \frac{1}{2\pi} \int_{-\infty}^{\infty} g_\mu(x_o(\tau), y_o(\tau)) d\tau \quad (139)$$

In particular for the two distinguished orbits $\widetilde{\mathcal{W}}^\pm$ this is easily calculated to be

$$w(\widetilde{\mathcal{W}}^-) = w_0 - \frac{1}{2\pi} (n_+ - 2\pi k_+ - s_-) = k_+ \quad w(\widetilde{\mathcal{W}}^+) = w_0 - \frac{1}{2\pi} (s_+ - (n_- + 2\pi k_-)) = k_- \quad (140)$$

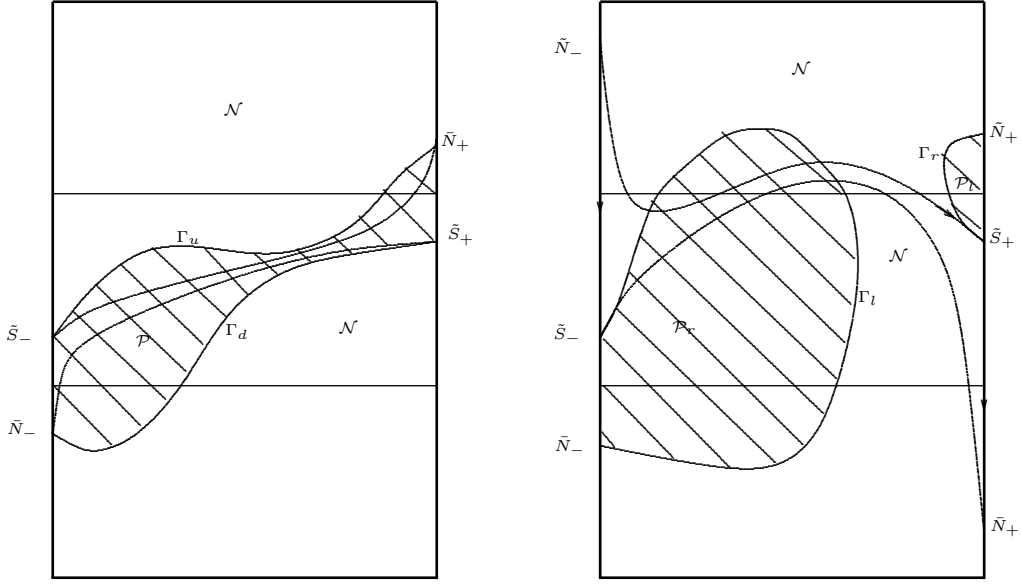


Figure 4: Change in the topology of nullclines: $\mu < \mu_c$ (left) and $\mu > \mu_c$ (right)

and is therefore an integer. We call these the *winding numbers*, of \mathcal{W}^- and \mathcal{W}^+ , respectively.

On the other hand it is easy to see that these two must in fact be equal. The reason is that, since $\widetilde{\mathcal{W}}^-$ goes from $\widetilde{S}^- = (x_-, s_-)$ to $\widetilde{N}^+ = (x_+, n_+ - 2\pi k_+)$, there are two copies of it in the universal cover that go from $(x_-, s_- + 2\pi k_+)$ to (x_+, n_+) and from $(x_-, s_- + 2\pi(k_+ - 1))$ to $(x_+, n_+ - 2\pi)$. Since $n_+ - 2\pi < s_+ < n_+$, these two copies of $\widetilde{\mathcal{W}}^-$ must sandwich any orbit that goes into \widetilde{S}_+ , in particular $\widetilde{\mathcal{W}}^+$. Thus the α -limit point of $\widetilde{\mathcal{W}}^+$ has to be $(x_-, n_- + 2\pi k_+)$, and therefore $k_+ = k_-$.

Definition 7.3. The *winding number of the corridor* \mathcal{K}_1 is the common value of the winding numbers of $\widetilde{\mathcal{W}}^\pm$.

Figure 4 shows a corridor of winding number zero when $\mu < \mu_c$ (left) and one of winding number equal to one for $\mu > \mu_c$ (right).

7.3.5 Continuity argument for existence of saddle connectors

As the parameter μ varies, the two distinguished orbits $\widetilde{\mathcal{W}}^\pm$, and hence the corridor \mathcal{K}_1 that they form will also vary. Let $\mathcal{K}_1(\mu)$ denote the corresponding corridor for parameter value μ (if it exists). Since the winding number $w(\mathcal{K}_1(\mu))$ is integer-valued, if \mathcal{K}_1 varies continuously with respect to μ , its winding number would have to remain constant. It is however possible that for some value of μ the two orbits $\widetilde{\mathcal{W}}^\pm$ coincide and the corridor $\mathcal{K}_1(\mu)$ disappears, leaving a saddle connector behind (for which the winding number will not be an integer). We would like to show that this is the only way for the winding number of $\mathcal{K}_1(\mu)$ to be different for two different values of the parameter μ . In other words:

PROPOSITION 7.4. *Suppose there exist two values $\mu_0 < \mu_1$ in the interval I for each of which a non-empty corridor \mathcal{K}_1 of finite winding number exists, and such that*

$$w(\mathcal{K}_1(\mu_0)) \leq 0 \quad \text{and} \quad w(\mathcal{K}_1(\mu_1)) \geq 1. \quad (141)$$

Then there exists $\mu_s \in (\mu_0, \mu_1)$ such that the flow (133) with $\mu = \mu_s$ has a saddle connector $\mathcal{S}(\mu_s)$, whose lift to the universal cover $\tilde{\mathcal{C}}$ connects the saddle-node \tilde{S}_- to the saddle-node \tilde{S}_+ .

Proof. Let $\mathbf{a}(\mu)$ denote the signed area of $\mathcal{K}_1(\mu)$, defined via Green's theorem:

$$\mathbf{a}(\mu) = \oint_{\partial\mathcal{K}_1(\mu)} (-y)dx = \int_{x_-}^{x_+} (y_\mu^+ - y_\mu^-) dx. \quad (142)$$

Here y_μ^\pm are the y -components of the orbits $\tilde{\mathcal{W}}^\pm$, thought of as functions of x (which is always possible since $x(\tau)$ is a monotone increasing function of flow parameter τ .) The following facts about $\mathbf{a}(\mu)$ are easily verified:

1. $\mathbf{a}(\mu) > 0$ if and only if $w(\mathcal{K}_1(\mu)) \geq 1$.

Proof: Since $\tilde{\mathcal{W}}^+$ and $\tilde{\mathcal{W}}^-$ cannot intersect without coinciding, it is clear from (142) that $\mathbf{a}(\mu) > 0$ if and only if $\tilde{\mathcal{W}}^+$ is above $\tilde{\mathcal{W}}^-$, which is equivalent to $\tilde{S}_+ = (x_+, s_+)$ being above $\tilde{N}_+ = (x_+, n_+ - 2\pi k)$ where $k = w(\mathcal{K}_1(\mu))$. Since $n_+ > s_+$ we must have $k \geq 1$.

2. $\mathbf{a}(\mu) < 0$ if and only if $w(\mathcal{K}_1(\mu)) \leq 0$.

Proof: Similar to above.

3. $\mathbf{a}(\mu) = 0$ if and only if $\mathcal{K}_1(\mu) = \emptyset$.

Proof: From the definition (142) it is clear that the only way for $\mathbf{a}(\mu)$ to be zero is for the two orbits $\tilde{\mathcal{W}}^+$ and $\tilde{\mathcal{W}}^-$ to coincide.

4. $\mathbf{a}(\mu)$ is a continuous function of μ for all $\mu \in [\mu_0, \mu_1]$.

Proof: Let $\epsilon > 0$ be given. For $\mu, \mu' \in [\mu_0, \mu_1]$ and $\xi > 0$ small enough, we have

$$\begin{aligned} |\mathbf{a}(\mu') - \mathbf{a}(\mu)| &= \left| \int_{x_-}^{x_+} (y_{\mu'}^+ - y_\mu^+) - (y_{\mu'}^- - y_\mu^-) dx \right| \\ &\leq \int_{x_-}^{x_- + \xi} |y_{\mu'}^+| + |y_\mu^+| dx + \int_{x_- + \xi}^{x_+ - \xi} |y_{\mu'}^+ - y_\mu^+| dx + \int_{x_+ - \xi}^{x_+} |y_{\mu'}^+| + |y_\mu^+| dx \\ &\quad + \int_{x_-}^{x_- + \xi} |y_{\mu'}^-| + |y_\mu^-| dx + \int_{x_- + \xi}^{x_+ - \xi} |y_{\mu'}^- - y_\mu^-| dx + \int_{x_+ - \xi}^{x_+} |y_{\mu'}^-| + |y_\mu^-| dx \\ &= \text{I} + \text{II} + \text{III} + \text{IV} + \text{V} + \text{VI}. \end{aligned}$$

By Lemma 7.2 the functions y_μ^\pm are monotone in μ , thus for all $\tau \in \mathbb{R}$ we have

$$|y_\mu^\pm(\tau)| \leq \max\{|y_{\mu_0}^\pm(\tau)|, |y_{\mu_1}^\pm(\tau)|\}$$

Thus in particular, using the continuity of orbits of (133) in the flow parameter τ , for x near the boundary points x_\pm and all $\mu \in [\mu_0, \mu_1]$ we have $|y_\mu^\pm(x)| \leq C$, where $C > 0$ is a constant depending only on $y_{\mu_0}^\pm(x_\pm)$ and $y_{\mu_1}^\pm(x_\pm)$, or in other words, on the finite winding numbers of the corridors $\mathcal{K}_1(\mu_0)$ and $\mathcal{K}_1(\mu_1)$. Thus, given ϵ , we may choose ξ small enough (depending on ϵ) such that I, III, IV and VI are all less than $\epsilon/6$.

Fixing ξ in this way, by continuous dependence of orbits of the flow (133) on the parameter μ (and since we are on a compact interval in the flow parameter τ), it is possible for $\delta > 0$ to be chosen small enough (depending on ϵ and ξ), such that $|\mu' - \mu| < \delta$ implies $|y_{\mu'}^+(x) - y_\mu^+(x)| < (x_+ - x_-)^{-1}\epsilon/6$ for all $x \in [x_- + \xi, x_+ - \xi]$. Therefore II $< \epsilon/6$. Similarly, V $< \epsilon/6$ as well, and we are done.

Having established the above properties for the signed area $\mathbf{a}(\mu)$, the standard continuity argument can now be applied: By assumption we have $\mathbf{a}(\mu_0) < 0$ and $\mathbf{a}(\mu_1) > 0$. Thus by the Intermediate Value Theorem there exists $\mu_s \in (\mu_0, \mu_1)$ such that $\mathbf{a}(\mu_s) = 0$, which is equivalent to the existence of a connector. Since the corridors $\tilde{\mathcal{K}}_1(\mu)$ by definition always have \tilde{S}_\pm on their boundary, the saddle connector also has to go from \tilde{S}_- to \tilde{S}_+ . \square

We will use this proposition to establish the existence of saddle connectors for both the Θ (118) and the Ω (122) equations. In each case we need to show that the flow satisfies the assumptions we have made in this subsection about the flow on a cylinder (133).

7.3.6 Existence of saddle connectors for the Θ equation

For simplicity we are only going to consider the case $\kappa = \frac{1}{2}$. To see that (118) is a flow on a cylinder of the type we have considered in the above, we make the following identifications: $x = \theta$, $y = \Theta$, $x_- = 0$, $x_+ = \pi$. We have

$$f(\theta) = \sin \theta, \quad g(\theta, \Theta) = -2a \sin \theta \cos \theta \cos \Theta + (2aE \sin^2 \theta - 1) \sin \Theta + 2\lambda \sin \theta.$$

Therefore, the equilibria are at

$$S_- = (0, 0), \quad N_- = (0, \pi) \quad S_+ = (\pi, -\pi), \quad N_+ = (\pi, 0);$$

all of these are hyperbolic. Condition (131) is clearly satisfied: $s_- - n_- = -\pi$ and $n_+ - s_+ = \pi$.

Here we make a note of the fact, easily verified, that the Θ flow (118) possesses a discrete symmetry: Let $\mathcal{O} = (\theta(\tau), \Theta(\tau))$ be any orbit of the flow (118). Then $\mathcal{O}' = (\pi - \theta(-t), \pi - \Theta(-t))$ is also an orbit of (118). This is simply due to the fact that

$$f(\theta) = f(\pi - \theta), \quad g(\theta, \Theta) = g(\pi - \theta, \pi - \Theta).$$

This symmetry will prove useful in constructing corridors of given winding number for the flow. Next we check that the assumptions we made in the previous subsection about the topology of nullclines hold in the case of (118).

7.3.7 Topology of the nullclines

Let $T := \tan(\Theta/2)$. Then,

$$g_{E,\lambda}(\theta, \Theta) = \frac{1}{1+T^2} (2 \sin \theta (\lambda - a \cos \theta) T^2 + 2(2aE \sin^2 \theta - 1) T + 2 \sin \theta (\lambda + a \cos \theta)) =: \frac{q(T)}{1+T^2}$$

where q is a quadratic polynomial in T , whose discriminant we calculate to be

$$\Delta_q := (2aE \sin^2 \theta - 1)^2 - 4\lambda^2 \sin^2 \theta + 4a^2 \sin^2 \theta \cos^2 \theta.$$

Setting $\tau := -\cot \theta$ we thus obtain

$$\Delta_q = \frac{\delta_q(\tau)}{(1+\tau^2)^2}, \quad \delta_q(\tau) = \tau^4 + 2(1 - 2aE + 2a^2 - 2\lambda^2)\tau^2 + (1 - 2aE)^2 - 4\lambda^2$$

Therefore δ_q is a quartic polynomial which is quadratic in $s = \tau^2$. The discriminant of $\delta_q(s)$ is

$$\Delta_{\delta_q} = \lambda^4 + 2a(E - a)\lambda^2 + a^2(a^2 - 2aE + 1) = (\lambda^2 + a(E - a))^2 + a^2(1 - E^2)$$

which is always positive. Thus $\delta_q(s) = 0$ will always have two roots, which will be of opposite signs if $|\lambda| \geq \frac{1}{2} - aE$. In that case there exists $s_1 = \tau_1^2 > 0$ such that $\Delta_q(\pm\tau_1) = 0$, and $\Delta_q(\tau) < 0$ for $\tau \in (-\tau_1, \tau_1)$. It follows that the quadratic equation $q(T) = 0$ will have two roots so long as $|t| > \tau_1$, it will have repeated roots when $t = \pm\tau_1$, and no roots when $|t| < \tau_1$.

If on the other hand $|\lambda| < \frac{1}{2} - aE \leq \frac{1}{2}$ then the roots of $\delta_q(s) = 0$ will be of the same sign. Since the sum of these roots will be $4(\lambda^2 - a^2 - (\frac{1}{2} - aE))$ it is easy to see that both will be negative, thus they will not correspond to real roots of δ_q . Therefore $\Delta_q > 0$ and hence $q(T) = 0$ will always have two roots. Thus the critical value of the parameter λ (thinking of the other parameters a and E as given and fixed) is

$$\lambda_c := -\frac{1}{2} + aE.$$

Now any zero of q will be a zero of g , and thus will give us a point on the nullcline Γ . In addition, g may also have a zero at $\Theta = 0$ or $\pm\pi$, where $T = \pm\infty$. For g to be zero there we need the coefficient of T^2 in q to vanish, i.e. either $\sin \theta = 0$, which will give us the equilibrium points, or $\cos \theta = \lambda/a$ which is impossible so long as

$$\lambda < -a.$$

Under this condition therefore, the nullclines have the topology we assumed in the previous subsection, with λ playing the role of the parameter μ , i.e., given $a \in (0, \frac{1}{2})$, $E \in [0, 1]$, $\lambda \in (-\infty, -a)$ and λ_c as in the above, the topology of Θ -nullclines for the flow (118) changes across $\lambda = \lambda_c$ in the manner described in assumption **(A)**. Figure 5 shows Maple plots of the Θ -nullclines for values of λ below and above the critical value.

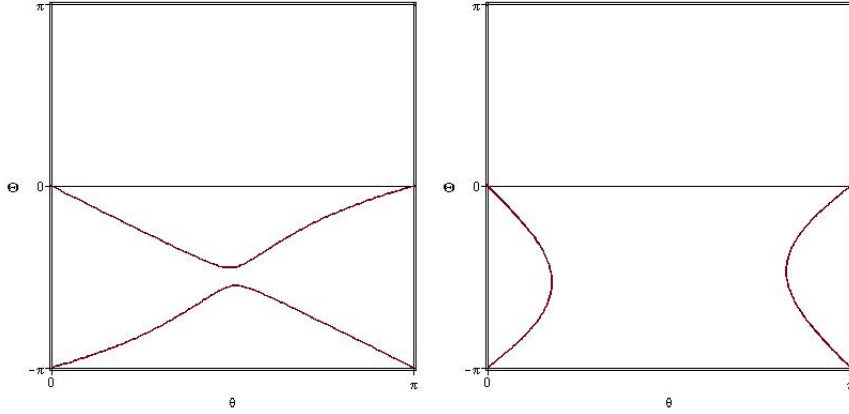


Figure 5: Θ nullclines for $\lambda = -0.4$ (left) and $\lambda = -0.9$ (right), with $a = 0.1$, $E = 0.95$.

7.3.8 Explicit solutions of the Θ equation

One easily verifies that given $a \in [0, \frac{1}{2})$, there is an explicit solution of (114), for $E = 1$ and $\lambda = -1 + a$, namely

$$\Theta_0(\theta) = -\theta.$$

This furthermore generates a saddle connector for (118): $\mathcal{S}_0 := (\theta_0(\tau), \Theta_0(\theta_0(\tau)))$ where $\theta_0(\tau)$ is the unique solution to the ODE $\dot{\theta} = \sin \theta$ with $\theta(-\infty) = 0$ and $\theta(\infty) = \pi$.

This solution will help us get the iteration started.

7.3.9 Existence of corridors with unequal winding number

Throughout this section, a will be a fixed number in $[0, \frac{1}{2})$. The following two propositions will start things off:

PROPOSITION 7.5. *Given $E \in [0, 1]$ and $\lambda \leq \lambda_l := -1 - a$ the flow (118) has a corridor $\mathcal{K}_1(E, \lambda)$ with $w(\mathcal{K}_1) \geq 1$.*

Proof. The linearization of the flow at S_- gives us eigenvalues $\lambda_1 = 1$, $\lambda_2 = -1$ and a corresponding set of eigenvectors is $v_1 = \begin{pmatrix} 1 \\ \lambda - a \end{pmatrix}$ and $v_2 = \begin{pmatrix} 0 \\ 1 \end{pmatrix}$. The orbit \mathcal{W}^- , being the unstable manifold of S_- is tangent to the unstable direction v_1 , therefore the slope of \mathcal{W}^- at S_- is $\frac{d\Theta}{d\theta}|_{\mathcal{W}^-}(S_-) = \lambda - a < 0$. The slope of \mathcal{W}^+ at S_+ is similarly calculated, and it turns out to be the same $\frac{d\Theta}{d\theta}|_{\mathcal{W}^+}(S_+) = \lambda - a$. On the other hand, since $\lambda_l < \lambda_c$, we have $\Gamma = \Gamma_l \cup \Gamma_r$. S_- is a terminal

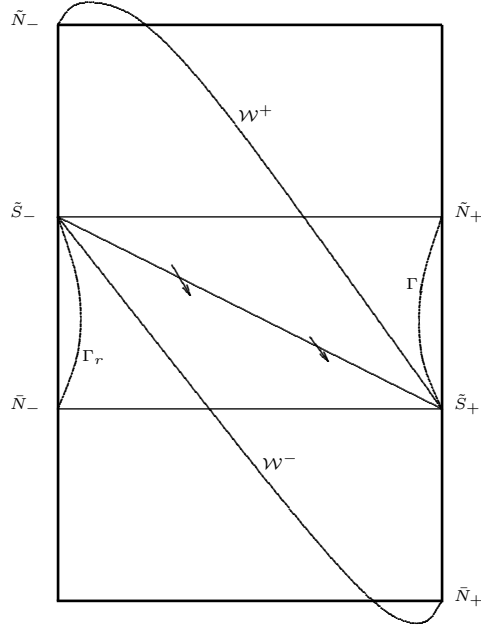


Figure 6: Existence of corridor with $w(\mathcal{K}_1) \geq 1$

point of the curve Γ_l , on which $g_{E,\lambda}(\theta, \Theta) = 0$, hence the slope of Γ_l at S_- can be calculated from implicit function theorem to be

$$\frac{d\Theta}{d\theta}\Big|_{\Gamma_l}(S_-) = -\frac{\partial_\theta g_{E,\lambda}}{\partial_\Theta g_{E,\lambda}}(S_-) = 2(\lambda - a) < \lambda - a.$$

Same is true for the slope of Γ_r at S_+ , as can be easily verified. Thus \mathcal{W}^- starts off inside \mathcal{N} , which is connected, and its initial slope is $\lambda - a < -1$ for $\lambda < -1 + a$. Consider the diagonal line segment S_-S_+ , on which $\Theta = -\theta$. Let us compute the slope of orbits that cross this line, and compare it to the slope of the line:

$$\frac{g_{E,\lambda}(\theta, -\theta)}{f(\theta)} - (-1) = -2a \sin \theta (\cos \theta + E \sin \theta) + 2\lambda + 2 \leq 0 \quad \text{for } \lambda \leq -1 - a.$$

Thus S_-S_+ acts as a “barrier”, not allowing \mathcal{W}^- to cross it from below to above. Hence the ω -limit of \mathcal{W}^- cannot be $(\pi, 0)$. The first possible terminal point is then $(\pi, -2\pi)$, hence $w(\mathcal{W}^-) \geq 1$. \square

PROPOSITION 7.6. *Suppose that for some $E \in (0, 1]$ and some $\lambda \leq \lambda_0 := -1 + a$, the flow (118) has a saddle connector $\mathcal{S}^\Theta(E, \lambda)$ whose lift to the universal cover of the cylinder goes from $\tilde{S}_- = (0, 0)$ to $\tilde{S}_+ = (\pi, -\pi)$. Then, for all $E' \in [0, E)$, there exists a corridor $\mathcal{K}_1(E', \lambda)$ of winding number $w(\mathcal{K}_1) = 0$ for (118).*

Proof. Let $\mathcal{S}(E, \lambda) = (\theta(\tau), \Theta_{E, \lambda}(\tau))$. Since $\mathcal{S}^\Theta = \mathcal{W}^- = \mathcal{W}^+$, by the calculation done in the proof of the previous Proposition, the graph of $\Theta_{E, \lambda}$ is entirely contained in \mathcal{N} , thus it has to be monotone decreasing, and thus $\sin \Theta_{E, \lambda}(\tau) \leq 0$ for all τ .

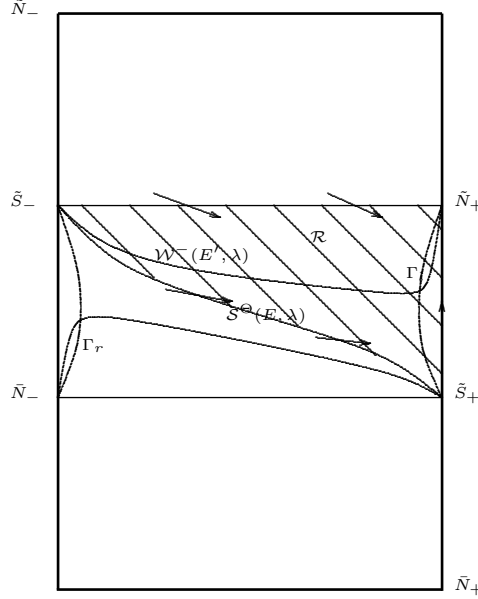


Figure 7: trapped region \mathcal{R} and the corridor with zero winding number.

Consider the region \mathcal{R} in $\tilde{\mathcal{C}}$ whose boundary consists of the following three curves: (i) From S_- to S_+ along $\Theta_{E, \lambda}$. (ii) From S_+ to N_+ along the right boundary, and (iii) from N_+ to S_- along a horizontal line segment. We claim that \mathcal{R} is a trapped region for (118) at parameter values (E', λ) provided $E' < E$. Since (ii) is always an orbit, and (iii) is entirely contained in \mathcal{N} , this only needs to be checked for (i). We have

$$g_{E', \lambda}(\theta, \Theta_{E, \lambda}) = g_{E, \lambda}(\theta, \Theta_{E, \lambda}) + 2a(E' - E) \sin^2 \theta \sin \Theta_{E, \lambda} \geq \dot{\Theta}_{E, \lambda}$$

which shows that the new flow crosses the old solution from left to right. Thus \mathcal{R} is trapped and since \mathcal{W}^- starts in \mathcal{R} , it must terminate in $\tilde{N}_+ = (\pi, 0)$, so that $w(\mathcal{W}^-) = 0$. \square

Setting $E_0 = 1$ and $\lambda_0 = -1 + a$, let $\mathcal{S}_0 := \mathcal{S}(E_0, \lambda_0)$ denote the explicit solution found in 7.3.8. For all $E_1 \in (0, 1)$, the above two propositions, together with the following immediate corollary of Proposition 7.4, establish the existence of a saddle connector $\mathcal{S}_1 := \mathcal{S}(E_1, \lambda_1)$ for (118), for some $\lambda_1 \in (\lambda_l, \lambda_0)$:

COROLLARY 7.7. *Let $E \in [0, 1]$ be fixed. Suppose that there exists $\lambda_1 < \lambda_2 < 0$ such that the flow (118) has corridors $\mathcal{K}_1(E, \lambda_1)$ and $\mathcal{K}_1(E, \lambda_2)$ with $w(\mathcal{K}_1(E, \lambda_2)) = 0$ and $w(\mathcal{K}_1(E, \lambda_1)) \geq 1$. Then there is a $\lambda \in (\lambda_1, \lambda_2)$ such that (118) has a saddle connector $\mathcal{S}^\Theta(E, \lambda)$ going from $(0, 0)$ to $(\pi, -\pi)$.*

Proof. Proposition 7.4 applies, with $-\lambda$ playing the role of the parameter μ . \square

Proceeding iteratively, suppose that given $E_n \in [0, 1]$ a saddle connector $\mathcal{S}_n^\Theta = \mathcal{S}^\Theta(E_n, \lambda_n)$ has been found for (118), for some $\lambda_n < \lambda_0$. In the next subsection we shall see how the newly-found λ_n can be used to prove the existence of a saddle connector for the Ω flow (122), namely $\mathcal{S}_n^\Omega := \mathcal{S}_n^\Omega(E_{n+1}, \lambda_n)$ for some $E_{n+1} \in (0, 1)$. Coming back to the Θ flow then, a new saddle connector \mathcal{S}_{n+1}^Θ needs to be found with the updated energy E_{n+1} , *given that a saddle connector $\mathcal{S}_n^\Theta(E_n, \lambda_n)$ already exists*. Since E_{n+1} can be on either side of E_n , in addition to Prop. 7.6 we also need the following:

THEOREM 7.8. *Given any $a \in [0, \frac{1}{2})$ and $E \in [0, 1]$, there exists a unique*

$$\lambda = \Lambda(E) \in [-1 - a, -1 + a]$$

such that (118) has a saddle connector $\mathcal{S}(E, \lambda)$ going from $(0, 0)$ to $(\pi, -\pi)$. Moreover, Λ is an increasing C^1 function, and $\frac{\partial \Lambda}{\partial E} < a$.

Proof. If $E = 1$ then $\lambda = -1 + a$ works and $\mathcal{S} = \mathcal{S}_0 = (\theta, -\theta)$. For $E < 1$ existence is guaranteed by Prop. 7.6, Prop. 7.5, and Corollary 7.7. To prove uniqueness, suppose that for a given E , there are two saddle connectors $\mathcal{S}(E, \lambda)$ and $\mathcal{S}'(E, \lambda')$, with $\lambda' > \lambda$. Let $\Theta_{E, \lambda}$ and $\Theta_{E, \lambda'}$ denote the corresponding Θ -components of \mathcal{S} and \mathcal{S}' , respectively. For $\theta \in (0, \pi)$,

$$g_{E, \lambda'}(\theta, \Theta_{E, \lambda}) = g_{E, \lambda}(\theta, \Theta_{E, \lambda}) + 2(\lambda' - \lambda) \sin \theta > \dot{\Theta}_{E, \lambda}.$$

Thus orbits of the (E, λ') flow can only cross \mathcal{S} from below to above. On the other hand, since \mathcal{S}' is a saddle connector, near S_- it coincides with $\mathcal{W}^-(E, \lambda')$, and near S_+ it coincides with \mathcal{W}^+ . Thus from the linearization of the flow at S_\pm ,

$$\left. \frac{d\Theta_{E, \lambda'}}{d\theta} \right|_{S_\pm} = \lambda' - a > \lambda - a.$$

Therefore \mathcal{S}' must be above \mathcal{S} near S_- and below it near S_+ , so \mathcal{S}' would have to cross \mathcal{S} from above to below, which is a contradiction, unless they coincide.

Given E then, let $\Lambda(E)$ denote the unique value of λ for which a saddle connector $\mathcal{S}(E, \lambda)$ exists. The fact that Λ is continuously differentiable (in fact analytic), and the bound on the derivative, have already been shown in [38] and [3] using analytic perturbation theory. Here we give a simple proof of monotonicity of Λ which also establishes the bound on the derivative:

Given $E \in [0, 1]$ let $\lambda = \Lambda(E)$ and let $(\theta(\tau), \Theta_E(\tau))$ denote the unique (modulo translations in τ) saddle connector for (118) whose existence we have established. Let $u := \frac{\partial \Theta_E}{\partial E}$. By differentiating the Θ equation in (118) with respect to E we obtain an equation for u :

$$\frac{du}{d\tau} = P(\tau)u + Q(\tau), \quad \begin{cases} P & := 2a \sin \theta \cos \theta \sin \Theta_E + (2aE \sin^2 \theta - 1) \cos \Theta_E \\ Q & := 2a \sin^2 \theta \sin \Theta_E + 2 \frac{d\Lambda}{dE} \sin \theta. \end{cases} \quad (143)$$

Let

$$U(\tau_1, \tau_2) := e^{-\int_{\tau_1}^{\tau_2} P(\tau) d\tau}.$$

Thus we have

$$U(\tau_2, \tau_1) = \frac{1}{U(\tau_1, \tau_2)}, \quad U(\tau_1, \tau_2)U(\tau_2, \tau_3) = U(\tau_1, \tau_3).$$

Solving the first-order linear ODE (143) for u we obtain

$$u(\tau) = U(\tau, \tau_1)u(\tau_1) + \int_{\tau_1}^{\tau} U(\tau, \tau')Q(\tau')d\tau'. \quad (144)$$

Note that P and Q are bounded functions of τ and

$$\lim_{\tau \rightarrow -\infty} P(\tau) = -1, \quad \lim_{\tau \rightarrow \infty} P(\tau) = 1.$$

Therefore, for any fixed τ ,

$$U(\tau, \tau_1) \rightarrow 0 \text{ as } \tau_1 \rightarrow -\infty, \quad U(\tau, \tau_1) \rightarrow 0 \text{ as } \tau_1 \rightarrow \infty.$$

Moreover $u(\pm\infty) = 0$. For any finite τ from (144) we thus obtain two equivalent expressions for $u(\tau)$. For example, setting $\tau = 0$,

$$u(0) = \int_{-\infty}^0 U(0, \tau')Q(\tau')d\tau' = - \int_0^{\infty} U(0, \tau')d\tau'.$$

Thus in particular

$$0 = \int_{-\infty}^{\infty} U(0, \tau')Q(\tau')d\tau' = 2a \int_{\mathbb{R}} U(0, \tau') \sin^2 \theta(\tau') \sin \Theta_E(\tau') d\tau' + 2 \frac{d\Lambda}{dE} \int_{\mathbb{R}} U(0, \tau') \sin \theta(\tau') d\tau'$$

Therefore

$$\frac{d\Lambda}{dE} = a \frac{\int_{\mathbb{R}} U(0, \tau') \sin^2 \theta(\tau') (-\sin \Theta_E(\tau')) d\tau'}{\int_{\mathbb{R}} U(0, \tau') \sin \theta(\tau') d\tau'} \geq 0.$$

We have already shown that $\Theta_E(\tau) \in (-\pi, 0)$ and $\theta(\tau) \in (0, \pi)$ for all τ . Thus the numerator in the above fraction is strictly less than the denominator, hence

$$0 \leq \frac{\partial \Lambda}{\partial E} < a \quad (145)$$

□

7.3.10 Existence of saddle connectors for the Ω equation

We now show that the flow (122) also satisfies all the hypotheses we had made about flows on a cylinder. Once again, for simplicity we are only going to consider the case $\kappa = \frac{1}{2}$. The situation is somewhat more complicated than what we have done in the above for the Θ equation, due to the presence of an extra parameter, namely γ , which breaks the symmetry that was present for the Θ flow, as well as the fact that the equilibria of the Ω flow are degenerate (non-hyperbolic).

Let us make the identifications $x = \xi$ and $y = \Omega$. Thus $x_- = -\pi/2$, $x_+ = \pi/2$, and we now have

$$f(\xi) = \cos^2 \xi, \quad g_{E,\lambda}(\xi, \Omega) = 2a \sin \xi \cos \Omega + 2\lambda \cos \xi \sin \Omega + 2\gamma \sin \xi \cos \xi + \cos^2 \xi - 2aE$$

Therefore, for $E \in [0, 1)$,

$$s_- = -\pi + \cos^{-1}(E), \quad n_- = \pi - \cos^{-1}(E), \quad s_+ = -\cos^{-1}(E), \quad n_+ = \cos^{-1}(E),$$

where by \cos^{-1} we mean the principal branch of the arccosine, $0 \leq \cos^{-1} x \leq \pi$, and

$$S_{\pm} = (\pm \frac{\pi}{2}, s_{\pm}), \quad N_{\pm} = (\pm \frac{\pi}{2}, n_{\pm})$$

as before. We note that this time, all the equilibria are non-hyperbolic, since $f'(\pm\pi/2) = 0$, and that for $\gamma = 0$, there is a discrete symmetry: $f(-\xi) = f(\xi)$ and $g|_{\gamma=0}(-\xi, \pi - \Omega) = g|_{\gamma=0}(\xi, \Omega)$, which is broken when γ is turned on.

Also in the case $E = 1$ there is a further degeneracy: the two equilibria on each side coalesce into one singular point with both eigenvalues equal to zero. For these type of singular points center manifolds can be non-unique, so that the distinguished orbits \mathcal{W}^\pm and the index theory we have developed for the corridor they form, are not directly relevant to this case.

We now check the hypotheses about the topology of the nullclines.

7.3.11 Topology of the null-clines

Let $T := \tan(\Omega/2)$. Then

$$g_{E,\lambda}(\xi, \Omega) = \frac{2q(T)}{1 + T^2},$$

where

$$q(T) := \left(\gamma \sin \xi \cos \xi + \frac{1}{2} \cos^2 \xi - aE - a \sin \xi \right) T^2 + 2\lambda \cos \xi T + \left(\gamma \sin \xi \cos \xi + \frac{1}{2} \cos^2 \xi - aE + a \sin \xi \right).$$

Thus q is a quadratic polynomial in T with coefficients that are functions of ξ . The discriminant of q is

$$\Delta_q(\xi) := \lambda^2 \cos^2 \xi - \left(\gamma \sin \xi \cos \xi + \frac{1}{2} \cos^2 \xi - aE \right)^2 + 4a^2 \sin^2 \xi.$$

Let $\tau := \tan \xi$. Thus $-\infty < \tau < \infty$, and

$$\Delta_q(\tau) = \frac{p(\tau)}{(1 + \tau^2)^2}$$

with

$$p(\tau) := a^2(1 - E^2)\tau^4 + 2\gamma aE\tau^3 + \left(\lambda^2 - \gamma^2 + a^2 + 2a\left(\frac{1}{2} - aE\right) \right) \tau^2 - 2\gamma\left(\frac{1}{2} - aE\right)\tau + \lambda^2 - \left(\frac{1}{2} - aE\right)^2.$$

Consider first the case $0 \leq E < 1$. Then p is an irreducible quartic in τ . Let us write it as $p(\tau) = \sum_{j=0}^4 c_j \tau^j$. Suppose that

$$a \in [0, \frac{1}{2}), \quad \gamma \in (-\frac{1}{2}, 0), \quad \lambda < -\frac{1}{2}. \quad (146)$$

It then follows that

$$c_4 > 0, \quad c_3 < 0, \quad c_2 > 0, \quad c_1 > 0, \quad c_0 > 0.$$

By Descartes' Rule of Signs, then, $p(\tau)$ has either two or no real positive roots, and either two or no real negative roots, counting multiplicity. For a more accurate root count, one needs to use the discriminant of the quartic. Since the discriminant theory for general quartics is somewhat complicated, here we opt for a simpler analysis by estimating p from above and below with two *reducible* quartics. To this end, first we note that $p(\tau) = Q(\tau) - (q_1(\tau))^2$ with

$$Q(\tau) = \lambda^2 + (\lambda^2 + a^2)\tau^2 + a^2\tau^4, \quad q_1(\tau) = \frac{1}{2} - aE + \gamma\tau - aE\tau^2.$$

Thus on the one hand, by completing the square,

$$Q(\tau) = \left(a\tau^2 + \frac{\lambda^2 + a^2}{2a} \right)^2 - \frac{(\lambda^2 - a^2)^2}{4a^2} \leq (q_2^+(\tau))^2, \quad q_2^+(\tau) := a\tau^2 + \frac{\lambda^2 + a^2}{2a}$$

and on the other hand,

$$Q(\tau) \geq \lambda^2 + 2|\lambda|a\tau^2 + a^2\tau^4 = (q_2^-(\tau))^2, \quad q_2^-(\tau) = a\tau^2 + |\lambda|$$

Thus we have upper and lower bounds for p in terms of factorizable quartics $Q^\pm(\tau)$:

$$Q^-(\tau) := (q_2^-(\tau))^2 - (q_1(\tau))^2 \leq p(\tau) \leq (q_2^+(\tau))^2 - (q_1(\tau))^2 =: Q^+(\tau)$$

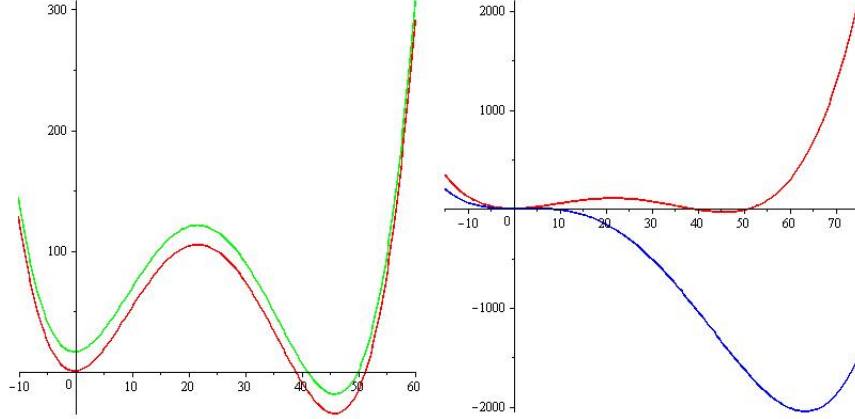


Figure 8: The bounding quartics: Q^+ and p (left), p and Q^- (right) for parameter values $a = 0.1$, $\gamma = -0.2$, $\lambda = -0.9$, $E = 0.978$.

Consider first the upper quartic Q^+ . We have $Q^+ = (q_2^+ + q_1)(q_2^+ - q_1) = q^+q^-$ where q^\pm are two quadratic polynomials. It is clear that if q^+ has any real roots, they will be positive, and if q^- has any roots, they will be negative (recall that we are assuming $\gamma < 0$). The discriminants of q^\pm are computed to be

$$\Delta_\pm := \gamma^2 - 2(1 \mp E)(\lambda^2 + a^2 \pm a(1 - 2aE)).$$

Similarly, $Q^- = (q_2^- + q_1)(q_2^- - q_1) = \tilde{q}^+\tilde{q}^-$ where \tilde{q}^\pm are two quadratics, and once again, any real roots of \tilde{q}^+ must be positive and any real root of \tilde{q}^- , negative. The discriminants of \tilde{q}^\pm are

$$\tilde{\Delta}_\pm = \gamma^2 - 2(1 \mp E)(-2a\lambda \pm a(1 - 2aE)) \geq \Delta_\pm$$

It thus follows that there are two subsets of the (a, γ, λ) parameter space (146) that are of interest: (R1) where both $\tilde{\Delta}_+$ and $\tilde{\Delta}_-$ are negative; and (R2) where $\Delta_+ > 0$ and $\tilde{\Delta}_- < 0$. For parameter values in the region (R1) the quartic Q^- will have no real zeros, and will be always positive, while for those in (R2) the quartic Q^+ will have exactly two positive roots and no negative root.

Let us fix a, γ, λ as in (146). We find that the range (R1) corresponds to $0 \leq E \leq E_l$, where

$$E_l(\lambda) := \begin{cases} \frac{1}{2a} \left[\lambda - a + \frac{1}{2} + \sqrt{(\lambda + a + \frac{1}{2})^2 + \gamma^2} \right] & 0 < -\gamma \leq \sqrt{2a(-2\lambda - 1)} \\ \frac{1}{2a} \left[-\lambda + a + \frac{1}{2} - \sqrt{(\lambda + a - \frac{1}{2})^2 + \gamma^2} \right] & \sqrt{2a(-2\lambda - 1)} < -\gamma \leq \sqrt{2a(-2\lambda + 1)} \\ 0 & \sqrt{2a(-2\lambda + 1)} < -\gamma < \frac{1}{2} \end{cases}$$

while range (R2) corresponds to $E_h < E \leq 1$, with

$$E_h(\lambda) = \frac{1}{4a^2} \left[\lambda^2 + 3a^2 + a - \sqrt{(\lambda^2 - a^2 + a)^2 + 4a^2\gamma^2} \right]$$

Note that $0 < E_h < 1$ and $0 \leq E_l < E_h$ for all values of a, γ, λ as in (146).

For $E > E_h$, therefore, since Q^+ has two positive roots, the quartic $p(\tau)$ must also have at least two roots, one of which will definitely be positive. Thus by the Rule of Signs, $p(\tau)$ has exactly two positive roots. We call them τ_1 and τ_2 , and $p(\tau) < 0$ for $\tau_1 < t < \tau_2$. It follows that the quadratic $q(T)$ will have two roots for $t \notin [\tau_1, \tau_2]$, double roots at τ_1 and at τ_2 , and no real roots for $\tau \in (\tau_1, \tau_2)$.

For $E \in [0, E_l]$ since Q^- has no real roots, p cannot have any either. Thus p is always positive and $q(T) = 0$ will have two roots for all $\tau \in \mathbb{R}$.

Combining these two, one concludes that a critical value for the energy $E = E_c(a, \gamma, \lambda)$ exists, $E_c \in (E_l, E_h)$, such that assumption **(A)** is satisfied, with the role of parameter μ played by E .

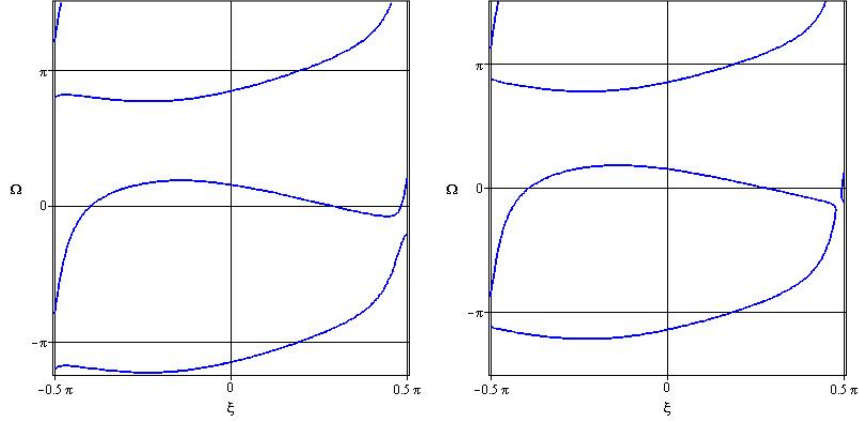


Figure 9: Ω -nullclines for parameter values $E = 0.8$ (left) and $E = 0.93$ (right), with $a = 0.1$, $\gamma = -0.4$, and $\lambda = -0.9$.

7.3.12 Existence of corridors with unequal winding number

Throughout this section, a will be a fixed number in $[0, \frac{1}{2})$ and γ a fixed number in $(-\frac{1}{2}, 0)$. The following two propositions help us get started:

PROPOSITION 7.9. *Given $\lambda \leq -1 + a$ there exists $\bar{E} \in (E_h(\lambda), 1)$ such that for all $E \in [\bar{E}, 1)$ the corridor $\mathcal{K}_1(E, \lambda)$ of the flow (122) has winding number greater than or equal to one.*

Proof. We compute the slope of solution orbits that cross the following line in $\bar{\mathcal{C}}$

$$L := \{(\xi, \Omega) \in \bar{\mathcal{C}} \mid -\frac{\pi}{2} \leq \xi \leq \frac{\pi}{2}, \quad \Omega = \frac{\pi}{2} - \cos^{-1} E - \xi\}$$

and compare it to the slope of L . Note that L passes through \tilde{N}_- and \tilde{S}_+ . We have

$$\frac{g_{E,\lambda}(\xi, \Omega)}{f(\xi)} - (-1) = 2 \left(1 - (a - \lambda)E + [\sqrt{1 - E^2}(a - \lambda) + \gamma] \tan \xi \right) \quad (147)$$

Consider first the case $\lambda = -1 + a$. Let

$$\xi_0(E) := \tan^{-1} \frac{1 - E}{-\gamma - \sqrt{1 - E^2}} \searrow 0 \text{ as } E \nearrow 1.$$

For $\xi \geq \xi_0(E)$, the slope of any orbit of the flow crossing L is less than (i.e. more negative than) the slope of L . Hence on the portion of L where $\xi \geq \xi_0$ orbits can only cross L from above to below.

On the other hand, suppose $\lambda < -1 + a$. Let

$$E_m := \max \left\{ \frac{1}{a - \lambda}, \sqrt{1 - \frac{\gamma^2}{(a - \lambda)^2}}, E_h(\lambda) \right\} \in (E_h, 1)$$

For $E \in (E_m, 1)$ we have $\xi_0(E) \in (-\frac{\pi}{2}, 0)$. Let

$$\eta_0 := \begin{cases} \cos^{-1} E & \lambda < -1 + a \\ \cos^{-1} E + \xi_0(E) & \lambda = -1 + a \end{cases}$$

We note that $\eta_0 \rightarrow 0$ as $E \rightarrow 1$. Let us consider the horizontal line

$$L' := \{(\xi, \Omega) \in \bar{\mathcal{C}} \mid -\frac{\pi}{2} \leq \xi \leq \frac{\pi}{2}, \quad \Omega = \frac{\pi}{2} - \eta_0\}$$

We compute the slope of orbits crossing L' :

$$h(\xi) := g_{E,\lambda}(\xi, \frac{\pi}{2} - \eta_0) = 2a \sin \xi \sin \eta_0 + 2\lambda \cos \xi \cos \eta_0 + 2\gamma \sin \xi \cos \xi + \cos^2 \xi - 2aE.$$

Clearly

$$\begin{aligned} h(\xi) &\leq 2a(\sin \eta_0 - E) + |\gamma| + 2\lambda \cos \xi \cos \eta_0 + \cos^2 \xi \\ &\leq 2a(\sin \eta_0 - E) + |\gamma| + 2\lambda \cos \eta_0 + 1 \\ &\rightarrow -2a + |\gamma| + 2\lambda + 1 \quad \text{as } E \rightarrow 1 \\ &\leq |\gamma| - 1 < 0 \end{aligned}$$

Thus there exists E sufficiently close to 1 such that $h(\xi) < 0$ for $\xi \in (-\frac{\pi}{2}, \frac{\pi}{2})$.

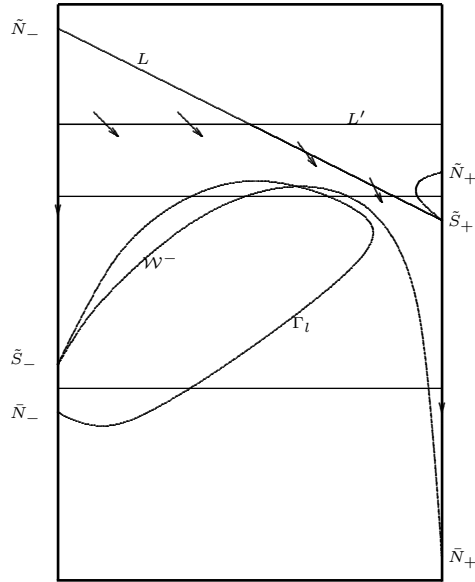


Figure 10: Construction of a barrier

Let γ be the curve in \mathcal{C} consisting of two line segments: From $(-\frac{\pi}{2}, \frac{\pi}{2} - \eta_0)$ along the horizontal line L' , upto the intersection point of L and L' , and then from there along L to \tilde{S}_+ . The curve γ provides a barrier for the flow: no orbit can cross it from the region below γ into the region above γ . Let $\tilde{\mathcal{W}}^-$ be the unstable manifold of \tilde{S}_- . Thus $\tilde{\mathcal{W}}^-$ must stay below γ , and as a result the ω -limit of \mathcal{W}^- cannot be \tilde{N}_+ so therefore its winding number is not zero or negative. \square

PROPOSITION 7.10. *Given $\lambda \leq -1 + a$ and $E \in [0, E_l(\lambda)]$ the corridor $\mathcal{K}_1(E, \lambda)$ of the flow (122) has winding number equal to zero.*

Proof. Once again we find a barrier that prevents \mathcal{W}^- from going down: Let us compute the slope of orbits crossing the line

$$L = \{(\xi, \Omega) \mid \Omega = \xi - \frac{\pi}{2}\}$$

and compare it to the slope of this line.

$$j(\xi) := g_{E,\lambda}(\xi, \xi - \frac{\pi}{2}) - \cos^2 \xi = 2a(1 - E) - 2(\lambda + a) \cos^2 \xi + 2\gamma \sin \xi \cos \xi.$$

Thus $j(\pm\pi/2) = 2a(1 - E) > 0$. Any interior minimum of j must be achieved at a critical point:

$$j'(\xi_0) = 0 \implies \xi_0 = \frac{1}{2} \tan^{-1} \frac{\gamma}{\lambda + a} > 0.$$

However we have

$$j(\xi_0) = 2a(1 - E) - (\lambda + a) - (\lambda + a) \cos 2\xi_0 + \gamma \sin 2\xi_0 = 2a(1 - E) - \frac{2(\lambda + a)^3}{(\lambda + a)^2 + \gamma^2} > 2a(1 - E).$$

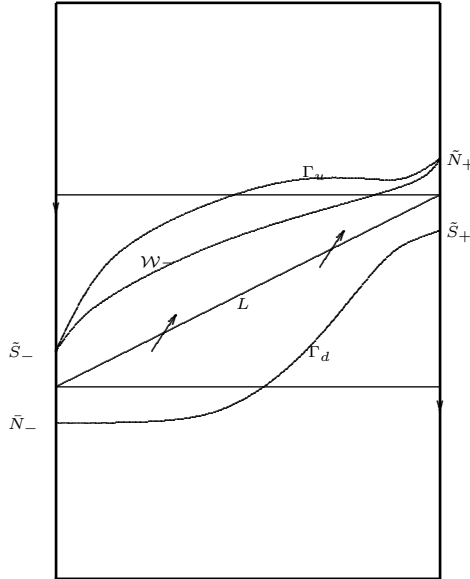


Figure 11: Barrier for zero winding number

Thus $j(\xi) > 0$ for all $\xi \in (-\frac{\pi}{2}, \frac{\pi}{2})$. It follows that the slope of any orbit crossing the line is greater than the slope of the line. Thus orbits cannot cross this line from above to below. In particular, the orbit \mathcal{W}^- starts at \tilde{S}_- , which is above this line. Hence \mathcal{W}^- cannot end at any

copy of the node N_+ other than \tilde{N}_+ , so its winding number cannot be positive. Since the region \mathcal{N} is connected, once \mathcal{W}^- leaves \mathcal{P} and enters \mathcal{N} , its Ω must decrease, hence \mathcal{W}^- cannot end at any copy of the node N_+ that is higher than \tilde{N}_+ either, and therefore the winding number of \mathcal{W}^- cannot be negative, hence $w(\mathcal{W}^-) = 0$. \square

Let $\lambda_0 = -1 + a$. The above two propositions, in conjunction with the following immediate corollary of Proposition 7.4, establish the existence a saddle connector $\mathcal{S}_0^\Omega(E_1, \lambda_0)$ for the flow (122), for some $E_1 \in (0, 1)$:

COROLLARY 7.11. *Let $\lambda \leq -1 + a$ be fixed. Suppose that there exists $0 \leq E_1 < E_2 < 1$ such that the flow (122) has corridors $\mathcal{K}_1(E_1, \lambda)$ and $\mathcal{K}_1(E_2, \lambda)$ with $w(\mathcal{K}_1(E_1, \lambda)) = 0$ and $w(\mathcal{K}_1(E_2, \lambda)) \geq 1$. Then there is an $E \in (E_1, E_2)$ such that (122) has a saddle connector $\mathcal{S}(E, \lambda)$.*

Proof. Proposition 7.4 applies, with E playing the role of the parameter μ . \square

Let $n \geq 1$. Suppose that given $\lambda_{n-1} \leq -1 + a$ a saddle connector $\mathcal{S}_{n-1}^\Omega = \mathcal{S}^\Omega(E_n, \lambda_{n-1})$ has been found for (122), for some $E_n \in (0, 1)$. In the previous subsection we saw how this newly-found E_n can be used to prove the existence of a saddle connector for the Θ flow (118), namely $\mathcal{S}_n^\Theta := \mathcal{S}_n^\Theta(E_n, \lambda_n)$ for some $\lambda_n \leq -1 + a$. Coming back to the Ω flow then, given the updated value $\lambda = \lambda_n$, a new saddle connector \mathcal{S}_n^Ω needs to be found with an updated energy E_{n+1} , *given that a saddle connector $\mathcal{S}_{n-1}^\Omega(E_n, \lambda_{n-1})$ already exists.* More generally, we have

THEOREM 7.12. *Fix $a \in [0, \frac{1}{2})$ and $\gamma \in (-\frac{1}{2}, 0)$. Then given any $\lambda \in [-1 - a, -1 + a]$, there exists a unique*

$$E = \mathcal{E}(\lambda) \in (E_l(\lambda), 1)$$

such that (122) has a saddle connector $\mathcal{S}^\Omega(E, \lambda)$. Moreover, \mathcal{E} is a C^1 function, and $|\frac{\partial \mathcal{E}}{\partial \lambda}| < \frac{1}{a}$.

Proof. Existence of a saddle connector is guaranteed by Propositions 7.9 and 7.10, and Corollary 7.11. To see uniqueness, suppose that there exists two saddle connectors $\mathcal{S}^\Omega(E, \lambda)$ and $\mathcal{S}^\Omega(E', \lambda)$ for (122) for E and E' in $(E_b(\lambda), 1)$, and suppose $E < E'$. Let $\Omega_{E, \lambda}$ be the Ω component of $\mathcal{S}^\Omega(E, \lambda)$. We have

$$g_{E', \lambda}(\xi, \Omega_{E, \lambda}) = g_{E, \lambda}(\xi, \Omega_{E, \lambda}) - 2a(E' - E) < \dot{\Omega}_{E, \lambda}.$$

It thus follows that orbits of the (E', λ) flow can only cross $\mathcal{S}^\Omega(E, \lambda)$ from above to below. On the other hand, since $E' > E$ the equilibrium point $\tilde{S}_-(E', \lambda)$ is situated below $\tilde{S}_-(E, \lambda)$, while $\tilde{S}_+(E', \lambda)$ is above $\tilde{S}_+(E, \lambda)$. Since $\mathcal{S}^\Omega(E', \lambda)$ coincides with both $\mathcal{W}^-(E', \lambda)$ and $\mathcal{W}^+(E', \lambda)$ it begins below $\mathcal{S}^\Omega(E, \lambda)$ and it ends above it, which is a contradiction, hence $E' = E$.

Given λ , let $\mathcal{E}(\lambda)$ denote the unique value of E for which a saddle connector $\mathcal{S}^\Omega(\mathcal{E}(\lambda), \lambda) = (\xi(\tau), \Omega_{\mathcal{E}(\lambda), \lambda}(\tau))$ exists. We now prove that \mathcal{E} is a C^1 function: Consider the two initial value problems for $\Omega^\pm(\tau)$, the Ω components of \mathcal{W}^\pm :

$$\dot{\Omega}^\pm = g_{E, \lambda}(\xi(\tau), \Omega^\pm), \quad \Omega^-(-\infty) = -\pi + \cos^{-1} E, \quad \Omega^+(\infty) = -\cos^{-1} E$$

By standard ODE theory these two problems have unique smooth solutions $\Omega_{E, \lambda}^\pm(\tau)$ which also depend smoothly on the parameters λ and E (so long as $E < 1$) for any finite τ .

Next recall that \mathcal{W}^- is a saddle connector if it coincides with \mathcal{W}^+ , which will be the case if these two orbits intersect at one point, e.g. if $\Omega_{E, \lambda}^+(\tau) = \Omega_{E, \lambda}^-(\tau)$. Let us define a smooth function

$$\Phi(E, \lambda) := \Omega_{E, \lambda}^+(\tau) - \Omega_{E, \lambda}^-(\tau) \tag{148}$$

Let $\lambda_0 \in [-1 - a, -1 + a]$ be fixed, and set $E_0 = \mathcal{E}(\lambda_0)$. Then $\Phi(E_0, \lambda_0) = 0$. By the Implicit Function Theorem, if

$$\frac{\partial \Phi}{\partial E}(E_0, \lambda_0) \neq 0, \quad (149)$$

then there is a neighborhood \mathcal{I} of λ_0 and a C^1 function $\tilde{\mathcal{E}}$ defined on \mathcal{I} such that $\Phi(\tilde{\mathcal{E}}(\lambda), \lambda) = 0$ for all $\lambda \in \mathcal{I}$. By the uniqueness result we have already shown, we must have $\tilde{\mathcal{E}} = \mathcal{E}$. Thus we only need to verify the condition (149).

For $\lambda \in [-1 - a, -1 + a]$ and $E \in (E_l(\lambda), 1)$, let

$$u_{\pm}(\tau) := \frac{\partial}{\partial E} \Omega_{E,\lambda}^{\pm}(\tau).$$

Then u_{\pm} satisfy the linear ODEs

$$\frac{du_{\pm}}{d\tau} = P_{\pm}(\tau)u_{\pm} - 2a, \quad P_{\pm}(\tau) := -2a \sin \xi(\tau) \sin \Omega_{E,\lambda}^{\pm}(\tau) + 2\lambda \cos \xi(\tau) \cos \Omega_{E,\lambda}^{\pm}(\tau) \quad (150)$$

together with the initial conditions

$$u_{-}(-\infty) = \frac{-1}{\sqrt{1-E^2}}, \quad u_{+}(\infty) = \frac{1}{\sqrt{1-E^2}}.$$

Moreover,

$$\Phi(E, \lambda) = u_{+}(0) - u_{-}(0).$$

For $\tau_1, \tau_2 \in \mathbb{R}$ let

$$U_{\pm}(\tau_1, \tau_2) := e^{-\int_{\tau_1}^{\tau_2} P_{\pm}(\tau) d\tau}.$$

Solving the ODEs (150) for u_{\pm} we obtain

$$U_{\pm}(\tau_1, \tau_2)u_{\pm}(\tau_2) = u_{\pm}(\tau_1) - 2a \int_{\tau_1}^{\tau_2} U_{\pm}(\tau_1, \tau) d\tau. \quad (151)$$

Note that

$$\lim_{\tau \rightarrow -\infty} P_{-}(\tau) = -2a\sqrt{1-E^2} < 0, \quad \lim_{\tau \rightarrow \infty} P_{+}(\tau) = 2a\sqrt{1-E^2} > 0.$$

Thus for any fixed τ ,

$$U_{-}(\tau, \tau_1) \rightarrow 0 \text{ as } \tau_1 \rightarrow -\infty, \quad U_{+}(\tau, \tau_2) \rightarrow 0 \text{ as } \tau_2 \rightarrow \infty.$$

And so, from (151) we obtain

$$u_{-}(\tau) = -2a \int_{-\infty}^{\tau} U_{-}(\tau, \tau') d\tau', \quad u_{+}(\tau) = 2a \int_{\tau}^{\infty} U_{+}(\tau, \tau') d\tau'. \quad (152)$$

We know that $\Omega_{E_0, \lambda_0}^{+}(\tau) = \Omega_{E_0, \lambda_0}^{-}(\tau)$ for all τ . Hence $U_{+}(\tau_1, \tau_2) = U_{-}(\tau_1, \tau_2) =: U(\tau_1, \tau_2)$ and thus

$$\frac{\partial \Phi}{\partial E}(E_0, \lambda_0) = 2a \int_{-\infty}^{\infty} U(0, \tau') d\tau' > 0$$

so that (149) is clearly satisfied. Since λ_0 was arbitrary we have shown that $\mathcal{E} \in C^1((-1-a, -1+a))$.

We can furthermore compute the derivative of \mathcal{E} by implicit differentiation. Let

$$v_{\pm}(\tau) := \frac{\partial}{\partial \lambda} \Omega_{E,\lambda}^{\pm}(\tau).$$

Then v_{\pm} satisfy

$$\dot{v}_{\pm} = P_{\pm}(\tau)v_{\pm} + 2 \cos \xi(\tau) \sin \Omega_{E,\lambda}^{\pm}(\tau), \quad v_{-}(-\infty) = 0, \quad v_{+}(\infty) = 0.$$

Thus by a similar argument to above,

$$v_{-}(\tau) = \int_{-\infty}^{\tau} U_{-}(\tau, \tau') \cos \xi(\tau') \sin \Omega_{E,\lambda}^{-}(\tau') d\tau', \quad v_{+}(\tau) = \int_{\tau}^{\infty} U_{+}(\tau, \tau') \cos \xi(\tau') \sin \Omega_{E,\lambda}^{+}(\tau') d\tau',$$

so that

$$\frac{\partial \Phi}{\partial \lambda}(E_0, \lambda_0) = -2 \int_{-\infty}^{\infty} U(0, \tau') \cos \xi(\tau') \sin \Omega_{E_0, \lambda_0}(\tau') d\tau',$$

and thus

$$\frac{d\mathcal{E}}{d\lambda} = -\frac{\partial \Phi / \partial \lambda}{\partial \Phi / \partial E} = \frac{\int_{-\infty}^{\infty} U(0, \tau) \cos \xi(\tau) \sin \Omega_{E,\lambda}(\tau) d\tau}{a \int_{-\infty}^{\infty} U(0, \tau) d\tau}.$$

Moreover, clearly

$$\left| \int_{-\infty}^{\infty} U(0, \tau) \cos \xi(\tau) \sin \Omega_{E,\lambda}(\tau) d\tau \right| < \int_{-\infty}^{\infty} U(0, \tau) d\tau.$$

so that

$$\left| \frac{d\mathcal{E}}{d\lambda} \right| < \frac{1}{a}.$$

□

7.3.13 The iteration argument

THEOREM 7.13. *There exists a $\lambda \in (-1 - a, -1 + a)$ and $E \in (0, 1)$ such that (118) has a saddle connector $\mathcal{S}^{\Theta}(E, \lambda)$ and (122) has a saddle connector $\mathcal{S}^{\Omega}(E, \lambda)$.*

Proof. Set $\lambda_0 = -1 + a$. For $n \geq 1$ let

$$E_n := \mathcal{E}(\lambda_{n-1}) \in (0, 1), \quad \lambda_n := \Lambda(E_n) \in (-1 - a, -1 + a).$$

Thus

$$E_{n+1} = \mathcal{E}(\Lambda(E_n)).$$

By Theorems 7.12 and 7.8 we have

$$|E_{n+1} - E_n| \leq \delta |E_n - E_{n-1}|$$

where

$$\delta := \max_{0 \leq E \leq 1} \left| \frac{d\Lambda}{dE} \right| \max_{-1-a \leq \lambda \leq -1+a} \left| \frac{d\mathcal{E}}{d\lambda} \right| < a \cdot \frac{1}{a} = 1$$

Thus by the contraction mapping theorem, the sequence E_n converges, and thus so does the sequence λ_n . Let $\lambda = \lim_{n \rightarrow \infty} \lambda_n \in [-1 - a, -1 + a]$ and $E := \lim_{n \rightarrow \infty} E_n$. We must have $E = \mathcal{E}(\lambda)$ and thus $E < 1$. □

8 Summary and Outlook

In this paper we have studied the Dirac equation for a point electron in static, electromagnetic, flat spacetimes with Zipoy topology which include the zero-gravity limit of the electromagnetic Kerr–Newman spacetimes as special case, but which in general can feature generalization of the Appell–Sommerfeld electromagnetic fields with any charge Q and current I one wants; the zero- G Kerr–Newman spacetimes here correspond to $Q = I\pi a$. In contrast to similar-spirited studies of the Dirac equation for a point electron on the Kerr–Newman spacetime, which are plagued by the presence of a Cauchy horizon and regions of closed timelike loops, [20, 5, 14, 15, 16, 17, 3, 39, 40, 4], our zero- G spacetimes do not possess any such physically troublesome features.

We proved with the help of an operator that anti-commutes with the pertinent Dirac Hamiltonian that the spectrum of any of its self-adjoint extensions is symmetric about zero; this result holds for any charge Q and current I of the electromagnetic spacetimes. By adapting an argument of Winklmeier–Yamada for the Dirac equation of a point electron in the outer region of the Kerr–Newman black hole spacetime, we have shown that the formal Dirac Hamiltonian is essentially self-adjoint on a complete spacelike slice of the maximal analytically extended, static $zGKN$ spacetime. Then we exploited the Chandrasekhar–Page–Toop separation-of-variables theorem for Dirac’s equation on a general Kerr–Newman spacetime, and the Prüfer transform, and showed that the self-adjoint Dirac operator on the $zGKN$ spacetime has a continuous spectrum with a gap about zero that contains a pure point spectrum. The pure point spectrum is associated with time-periodic L^2 spinor fields, representing bound states of Dirac’s point electron in the electromagnetic field of the ring singularity of the $zGKN$ spacetime.

Our results are far from exhaustive. In the following we list a number of interesting open problems which we hope will be solved in some future work.

We begin with the Dirac point electron in $zGKN$ spacetimes:

- **Problem 1** Characterize the point spectrum of the Dirac Hamiltonian on $zGKN$ in complete detail; to the extent possible, compute it analytically, or at least numerically in some representative situations.

Remark 8.1. We suspect that there is a countably infinite set of energy eigenvalues which correspond in a one-to-one fashion with saddle connectors of arbitrary winding numbers $w \in \mathbb{Z}$. The possibility of such saddle connectors we already established, see Theorem 7.13.

- **Problem 2:** Discuss the generalized scattering problem for Dirac spinor fields on the $zGKN$ spacetimes. In particular, investigate the evolution when (part of) the Dirac spinor field “dives” through the ring from, say, the upper to the lower sheet.

We now come to the Dirac point electron in zGK spacetimes equipped with generalization of the Appell–Sommerfeld electromagnetic fields to arbitrary charge Q and current I :

- **Problem 3:** For the formal Dirac Hamiltonian on the (Q, I) -generalization of the $zGKN$ spacetime (given a), show that essential self-adjointness holds if the “coupling constant” $(Q - I\pi a)e$ is small in magnitude; perhaps using a so-called Hardy–Dirac type estimate.
- **Problem 4** Suppose essential self-adjointness fails if $|(Q - I\pi a)e|$ is too large. If so, what is the sharp constant for $|(Q - I\pi a)e|$? Determine the (Q, I) -parameter regimes (given a) in which the Dirac Hamiltonian on a generalization of the $zGKN$ spacetime with Sommerfeld has several self-adjoint extensions, respectively has no self-adjoint extension. Amongst the self-adjoint extensions, can one identify a distinguished one?
- **Problem 5** In the cases of self-adjointness, characterize the spectrum of the Dirac Hamiltonian in complete detail; to the extent possible, compute the spectrum analytically, or at least numerically (in particular the point spectrum) in some representative situations.
- **Problem 6** The continuous spectrum of quantum physical operator families is usually very robust. Show that Theorem 2.7 holds for the (Q, I) generalization of our Dirac operators, at least as long as $(Q - I\pi a)e$ is small in magnitude.
- **Problem 7** Same consideration as above, for Theorem 2.8; thus: Can an eigenvalue of \hat{H} on \mathbb{H} for $Q = I\pi a$ be continuously deformed into an eigenvalue of \hat{H} on \mathbb{H} for $Q \neq I\pi a$ as long as $(Q - I\pi a)e$ is sufficiently small in magnitude? If so, does the size of the neighborhood of $Q = I\pi a$ into which an eigenvalue can be continued depend on the eigenvalue, or can one have a uniform control on the spectrum w.r.t. the coupling constant $(Q - I\pi a)e$? Is it possible that the point spectrum disappears completely if $(Q - I\pi a)e$ becomes too large in magnitude?
- **Problem 8:** Same as Problem 3, now for Dirac spinor fields on zGK equipped with generalizations of the Appell–Sommerfeld fields to arbitrary Q and I .

So far our problems concern the Dirac equation on the zero-gravity limit case of the KN spacetimes, and its generalization to arbitrary Q and I . To make contact with the existing studies of Dirac’s electron in Kerr–Newman spacetimes, the following bifurcation problem suggests itself:

- **Problem 9:** Deform these zero-gravity spacetimes perturbatively by “switching on” G and discuss the Dirac equation on them perturbatively as well. In particular, is it possible to perturb into generalizations of the Kerr–Newman spacetime with gyromagnetic ratios amounting to g -factors $g \neq g_{\text{KN}} = 2$, or is such a perturbation feasible only if $g = g_{\text{KN}} = 2$, viz. if $Q = I\pi a$?

In all the above problems, the energy-momentum-stress tensor is of classical electromagnetic nature. The Dirac spinor field does not influence the spacetime structure (in the zero-gravity limit the electromagnetic fields do not influence the spacetime structure, either). More to the point, Dirac’s point electron is treated as a *test particle* in this paper and in all the above problems. Since test particles are only mathematical fiction, no matter how useful practically, an obvious task is to investigate the Dirac electron *not* as a test particle. Thus:

- **Problem 10:** Treat the quantum-mechanical interaction of the Dirac electron with the ring singularity of the zero- G Kerr–Newman spacetime symmetrically as a two-body problem.

Remark 8.2. Of course, the same problem has not even been solved yet for the simpler setting of “Dirac hydrogen” in Minkowski spacetime; i.e., how to go beyond the traditional textbook problem where the relativistic hydrogen problem is treated by solving the Dirac equation of a point electron in flat Minkowski spacetime containing an *infinitely massive* positive point charge (representing the proton). Typically the problem of finite mass of the proton (or nucleus, more generally) is addressed by perturbation theory, starting from Pauli’s two-body equation and adding “relativistic corrections” in powers of $1/c$; or by perturbative QED-type calculations. In a similar vein one may be able to take a finite “ADM mass” of the ring singularity into account.

- **Problem 12:** Same as Problem 11, but now with the (Q, I) generalizations of $zGKN$.
- **Problem 13:** Same as Problems 11 and 12, but now perturbatively for $G > 0$.
- **Problem 14:** Compute the feedback of the Dirac spinor field onto the spacetime structure perturbatively when $G > 0$. This amounts to the perturbative discussion of the so-called Einstein–Dirac system for small G . In this problem the energy-momentum-stress tensor, in addition to classical electromagnetic fields, also involves the Dirac spinor field which now influences the spacetime structure.

All these problems are difficult, and the amount of work needed to solve all of them can only be handled by the involvement of many mathematical physicists. In this vein we hope that our paper inspires some readers to join us in our pursuit. We ourselves have made some progress on problems 3, 10, and 11, which we plan to report in forthcoming publications.

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