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Relaxation oscillations and canards in the Jirsa-Kelso excitator model: global flow perspective

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Abstract. Fenichel's geometric singular perturbation theory and the blowup method have been very successful in describing and explaining global non-linear phenomena in systems with multiple time-scales, such as relaxation oscillations and canards. Recently, the blowup method has been extended to systems with flat, unbounded slow manifolds that lose normal hyperbolicity at infinity. Here, we show that transition between discrete and periodic movement captured by the Jirsa-Kelso excitator is a new example of such phenomena. We, first, derive equations of the Jirsa-Kelso excitator with explicit time scale separation and demonstrate existence of canards in the systems. Then, we combine the slow-fast analysis, blowup method and projection onto the Poincaré sphere to understand the return mechanism of the periodic orbits in the singular case, $\epsilon = 0$.

1 Introduction

The Jirsa-Kelso excitator model is a class of excitable planar systems proposed as a minimal model to describe generation of rhythmic and discrete human movement [7]. The model predictions concerning mechanism of discrete movement generation have been tested in a number of empirical studies [3, 6].

The model of a single oscillator is given as a system of two ODEs (Eq. (12) in [7]):

$$\begin{aligned} \frac{du}{dt} &= v(t), \\ \frac{dv}{dt} &= (1 - u(t)^2)Tv(t) - u(t) - b\left(\frac{v(t)}{T} - u(t) + \frac{1}{3}u(t)^3\right) + a + I(t). \end{aligned} \tag{1}$$

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Model (1) has mixed time-scales. To separate them we, first, rearrange the second ODE, factor out $-Tv(t)$ and substitute the time constant $T = 1/\sqrt{\epsilon}$, obtaining,

$$\begin{aligned} \frac{du}{dt} &= v(t), \\ \frac{dv}{dt} &= -\frac{b}{3}u(t)^3 + (b-1)u(t) + a + I(t) - (u(t)^2 - 1 + b\epsilon) \frac{v(t)}{\sqrt{\epsilon}}. \end{aligned} \quad (2)$$

In (1) and (2), $u(t)$ is interpreted as a position and $v(t)$ as a velocity of the movement, a and b are intrinsic model parameters, T and ϵ are a time scale constants and $I(t)$ is an external stimulus input (in the rest of the paper we consider the autonomous systems, i.e. $I = 0$). Then, the mixed time-scales of the system (2) can be separated by rescaling the time t and the variable v , respectively:

$$(t/\sqrt{\epsilon}, u, v/\sqrt{\epsilon}) \rightarrow (\tau, x, y). \quad (3)$$

The resulting system, with explicitly separated time-scales,

$$\begin{aligned} x' &= \epsilon y, \\ y' &= -\frac{b}{3}x^3 + (b-1)x + a - (x^2 - 1 + \epsilon b) y, \end{aligned} \quad (4)$$

has the classical *fast* representation, $x' = \epsilon f(x, y, \epsilon)$, $y' = g(x, y, \epsilon)$, where $(\cdot)' = \frac{d}{d\tau}$ denotes differentiation with respect to the fast time τ [9, 10]. In the rest of the paper we refer to (4) as the Jirsa-Kelso excitator (JKE).

The modelling approach introduced in [7] is motivated by excitable systems with time scale separation. In fact, transformation given in [7],

$$\begin{aligned} x &= x, \\ y &= \left(z + x - \frac{x^3}{3} \right) / \epsilon, \end{aligned} \quad (5)$$

allows to transform the JKE (4) into the seminal example of such a system, the FitzHugh-Nagumo model (FHN) (with $I = 0$, $z \rightarrow -z$ and $a \rightarrow -a$) [4, 14, 15],

$$\begin{aligned} x' &= x - \frac{x^3}{3} + z, \\ z' &= \epsilon (a - x - bz). \end{aligned} \quad (6)$$

Figure 1 illustrates the transformation (5) between the FHN (6) and the JKE (4) models. Figure 1(a) shows the hypersurface $H = \{(x, y, z) \in \mathbb{R}^3 : y = (z + x - \frac{x^3}{3}) / \epsilon\}$ together with projection of the common x -nullcline of the JKE (4) and the FHN (6) (blue curve) and y -nullcline of the JKE (4) (red curve). Additionally, it shows two representative periodic orbits (dark, $a \approx 0.78978455$, and light, $a \approx 0.78978433$, grey curves). Figure 1(b)-(d) demonstrates the invariance of the x variable, and the difference in time-courses of variables y and z ; the time-series correspond to the two periodic orbits in Fig.1(a). Although, the variable x is invariant under the transformation (5), it has different physical interpretations in the two models. In the JKE (4) x describes position and hence allows to interpret the variable y as velocity, while in the FHN (6) x has been commonly interpreted as a the membrane potential of a neuronal cell, while y represents a recovery variable associated with a slow dynamics (usually attributed to ion channels' gating) that controls the generation of action potentials [4].

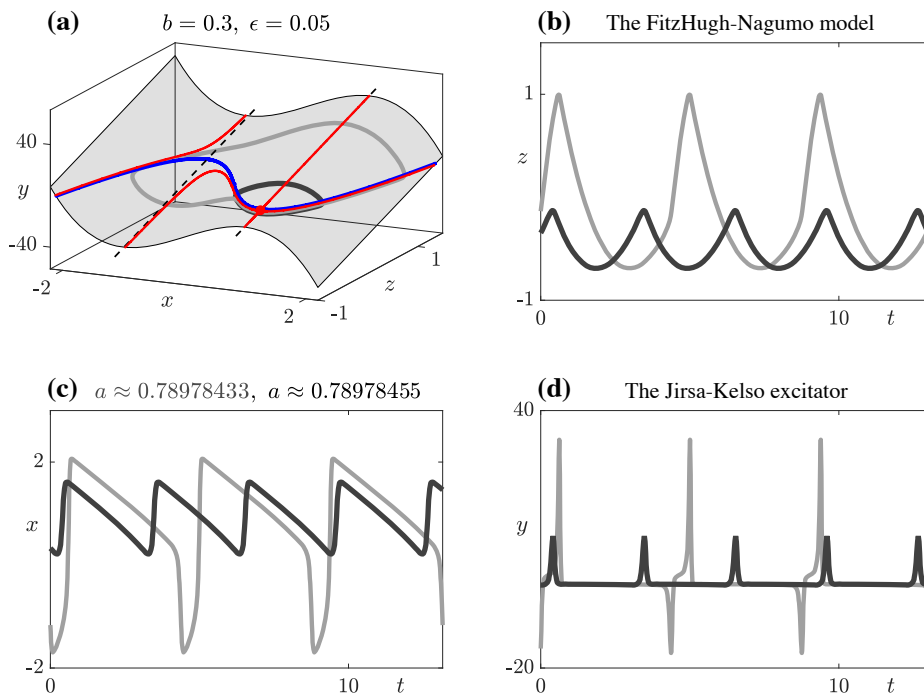


Fig. 1: Shared dynamics of the JKE and the FHN models. **a)** hypersurface H with periodic orbits and nullclines; *blue*, x -nullcline of the JKE and FHN; *red* y -nullcline of the JKE. **b),c)** and **d)** time series of the variables z, x and y , respectively. In all panels: *light grey*, big periodic orbit ($a \approx 0.78978433$); *dark grey*, small periodic orbit ($a \approx 0.78978455$); *dashed black* asymptotes of the y -nullcline; *red dot* unstable equilibrium.

Since the transformation (5) is homeomorphism (for $\epsilon > 0$) the intrinsic dynamics of the JKE (4) and FHN (6) systems are equivalent for $\epsilon > 0$ [7]. The equivalence between the FHN (6) and the JKE (4) means that the existing results related to the dynamics of the FHN model, e.g. existence of canard orbits or relaxation oscillations [14, 15], can be extended to the JKE model. However, the singularity in (5) for $\epsilon = 0$ implies that global mechanism responsible for the canard cycles and relaxation oscillations in the JKE is different than the one in the FHN; i.e. the two systems have different critical manifolds. In the rest of the paper we present a detailed study of the critical manifold C_0 of the uncoupled JKE model.

2 Critical manifold C_0 of the Jirsa-Kelso excitator model

As in the case of the FHN model the relaxation oscillations in the JKE appear through *canard explosion*, a rapid growth of periodic orbits' amplitude that happens in an exponentially small range of the control parameter [10]. However, in contrast to the FHN model, in the JKE the canard explosion is organised by a *non-generic branching point* at the y -nullcline [12], rather than by a fold point [5, 10]. Figure 2 depicts canard explosion in JKE model. Figure 2(a) shows the bifurcation diagram in the parameter a and Fig. 2(b) shows the (x, y) -phase portrait, in both panels $b = 0.3$ and $\epsilon = 0.05$.

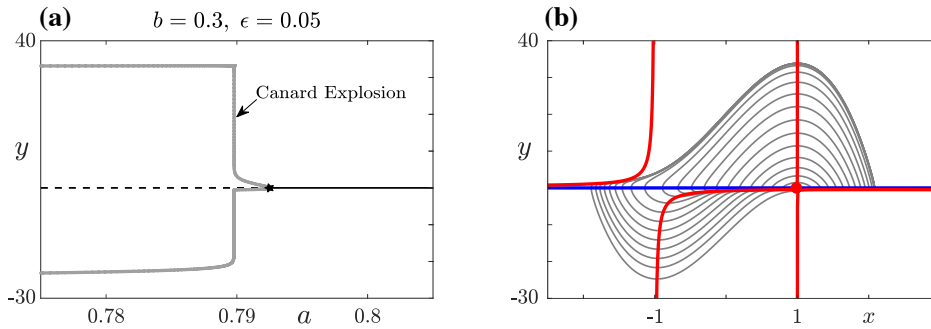


Fig. 2: Bifurcation diagram and corresponding phase portrait with periodic orbits for $a \in (a_H, 0.795)$ of the JKE model. **a)** *black solid* indicates stable branch of equilibria; *black dashed* indicates unstable branch of equilibria; *black star* indicates the Hopf point; *grey* indicates the max and min values of the variable y along the periodic orbits. **b)** *thin grey* periodic orbits for $a \in (a_H, 0.795)$; *blue* x -nullcline and *red* y -nullcline of the JKE model (4).

The periodic solutions appear at a Hopf bifurcation for,

$$a_H = \sqrt{1 - \epsilon b} \left(1 - \frac{2}{3}b - \epsilon \frac{b^2}{3} \right),$$

for which the x -coordinate is given as $x_H = \sqrt{1 - \epsilon b}$; the methodology of computations of a_H and x_H can be found for example in [1]. The bifurcation diagram and the phase portrait in Fig. 2 have been computed using AUTO continuation software [2].

We, now, focus on analysing the dynamics of the JKE system (4) in the singular limit $\epsilon = 0$, to this end we follow [1, 10]. Setting $\epsilon = 0$ in the fast system (4) defines the *layer problem* of the JKE system:

$$\begin{aligned} x' &= 0, \\ y' &= \left(-\frac{b}{3}x^3 + (b-1)x + a \right) - (x^2 - 1). \end{aligned} \quad (7)$$

Since $x' = 0$, we consider (7) as a one dimensional system with flow in the y -direction parametrised by the values of x . In other words, each value of x defines flow (7) along 1D invariant manifold (layer) in the xy -plane. The *critical manifold* of the JKE is given by the set of equilibria of the layer problem (7),

$$C_0 = \left\{ (x, y) \in \mathbb{R}^2 : y = \frac{\left(-\frac{b}{3}x^3 + (b-1)x + a \right)}{(x^2 - 1)} \right\}. \quad (8)$$

The curve C_0 has three branches separated by two vertical asymptotes at $x = \pm 1$. Note that, the critical manifold C_0 is the $\epsilon \rightarrow 0$ limit of the y -nullcline of JKE system (4), which has asymptotes at $x = \mp \sqrt{1 - \epsilon b}$. The stability of the equilibria that form C_0 can be determined from the sign of the eigenvalue $\lambda = 1 - x^2$ of the layer problem (7). For $x \in (-\infty, -1) \cup (1, \infty)$ the equilibria are stable, and for $x \in (-1, 1)$ they are unstable, for $x = \pm 1$ $\lambda = 0$ and C_0 is non-hyperbolic. Similarly to the FHN model the fast subsystem of (7) is tangent to the to the C_0 at $x = \pm 1$, however, in the case of JKE the tangency occurs at the infinity.

At the singular Hopf bifurcation, $a_{H_0} = a_H(\epsilon = 0) = 1 - 2b/3$, the middle and right parts of the critical manifold connect via a non-generic branching point [12] at $(x, y) = (1, -1/2)$ and the critical manifold takes the form:

$$C_0 = \left\{ (x, y) \in \mathbb{R}^2 : y = -\frac{bx}{3} - \frac{a_{H_0}}{(x+1)} \quad \text{or} \quad x = 1 \right\}. \quad (9)$$

Note that, for $x = \pm 1$ the stationary solutions of (7) are given as $y(t; x) = (\mp(1 - 2b/3) + a)t + \text{const.}$, meaning that, for $a = a_{H_0}$, there is no vertical flow on the $x = 1$ layer.

To analyse the flow on the critical manifold C_0 , we rescale (4) by changing the fast time τ to the slow time $t_s = \epsilon\tau$,

$$\begin{aligned} \dot{x} &= y, \\ \epsilon \dot{y} &= -\frac{b}{3}x^3 + (b-1)x + a - (x^2 - 1 + \epsilon b)y. \end{aligned} \quad (10)$$

In this way we obtain the so-called *slow* or *reduced system*: $\dot{x} = f(x, y, \epsilon), \epsilon \dot{y} = g(x, y, \epsilon)$, where $(\dot{\cdot}) = \frac{d}{dt_s}$ denotes differentiation with respect to the slow time t_s . Setting $\epsilon = 0$ and eliminating y from the differential equation (10), results in a one dimensional system

$$\dot{x} = C_0. \quad (11)$$

For $0 < b < 1$ and $-1 < a < 1$, equation (10) has only one equilibrium, which corresponds to the equilibrium of the JKE. The equilibrium is stable for $a > a_{H_0}$ and unstable for $a < a_{H_0}$; $b \in (0, 1)$. The layer problem and flow on the critical manifold are illustrated in Fig. 3. For $a = a_{H_0}$, at the singular Hopf bifurcation, there is no flow along, $x = 1$, branch of the critical manifold. On the other branch of the critical manifold, the flow is from $x = \infty$ to $x = -1$. The two branches intersect at the non-generic branching point, $C_0(x = 1) = -1/2$. The flow through this point can be studied formally by blowing up the y -axis to the cylinder $\bar{x}^2 + \bar{\epsilon}^2 = 1$ using cylindrical blow-up transformation [9].

3 Global flow of the Jirsa-Keslo excitator model

In this section, we investigate the global behaviour of the layer and reduced problems of the JKE model (4). To this end, we represent the JKE (4) on the Poincaré sphere using the following transformation of variables,

$$X = \frac{x}{\sqrt{1+x^2+y^2}}, \quad Y = \frac{y}{\sqrt{1+x^2+y^2}}, \quad Z = \frac{1}{\sqrt{1+x^2+y^2}}. \quad (12)$$

Transformation (12) defines a one-to-one correspondence between the points (x, y) in the plane and the points (X, Y, Z) on one of the hemisphere (we consider $Z \geq 0$); it follows that $x = X/Z$ and $y = Y/Z$. The advantage of using the Poincaré sphere is that the critical points at infinity are spread out along its equator, $Z = 0, X^2 + Y^2 = 1$.

Following [13] (Theorem 1 in Chapter 3.10), the equilibria and flow of the projected system on the equator, $Z = 0$, can be determined using following equation,

$$G(X, Y) = X Q_*(X, Y, 0) - Y P_*(X, Y, 0) = -X^3 \left(\frac{b}{3}X + Y \right), \quad (13)$$

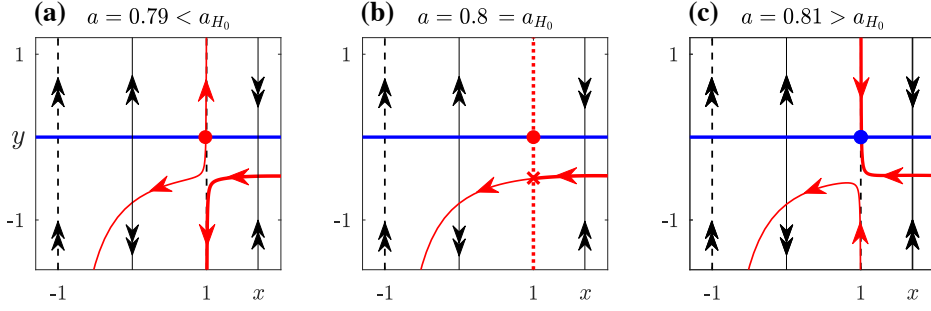


Fig. 3: The layer problem and flow on the critical manifold of the JKE model; $\epsilon = 0$, $b = 0.3$ and: **a)** $a = 0.81$, **b)** $a = a_{H_0} = 0.8$ and **c)** $a = 0.79$. *Red* the critical manifold C_0 (*thick* attractive branch, *thin* repelling branch and *dotted* zero vertical flow); *blue* the x -nullcline; *black* the layers of the fast subsystem; *single arrow* indicates slow flow and *double arrow* indicates the fast flow; *red dot* unstable node; *blue dot* stable node. Critical manifold has asymptotes at $x = \pm 1$. In **b)** the two branches of the critical manifold intersect at the non-generic branching point (*red cross*) at $(x, y) = (1, -1/2)$.

where,

$$\begin{aligned} P_*(X, Y, Z) &= Z^3 P(X/Z, Y/Z) = \epsilon Y Z^2 \\ Q_*(X, Y, Z) &= Z^3 Q(X/Z, Y/Z) \\ &= -X^2 \left(\frac{b}{3} X + Y \right) + Z^2 [(b-1)X + (1-\epsilon b)Y] + aZ^3, \end{aligned} \quad (14)$$

and,

$$\begin{aligned} P(x, y) &= \epsilon y, \\ Q(x, y) &= \left(-\frac{b}{3} x^3 + (b-1)x + a \right) - (x^2 - 1 + \epsilon b) y, \end{aligned} \quad (15)$$

is the JKE system (4); see [13] for details of derivation.

Solving Eq. (13) on the equator, $G(X, Y) = 0$, $X^2 + Y^2 = 1$, yields the four critical points

$$\begin{aligned} X &= 0, & Y &= \pm 1; \\ X &= \pm \frac{3}{\sqrt{9+b^2}}, & Y &= \mp \frac{b}{\sqrt{9+b^2}}, \end{aligned}$$

where the flow between these critical points is determined by the sign of (13). For $a < a_{H_0}$ the two points $X = 0$ are unstable nodes, and the two for $X \neq 0$ are saddles. In the singular case for $a = a_{H_0}$ and $\epsilon = 0$ the two points at $X = 0$ change from unstable nodes to non-hyperbolic nodes; the two nodes $(0, \pm 1)$ provide return mechanism for the unbounded segments of the singular canards. The flow of the JKE system on the Poincaré hemisphere for $Z \geq 0$ (projected orthogonally on (X, Y) -plane) is shown in Fig. 4.

To understand the return mechanism at the points $(0, \pm 1, 0)$ for $\epsilon = 0$, we project the flow on a plane tangent to the sphere at these points $(X, \pm 1, Z)$, and then we blow-up the non-hyperbolic nodes. We first observe that in the tangent plane, the flow in the neighbourhood of the non-hyperbolic nodes $(0, \pm 1, 0)$ is radial. Specifically, from

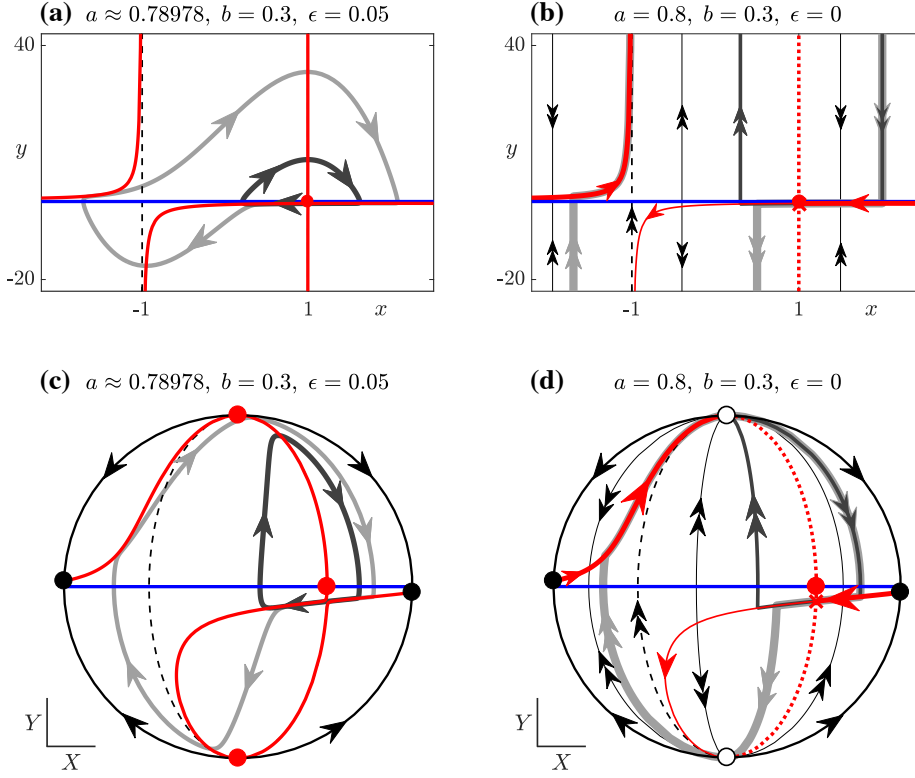


Fig. 4: The global phase portraits of the JKE system (4). **a)** phase space for $\epsilon > 0$; **b)** phase flow of the layer and reduced problems for $\epsilon = 0$. **c)** and **d)** projections of the phase space shown in **a**, and **b**, respectively, onto the upper Poincaré hemisphere ($Z \geq 0$). *Red* indicates y -nullcline in **a**, and **c** and the critical manifold C_0 in **b**, and **d** (*thick* attracting branch, *thin* repelling branch and *dotted* for zero vertical flow); *blue* the x -nullcline; *light gray* big and *dark gray* small canard orbits for $\epsilon > 0$ or a singular canard orbits for $\epsilon = 0$; *thin black* layers of the layer problem; *single arrow* slow flow on the critical manifold C_0 and *double arrow* fast flow; *black dashed* asymptotes of the critical manifold C_0 ; *thick black* the equator of the Poincaré sphere; *red dot* unstable node; *black dot* saddle; *black and white dot* non-hyperbolic node; *red cross* non-generic branching point.

[13] (Theorem 2 in chapter 3.10) we can deduce that in the neighbourhood of $(0, \pm 1, 0)$ the behaviour of the JKE model is equivalent to the behaviour of a system given by,

$$\begin{aligned} \pm X' &= -XQ_*(X, \pm 1, Z) + P_*(X, \pm 1, Z), \\ \pm Z' &= -ZQ_*(X, \pm 1, Z) \end{aligned} \quad (16)$$

with respective sign determined by the flow on the equator. Since $\epsilon = 0$ $P_* = 0$, hence the system (16) is a separable equation $dZ/dX = Z/X$ with solution $Z = CX$ for $C \in \mathbb{R}$, meaning that in the neighbourhood of $(0, \pm 1, 0)$ the flow has only radial direction.

To further investigate the dynamics near the projected nodes $(0, \pm 1, 0)$, and to understand how they can be attracting and repelling at the same time we use polar

blow-up transformation [11].

$$\Phi : (X, Z) \rightarrow (r \cos(\theta), r \sin(\theta)) \quad (17)$$

for $(\theta, r) \in [0, 2\pi) \times [0, r_0]$ for $r_0 > 0$. Transformation (17) replaces the non-hyperbolic equilibrium $(X, Z) = (0, 0)$ with a unit circle $\mathcal{S}^1 \times \{r = 0\}$.

Changing X', Z' , (16), to radial variables and solving for r' and θ' gives the set of the differential equations of the blown-up system,

$$\begin{aligned} r' &= \cos(\theta)X' + \sin(\theta)Z' \\ r\theta' &= -\sin(\theta)X' + \cos(\theta)Z' \end{aligned} \quad (18)$$

Using (16), with $\epsilon = 0$, we obtain,

$$\begin{aligned} r' &= \cos(\theta) [-XQ_*(X, 1, Z)] + \sin(\theta) [-ZQ_*(X, 1, Z)], \\ r\theta' &= -\sin(\theta) [-XQ_*(X, 1, Z)] + \cos(\theta) [-ZQ_*(X, 1, Z)], \end{aligned} \quad (19)$$

which simplifies to,

$$\begin{aligned} r' &= \sigma(\theta)r^3 - \mu(\theta)r^4, \\ \theta' &= 0, \end{aligned} \quad (20)$$

with,

$$\begin{aligned} \sigma(\theta) &= \cos^2(\theta) - \sin^2(\theta), \\ \mu(\theta) &= -\frac{b}{3} \cos^3(\theta) + (b-1) \cos(\theta) \sin^2(\theta) + a_{H_0} \sin^3(\theta). \end{aligned}$$

Finally, by rescaling time $t_s \mapsto r^2 t_s$ we obtain a set of equations that describes the dynamics of the layer problem in the neighbourhood of the $(X, Z) = (0, 0)$,

$$\begin{aligned} r' &= r(\sigma(\theta) - \mu(\theta)r), \\ \theta' &= 0. \end{aligned} \quad (21)$$

Figure 5 illustrates the extended critical manifold near infinity. Figure 5(a)-(b) shows qualitative phase portraits of the layer problem of the JKE model projected on the XZ -plane near the equilibrium $(X, Y, Z) = (0, 1, 0)$ and Fig. 5(c)-(d) illustrates the dynamic near the blown-up circle in $\bar{X}\bar{Z}$ -plane.

The layer problem at infinity (21) is parametrised by θ , and has three branches of equilibria: $B_1 = \{(\theta, r) : \theta \in [0, \pi), r = 0\}$, $B_2 = \{(\theta, r) : \theta = \pi/4, r \geq 0\}$ and $B_3 = \{(\theta, r) : \theta \in [3\pi/4, \pi), r = \sigma(\theta)/\mu(\theta)\}$, given by the solutions of $\sigma(\theta)/\mu(\theta) \geq 0$; see Fig. 5(c). The branch B_1 is the blown-up node at infinity (black and white circle in Fig. 5(a)) and can be interpreted as an additional part of the critical manifold C_0 . Branch B_2 corresponds to the line of no vertical flow (dotted line in Figs. 3-5) and branch B_3 corresponds to the projected critical manifold C_0 . Branch B_1 intersects with branch B_2 at $\theta = \pi/4$ (red cross in Fig. 5(c)) and with branch B_3 at $\theta = 3\pi/4$ (red square in Fig. 5(c)).

Stability of the equilibria along the branch B_1 is determined by the eigenvalue $\lambda = \sigma(\theta)$, they are radially attracting for $\theta \in (\pi/4, 3\pi/4)$, and radially repelling for $\theta \in (0, \pi/4) \cup (3\pi/4, \pi)$. Stability of the equilibria along the branch B_3 is determined by the eigenvalue $\lambda = \sigma(\theta) - 2\mu(\theta)(\sigma(\theta)/\mu(\theta)) = -\sigma(\theta)$, they are radially attracting for $\theta \in (3\pi/4, \pi)$. There is no radial flow along the B_2 branch.

Branch B_1 changes stability at the points of intersection with branches B_2 and B_3 . Orthogonal intersection with the branch B_2 , $(r, \theta) = (0, \pi/4)$ (red cross in Fig. 5(c)), is a non-generic branching point that is of the same nature as the point

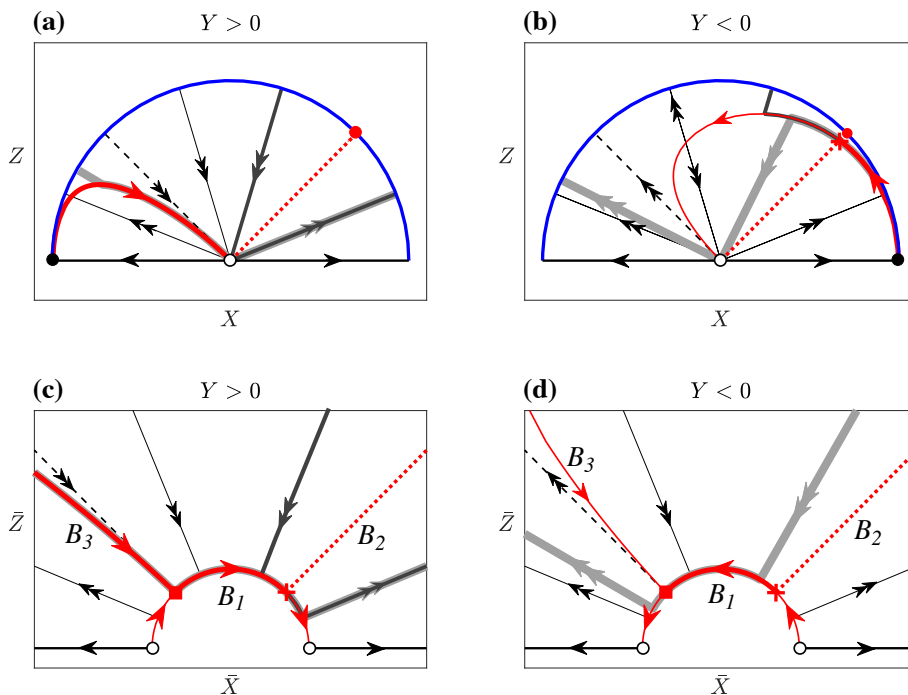


Fig. 5: The projection and blowup of dynamics of the JKE system near the equilibrium at infinity for $a = a_{H_0}$, $b = 0.3$ and $\epsilon = 0$. **a)** and **b)** the flow projected on $(X, 1, Z)$ and $(X, -1, Z)$, respectively. **c)** and **d)** dynamics near the blown-up circle for $(0, 1, 0)$ and $(0, -1, 0)$, respectively. Color coding as in Fig. 4.

$(x, y) = (1, -1/2)$. Branches B_1 and B_3 intersect and exchange stability in a transcritical bifurcation at $(r, \theta) = (0, 3\pi/4)$ (red square in Fig. 5(c)). The intersection of branches B_1 and B_2 can be analysed formally by means of the cylindrical blowup transformation [8], while the intersection of branches B_1 and B_3 can be analysed formally by the classical techniques described in [9].

To complete the analysis of the critical manifold, we determine the direction of the flow of the reduced problem along the branch B_1 . To this end we again use Theorem 2 from chapter 3.10 of [13] and blowup method. Since, for the singular case $\epsilon = 0$, the flow is defined only on the critical manifold, equations (15) for the reduced problem of the JKE model simplify to $P(x, y) = y$ and $Q(x, y) = 0$. Hence, the system (16) for the reduced problem is given by,

$$\begin{aligned} \dot{X} &= P_*(X, 1, Z) = Z^3 P(X/Z, 1/Z) = Z^2, \\ \dot{Z} &= 0. \end{aligned} \quad (22)$$

After polar blow-up Φ (17), and rescaling time $t \mapsto rt$, the flow along the branch B_1 can be determined from the system,

$$\begin{aligned} \dot{r} &= \cos(\theta) \sin^2(\theta) r, \\ \dot{\theta} &= -\sin^3(\theta). \end{aligned} \quad (23)$$

Therefore, the reduced problem at infinity has two non-hyperbolic nodes on the blown-up circle, the branch B_1 , $(r, \theta) = (0, 0)$ and $(r, \theta) = (0, \pi)$ and the flow between the two nodes is clockwise.

The behaviour near the node $(0, -1, 0)$ can be analysed in the same way as the node $(0, 1, 0)$, compare Fig. 5(b) and (d). After projection on the tangent plane and polar blow-up. The layer problem is given as,

$$\begin{aligned} r' &= r(\sigma(\theta) + \mu(\theta)r), \\ \theta' &= 0, \end{aligned} \tag{24}$$

while the dynamics on the blown-up circle is described by,

$$\begin{aligned} \dot{r} &= -\cos(\theta) \sin^2(\theta)r, \\ \dot{\theta} &= \sin^3(\theta). \end{aligned} \tag{25}$$

The layer problem at infinity (24) has again three branches of equilibria.

$B_1 = \{(\theta, r) : \theta \in [0, \pi), r = 0\}$, $B_2 = \{(\theta, r) : \theta = \pi/4, r \geq 0\}$ and $B_3 = \{(\theta, r) : \theta \in [0, 3\pi/4), r = -\sigma(\theta)/\mu(\theta)\}$, given by solutions of $-\sigma(\theta)/\mu(\theta) \geq 0$, see Fig. 5(d). Branches B_1 and B_2 , have the same stability and interpretation as the equivalent branches near the $(0, 1, 0)$ node. Branch B_3 again corresponds to the projected critical manifold and has the same eigenvalue $\lambda = -\sigma(\theta)$, but near the node $(0, -1, 0)$ it is defined over $\theta \in [0, 3\pi/4]$ rather than $\theta \in [3\pi/4, \pi)$ and hence is radially repelling. Branches B_1 and B_3 exchange stability at point $(r, \theta) = (0, 3\pi/4)$ in a transcritical bifurcation (red square in Fig. 5(d)). The reduced problem at infinity (25) has two non-hyperbolic nodes on the blown-up circle, $(r, \theta) = (0, 0)$ and $(r, \theta) = (0, \pi)$. The flow of the slow subsystem between the two nodes is counterclockwise.

4 Conclusion

In the paper we presented a detailed analysis of the mechanism that leads to appearance of relaxation oscillations in the Jirsa-Kelso excitator. We, first, demonstrated existence of canards in the the Jirsa-Kelso excitator and showed that they are organised by a non-generic branching point at the y -nullcline of the system. We, then, combined projection onto the the Poincaré sphere, the slow-fast analysis and blowup method to investigate the return mechanism of the periodic orbits in the singular case, $\epsilon = 0$. The analysis revealed that the return mechanism is organised by two non-hyperbolic nodes at infinity, for which hyperbolicity can be recovered with a polar blowup transformation. The blown-up flow at infinity showed the the singular canard cycles can indeed be represented as concatenation of orbits, respecting the direction of time, of the layer and reduced problems. Overall, presented analysis shows that although the Jirsa-Kelso excitator can be transformed into the FitzHugh-Nagumo model via a homeomorphic transformation (for $\epsilon > 0$), the two models have different critical manifolds (for $\epsilon = 0$) and different mechanisms of the global transitions.

From the point of view applications, presented analysis shows that the slow-fast nature of the transition between discrete and rhythmic movements described by the JKE can be analysed and understood, not only in the abstract phase space of the FHN, but also in the experimentally relevant phase-space of position and velocity. Our study, further, describes an additional scenario, besides changes in parameters a and b , for the transition between discrete and rhythmic movements. In this scenario, the transition is controlled by the time scale parameter ϵ . For small values of $\epsilon \ll 1$ the relaxation oscillations have a very long period and can be viewed as two almost steady states connected by fast transitions between them; note that the results of

stimulus applied to such oscillations would be very similar to the perturbations of the bi-stable regime of the JKE, for frequency of the stimulus higher than the frequency of the oscillations. As the ϵ increases the oscillations become more sinusoidal and can be viewed as a rhythmic movement. Transformation between equations (2) and (4) suggest that experimentally the time scale parameter ϵ could be related to changes in range of admitted velocities v . Further empirical studies would be required to validate these observations in relevant movement tasks.

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Authors' contributions

PS and SAR contributed equally to the paper. SAR performed the basic analysis and computations, PS refined the analysis and figures. SAR and PS wrote the paper. KTA supervised the work and revised the final manuscript.

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