Rotating Eights: I. The three $\Gamma_i$ families

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Abstract

We show that three families of relative periodic solutions bifurcate out of the Eight solution of the equal-mass three-body problem: the planar Hénon family, the spatial Marchal $P_{12}$ family and a new spatial family. The Eight, considered as a spatial curve, is invariant under the action of the 24-element group $D_6 \times Z_2$. The three families correspond to symmetry breakings where the invariance group becomes isomorphic to $D_6$, the three $D_6$s being embedded in the larger group in different ways. The proof of the existence of these three families relies on writing down the action integral in a rotating frame, viewing the angular velocity of the frame as a parameter, exploiting the invariance of the action under a group action which acts on the angular velocities as well as the curves and, finally, checking numerically the non-degeneracy of the Eight. Pictures and numerical evidence of the three families are presented at the end.

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1. Results: three families of rotating Eights

We consider the equal-mass Newtonian three-body problem in three-space:

$$\frac{d^2 x_i}{dt^2} = \frac{\partial U}{\partial x_i}(x_1, x_2, x_3), \quad U(x_1, x_2, x_3) = \sum_{i<j} \frac{1}{\|x_i - x_j\|}, \quad x_i \in \mathbb{R}^3.$$

The Eight (figure 1) is a recently discovered periodic planar solution to this problem (see [Mo, CM]). Due to the homogeneity of the potential function $U$, if $x(t)$ is a solution then so is $\mu^{-2/3} x(\mu t)$ for every $\mu > 0$; in particular to any periodic solution there corresponds a rescaled periodic solution with a given period $T$. Hence from now on, we shall stick to a given period $T$. 
The Eight has a large symmetry group, $\Gamma$ (see section 4.3). It is *choreographic*, meaning that the three masses of the Eight travel along a single curve in space. This space curve is a planar figure Eight, with three obvious symmetry axes, two in the plane of the Eight and the third orthogonal to that plane.

We can rotate (in space) any solution to the three-body problem to obtain a new solution. We can also translate time. As a consequence of these symmetries, the linearization of the return map ($=\text{time map of the flow}$) associated with a periodic solution of period $T$ will have 1 as an eigenvalue. Provided the solution is not a relative equilibrium (orbit of the rotation group), then the multiplicity of the eigenvalue will be at least $4:3$ for the rotation group and 1 for time translation. (We fix the centre of mass to be zero and so translations are not allowed.) We shall call that solution *Poincaré non-degenerate* if the dimension of the eigenspace associated with the eigenvalue 1 is exactly 4. Because the flow of the $n$-body problem is Hamiltonian and the angular momentum of the Eight is equal to 0, the algebraic multiplicity of the eigenvalue 1 is at least 8 (and in fact equal to 8, see section 4), but we shall show that the Eight is Poincaré non-degenerate by checking numerically its Jordan block structure. This non-degeneracy is consistent with the fact, proved in [KZ] using interval arithmetic, that, among $T$-periodic curves, the Eight is isolated modulo rotations, translations and scaling. It completes the numerical proof of stability of the Eight after reduction given in [S].

A solution is called *relatively $T$-periodic* if it is periodic up to a rigid rotation of the three bodies or, equivalently, if there is a rotation $R$ such that $x(T) = Rx(0)$ and $\dot{x}(T) = R\dot{x}(0)$ (where $R$ acts diagonally on the positions and velocities of the bodies). Upon writing $R = \exp(\omega T)$ for an angular velocity $\omega$, the solution becomes $T$-periodic when viewed with respect to a rotating frame which is rotating with constant angular velocity $\omega$.

The main goal of this paper is to prove the following result.

**Theorem 1 (Existence of the $\Gamma_i$ families).** Under the hypothesis, numerically verified, that the Eight is Poincaré non-degenerate, there are three families $\{x^i_\lambda\}$, $i = 1, 2, 3$, $\lambda \in \mathbb{R}$, $|\lambda| < \lambda_0$, of relatively $T$-periodic solutions bifurcating out of the Eight, one family for each symmetry axis $e_1$, $e_2$, $e_3$ of the Eight. Viewed in the frame rotating around the axis $e_i$ with angular speed $\lambda$, the solution $x^i_\lambda$ is $T$-periodic, depends analytically on $\lambda$, and has symmetry group $\Gamma_i$, described in section 4.3.

Two of these families were known. One was discovered by Michel Hénon [H2] using the numerical methods of [H1], another by Christian Marchal [Ma]. The third family is new. Pictures of the three families in the rotating frame, obtained numerically, are given at the end of this paper.

About the unicity of the $\Gamma_i$ families. After seeing a former version of this paper, David Chillingworth showed us how it would be possible to check using [Ch] that no other families
Outline of proof. Any relatively periodic solution is periodic in a frame which rotates with an appropriate angular velocity $\omega$. Treat $\omega$ as a bifurcation parameter. The problem becomes a problem in the theory of bifurcations with symmetry. There is a parametrized action principle $A_\omega$ whose non-collision critical points are solutions periodic in the frame rotating with constant angular velocity $\omega$. The Eight is a critical point for $A_0$. Applying the symmetries of rotation and time translation to the Eight yields a four-dimensional critical submanifold $E$ of Eights for $A_0$. The idea is to turn on $\omega$ and simultaneously move normally to $E$ so as to look for critical points to $A_\omega$ in the space normal to $E$ at a given Eight. This is the Lyapunov–Schmidt reduction (section 4.1) and it reduces the search for nearby relatively periodic solutions to the search for critical points of a certain function on $E$. Using group theory, in section 4.5 we reduce the latter critical point problem to a nonlinear eigenvalue problem. In section 5 we solve this nonlinear eigenvalue problem explicitly (again using group theory) under the assumption that $\omega$ is parallel to any one of the three symmetry axes of the Eight (section 5). Hence the three families. The Lyapunov–Schmidt method requires the non-degeneracy of the Hessian of $A_0$ in directions orthogonal to $E$. We show in lemma 2 of section 4.1 that this non-degeneracy is equivalent to the non-degeneracy of the Poincaré map associated with the Eight. We verify this latter non-degeneracy numerically in section 4.4.

2. The Lagrangian and the action in a rotating frame

The configuration space $X$ for the equal-mass spatial three-body problem consists of triples $x = (x_1, x_2, x_3)$ of vectors $x_i \in \mathbb{R}^3$ such that $x_1 + x_2 + x_3 = 0$. If $R : \mathbb{R}^3 \to \mathbb{R}^3$ is an orthogonal transformation then we write $Rx$ for $(Rx_1, Rx_2, Rx_3)$. To formulate the dynamics in a rotating frame, we take $R = R(t)$, a family of rotations depending on time, and use as variable $\xi(t)$ where

$$x(t) = R(t)\xi(t).$$

We are only interested in frames rotating with constant angular velocity $\omega$, in which case

$$R(t) = \exp(t\omega).$$

Some words are in order regarding the notation $\exp(t\omega)$. Throughout this paper we will be using the standard identifications of $\mathbb{R}^3$ with the Lie algebra $so(3)$ of skew-symmetric linear maps of $\mathbb{R}^3$ and with the spaces $\Lambda^2\mathbb{R}^3$ of exterior 2-forms and $\Lambda^2\mathbb{R}^3$ of 2-vectors. Under the first identification, $\omega \in \mathbb{R}^3$ defines the skew-symmetric linear operator $v \mapsto \omega \times v$. Under the second identification, $\omega$ defines the 2-form $(v, w) \mapsto \langle \omega, v \times w \rangle$. Under these identifications, if $(e_1, e_2, e_3)$ is an oriented basis for $\mathbb{R}^3$ then $e_3$ represents the rotation of angle $+\pi/2$ in the $(e_1, e_2)$ plane which in turn is identified with the bivector $e_1 \wedge e_2$. It is fundamental in the following to keep in mind that $\omega$ is a bivector, and so transforms as a bivector, not a vector.

The Lagrangian $L_\omega$ governing the dynamics in the moving frame is

$$L_\omega(\xi, \dot{\xi}) = \frac{1}{2}\|\omega \times \xi + \dot{\xi}\|^2 + U(\xi) = L_0(\xi, \dot{\xi}) + \frac{1}{2}\|\omega \times \xi\|^2 + \langle \omega \times \xi, \dot{\xi} \rangle.$$ 

Here we have used the invariance of $U$, the negative of the potential, under rotations: $U(R\xi) = U(\xi)$. This Lagrangian can be rewritten as

$$L_\omega(\xi, \dot{\xi}) = L_0(\xi, \dot{\xi}) + \frac{1}{2}\langle \omega, I(\xi)\omega \rangle + \langle \omega, J(\xi, \dot{\xi}) \rangle,$$
where $I$ is the inertia matrix of the solid defined by the configuration $\xi$ and $J$ is the angular momentum of the phase element $(\xi, \dot{\xi})$. Note that $L_0$ is the Lagrangian for the three-body problem in an inertial frame. The action in the rotating frame is given by the formula

$$A_\omega(\xi) = A_0(\xi) + \frac{1}{2}\langle \omega, I(\xi)\omega \rangle + \langle \omega, J(\xi) \rangle,$$

where $I$ and $J$ are the integrated versions of $I$ and $J$. We take the domain of the action $A_\omega$ to be the loop space $H^1 = H^1(S^1, X)$, consisting of those maps from the time circle $S^1 = \mathbb{R}/T\mathbb{Z}$ to the configuration space $X$, which, together with their derivative, are square integrable. The choice of $H^1$ does allow for collisions. However, because $H^1$ is contained in the space $C_0$ of continuous paths, and because all our analysis is local in an $H^1$-neighbourhood of the Eight, collisions will be avoided, at least for small values of the bifurcation parameter $\omega$.

### 3. Symmetries and symmetry breakings

We introduce a group $G$ which acts naturally on the space of relatively $T$-periodic solutions when the three masses are equal. For the notations, we follow Ferrario and Terracini [FT].

#### 3.1. Action on loops

The group $G = O(3) \times O(2) \times S_3$ acts on the loop space $H^1$ in the following way. The factor $O(3)$ acts diagonally on the configuration space $\chi$. The factor $O(2)$ acts on the time circle $\mathbb{R}/T\mathbb{Z}$. The third factor, $S_3$—the symmetric group on three letters—acts on the three indices 1, 2, 3 which label the three bodies. $G$ acts on $H^1$ according to

$$(\rho, \tau, \sigma) \cdot (\xi_1, \xi_2, \xi_3)(t) = (\rho \xi_{\sigma^{-1}(1)}(\tau^{-1}t), \rho \xi_{\sigma^{-1}(2)}(\tau^{-1}t), \rho \xi_{\sigma^{-1}(3)}(\tau^{-1}t)).$$

#### 3.2. Action on bivectors

We now define an action of $G$ on the space $\Lambda^2 \mathbb{R}^3$ of bivectors $\omega$, such that

$$A_g \omega(g\xi) = A_\omega(\xi)$$

for $g = (\rho, \tau, \sigma) \in G$. The permutation group factor $\sigma \in S_3$ acts trivially on bivectors. The group $O(2)$ of time isometries acts according to the sign homomorphism: $\tau \cdot \omega = \det(\tau)\omega$. Thus any orientation-preserving isometry of the time circle acts as the identity on $\omega$ while an orientation-reversing isometry reverses the sign of $\omega$. This reversal is seen to follow either from the definition of angular velocity as a time derivative, or by its duality with angular momentum and the formula for angular momenta in terms of $\xi \wedge \dot{\xi}$. Finally the linear isometry group $O(3)$ of $\mathbb{R}^3$ acts in the standard way by push-forward. If we identify $\Lambda^2 \mathbb{R}^3$ with $\mathbb{R}^3$ as above, then $\rho \in O(3)$ sends $\omega \in \mathbb{R}^3$ to $\rho \cdot \omega = \det(\rho)\rho\omega$.

#### 3.3. Space of Eights

Let $x_8$ be the figure Eight solution. Write $\mathcal{E} = Gx_8$ for the orbit of the Eight under the $G$-action. By definition, $\mathcal{E}$ is the space of Eights. It is a four-dimensional submanifold of $H^1$ isomorphic to $G/\Gamma$ where the finite group $\Gamma \subset G$ is the isotropy group of the Eight (see section 4.3 for more on $\Gamma$). The submanifold $\mathcal{E}$ admits a tubular neighbourhood $T$ in the full loop space $H^1$ (use normal projection), which allows us to introduce coordinates $(\eta, \zeta)$ in $T$, with $\eta \in \mathcal{E}$ and $\zeta \perp T_\eta \mathcal{E}$, and $\eta + \zeta$ sweeping out the neighbourhood $T$. To stress the dependence of the action on these variables, we will write

$$A_\omega(\eta + \zeta) = A(\eta, \zeta, \omega).$$
4. The Liapunov–Schmidt reduction

4.1. Initial reduction

The Euler–Lagrange equation \( dA_0(\eta + \zeta) = 0 \) is equivalent to the two equations
\[
\frac{\partial A}{\partial \eta}(\eta, \zeta, \omega) = 0, \quad \frac{\partial A}{\partial \zeta}(\eta, \zeta, \omega) = 0.
\]
These equations are identically satisfied on the space of Eights \( \zeta = 0 \) for \( \omega = 0 \). According to the implicit function theorem, we can solve the second equation for a
\[
\zeta = \zeta(\eta, \omega)
\]
defined in a neighbourhood \( |\omega| < C \) of \( \omega = 0 \) provided that \( (\partial^2 A/\partial \zeta^2)(\eta, 0, 0) \) is invertible. Then the critical point equation reduces to the finite-dimensional condition
\[
\frac{\partial A}{\partial \eta}(\eta, \zeta(\eta, \omega), \omega) = 0
\]
equivalent to the condition that \( \eta \) be a critical point of the function on \( \mathcal{E} \) defined by \( \eta \mapsto A(\eta, \zeta(\eta, \omega), \omega) \). This last equation is the Liapunov–Schmidt reduction. Solving it in a neighbourhood of \( \omega = 0 \) is the topic of sections 4.4 and 5.

Remark. According to the implicit function theorem, the invertibility condition at \( \eta \) implies the existence of the solution \( \zeta(\eta, \omega) \) for \( \omega \) in some neighbourhood \( |\omega| < C_\eta \) of \( \omega = 0 \). The constant \( C_\eta \) can be taken independent of \( \eta \), as can be seen by the equivariance of \( A \) or, if one prefers, by the compactness of \( \mathcal{E} \).

Before we continue, care is needed in understanding the above condition of invertibility of \( \partial^2 A/\partial \zeta^2 \). Note that we set \( \omega = 0 \) in writing this condition, so that we are dealing with the action functional \( A_0 \). Write \( H^1 = H^1(S^1, \mathcal{X}) \) for the Sobolev space of maps from the circle to \( \mathcal{X} \) which are square-integrable, together with their \( j \)th derivative, \( j = 1, \ldots, s \). Use the \( L^2 \)-pairing to convert differentials and quadratic forms to vectors and linear operators. Thus for \( z \in H^2 \), the 1-form \( dA_0(z) \) (a linear functional on \( H^1 \)) becomes the \( L^2 \)-gradient \( \nabla A_0(z) = -\zeta + \nabla U(z) \) (lying in \( L^2 = H^0 \)). The map \( z \mapsto \nabla A_0(z) \) is a nonlinear map from \( H^2 \) to \( L^2 = H^0 \). And the Hessian of the action becomes
\[
d^2 A_0(z)(w_1, w_2) = (L(z)w_1, w_2),
\]
where
\[
L(z)w = -\dot{w} + \nabla^2 U(z)w
\]
and where we require that \( w \in H^2 \). (The Hessian \( d^2 A_0(\eta) \) is defined for \( w_1, w_2 \in H^1 \). However, for \( L(\eta)w_1 \) to be defined we need \( w_1 \in H^2 \).) Freeze \( \zeta \) and write \( L = L(\zeta) \) for notational simplicity. \( L \) is the linearization of \( \nabla A_0(z) \) at \( z = (\eta, 0) \). \( L \) is formally self-adjoint and elliptic, and maps \( H^2 \) to \( L^2 \). \( G \)-invariance of the action implies that \( T_\eta \mathcal{E} \subset \ker(L) \). Self-adjointness of \( L \) implies that \( \text{im}(L) \subset T_\eta \mathcal{E} \perp \). The Poincaré non-degeneracy condition is precisely that \( \ker(L) = T_\eta \mathcal{E} \) which is four-dimensional. The index of \( L \) is zero by self-adjointness or, alternatively, because \( L \) is a compact perturbation of the operator \( -(d/dt)^2 \) which has index 0. It follows that \( \text{im}(L) = T_\eta \mathcal{E} \perp \cap L^2 \), the \( \mathcal{E} \perp \) being relative to \( L^2 \). Write \( L_\zeta \) for the restriction of \( L \) to \( T_\eta \mathcal{E} \perp \). \( L_\zeta \) is what we meant by \( \partial^2 A/\partial \zeta^2 \) in the paragraph above when we were invoking the implicit function theorem. Then
\[
L_\zeta : T_\eta \mathcal{E} \perp \cap H^2 \to T_\eta \mathcal{E} \perp \subset L^2
\]
is one-to-one and onto—an isomorphism. We have proved the following lemma.

**Lemma 2.** The Poincaré non-degeneracy of the Eight \( \eta \) is equivalent to the invertibility of \( (\partial^2 A/\partial \zeta^2)(\eta, 0, 0) \).
Summary. We have established that the Poincaré non-degeneracy of the Eight implies the existence of an $\omega$-dependent section $\zeta: T \times U \subset T \times \mathbb{R}^3 \to \mathcal{E}$, written $(\eta, \omega) \mapsto \zeta(\eta, \omega)$, of the normal bundle $T \to \mathcal{E}$ such that the second equation of (1) holds. The group $G$ acts on $T$, $\mathbb{R}^3$ and $\mathcal{E}$. The $G$-invariance of the action $A_\omega$ implies that this $\omega$-dependent section map is $G$-equivariant:

$$\zeta(g\eta, g\omega) = g \zeta(\eta, \omega).$$

Remark. The invertibility of $(\partial^2 A / \partial \zeta^2)(\eta, 0, 0)$ is nothing but the Bott non-degeneracy of the critical manifold $E$ (see [B]). Hence the two notions of non-degeneracy coincide.

The following two sections are devoted to proving the Poincaré non-degeneracy of the Eight.

4.2. The Jordan structure of the monodromy operator (1)

Recall the linearization operator $L$ of section 4.1:

$$(Lw)(t) = -\ddot{w} + \nabla^2 U(z(t))w(t),$$

where $z(t)$ denotes the standard Eight. Here, we do not insist that $w$ need be periodic in $t$. The equation $Lw = 0$ is a system of six linear differential equations of second order: six because we fix the centre of mass to be zero and are considering spatial solutions. The solution space $S$ to $Lw = 0$ is a 12-dimensional linear space which identifies with the three-body phase space $\mathbb{R}^{12} = \mathbb{R}^6 \oplus \mathbb{R}^6$ by associating to a solution $w$ its initial conditions $w(0)$, $\dot{w}(0)$.

Write $\Phi_t: \mathbb{R}^{12} \to \mathbb{R}^{12}$ for the map taking $(w(0), \dot{w}(0))$ to $(w(t), \dot{w}(t))$. This map is the linearization of the Newton flow $\phi_t$. The monodromy operator is

$$M = \Phi_T,$$

where $T$ is the Eight’s period. Now, return to the $L$ of section 4.1, whose domain consisted of vector-valued periodic functions: $w(t + T) = w(t)$. A function $w \in S$ is in $\ker(L)$ if and only if $w(T) = w(0)$ and $\dot{w}(T) = \dot{w}(0)$, i.e. if and only if $MZ = Z$ where $Z = (w(0), \dot{w}(0))$.

That is to say

$$\ker(L) = \{w \in S : Mw = w\}. \quad (*)$$

It follows that the Jordan block structure of the symplectic matrix $M$ is the key to non-degeneracy.

For each first integral $F$ of the flow $\phi_t$, the value $X = X_F(a)$ of its Hamiltonian vector field $X_F$ at a point $a$ of the Eight is an eigenvector of $M$ with eigenvalue 1 (i.e. $MX = X$). In the same way, $\ell = dF(a)$ is an eigencovector of $M$ with eigenvalue 1 (i.e. $\ell \circ M = \ell$).

In terms of matrix representation, $X$ corresponds to a column vector while $\ell$ corresponds to a row vector (or left eigenvector) (see for instance [Me, p 80, lemma 6.5.1]). Hence, with the four integrals of energy $H$ and angular momenta $J_1, J_2, J_3$ are associated four eigenvectors with eigenvalue 1. The following lemma and the fact that the angular momentum of the Eight is equal to 0 imply that the multiplicity of the eigenvalue 1 actually is at least 8.

Lemma 3. Let $M : V \to V$ be a linear operator of a complex vector space of finite dimension $n$. Let $X_1, \ldots, X_p$ (respectively $\ell_1, \ldots, \ell_q$) be $p$ independent eigenvectors (respectively $q$ independent eigencovectors) with eigenvalue 1. If moreover $\ell_i(X_j) = 0$ for all pairs $(i, j)$, $1 \leq i \leq q$, $1 \leq j \leq p$, the multiplicity of 1 as an eigenvalue of $M$ is at least $p + q$.

Proof. In a basis $\{e_1, \ldots, e_n\}$ of $V$ such that

1. $e_i = X_i$ for $1 \leq i \leq p$. 
2. $e_i = \ell_i$ for $1 \leq i \leq q$. 

Then $M e_i = e_i$ for $1 \leq i \leq p$ and $M e_i = \ell_i$ for $1 \leq i \leq q$. Therefore, $M e_i = e_i$ for $1 \leq i \leq p + q$, and the multiplicity of 1 as an eigenvalue of $M$ is at least $p + q$. 


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(2) $e_1, \ldots, e_{n-q}$ generate the kernel of $(\ell_1, \ldots, \ell_q) : \mathbb{R}^n \to \mathbb{R}^q$.

(3) $\ell_i(e_{n-q+j}) = \delta_{ij}$ for all $1 \leq i, j \leq q$.

the matrix of $M$ has the following form (the diagonal blocks are respectively of size $p, n - p - q, q$):

$$
\begin{pmatrix}
\text{Id} & * & * \\
0 & * & *
\end{pmatrix},
$$

which proves the lemma.

Now, two complex conjugate pairs $\exp(\pm i2\pi \lambda_1), \exp(\pm i2\pi \lambda_2)$ lying on the unit circle were already numerically computed by Simó in his study of the stability of the Eight after reduction of the first integrals [S]. (Simó studies the planar problem, but when the angular momentum is zero the motion is necessarily planar.) The angles $\lambda_j$ are approximately

0.298 092 529 004 and 0.008 422 724 708.

From this fact and lemma 3 (with $p = q = 4$), we deduce the following lemma.

**Lemma 4.** The spectrum of the monodromy operator $M$ consists of the real number 1 with multiplicity 8 and the two simple complex conjugate pairs $e^{\pm i2\pi \lambda_1}, e^{\pm i2\pi \lambda_2}$, each with multiplicity 1. Moreover, the Jordan blocks corresponding to the eigenvalue 1 are at most $2 \times 2$.

**Proof.** As the angular momentum of the Eight is equal to 0, the four first integrals $H, J_1, J_2, J_3$ commute and hence the eigenvectors and covectors deduced from them satisfy the hypotheses of lemma 3. It remains for us only to prove the assertion on the size of the Jordan blocks of $M$.

But the fact that the remaining eigenvalues are all different from 1 implies that we can find a new basis of $V$ which converts $M$ into a matrix of the form

$$
\begin{pmatrix}
\text{Id} & 0 & * \\
0 & 0 & * \\
0 & 0 & \text{Id}
\end{pmatrix}.
$$

It is now clear that the restriction $\text{Id} + N$ of $M$ to the eight-dimensional generalized eigenspace associated with the eigenvalue 1 is block triangular and hence such that $N^2 = 0$, which proves the assertion.

4.3. The Jordan structure of the monodromy operator (2)

In this section, which is not necessary for the following, we elaborate on the Jordanization of $M$, paying attention successively to the symplectic structure, the ‘energy-scaling’ block and the ‘horizontal-vertical’ decomposition.

4.3.1. Non-diagonal terms. In our symplectic setting, we can give a simple interpretation of the non-diagonal terms in the Jordan blocks associated with the eigenvalue 1 of the monodromy operator $M$.

The vector space $V$ is of dimension 12 and it splits as the sum $V = U \oplus E$ of two invariant symplectic subspaces, where the $U$ is the eight-dimensional generalized eigenspace associated with the eigenvalue 1. We shall still call $M : U \to U$ the restriction to $U$ of the monodromy map. The subspace $L$ of $U$ generated by the four eigenvectors $X_H(a), X_{J_1}(a), X_{J_2}(a), X_{J_3}(a)$ is Lagrangian, precisely because the angular momentum of the Eight equals 0. The restriction of $M$ to $L$ is the identity, and hence the following lemma applies.
Lemma 5. Let $M : U \to U$ be a symplectic linear mapping from the symplectic vector space $(U, \omega)$ of dimension $2n$ to itself. Let $L \subset U$ be a Lagrangian subspace on which $M$ induces the identity. Then
(1) the sole eigenvalue of $M$ is 1;
(2) the formula
$$s(u', u'') = \omega(Mu' - u', u'')$$
defines a symmetric bilinear form on the quotient space $U/L$. This bilinear form is non-degenerate if and only if the Jordan form of $M$ is made of $n$ non-trivial $2 \times 2$ blocks. More precisely, the signature of $s$ gives the number of blocks with positive (respectively negative) non-diagonal entry.

Proof. The fact that $s$ is a well-defined symmetric bilinear form is checked directly. It follows from the computations below with coordinates. If one chooses to identify $U/L$ with a Lagrangian supplementary $L'$ of $L$ in $U$, this form becomes the symplectic angular form of the pair $(L', M(L'))$ which appears at the beginning of [LMS]. It is non-degenerate if and only if $L'$ and $M(L')$ are transverse. (This transversality is independent of the choice of $L'$ because $M$ induces the identity on $L$.) If one chooses symplectic coordinates $(p, q)$ such that the $p$s generate $L$ and the $q$s generate $L'$, the matrix of $M$ will be of the form
$$\begin{pmatrix}
Id & S \\
0 & Id
\end{pmatrix},$$
where the symmetric matrix $S$ represents $s$. A linear symplectic change of variables preserving $L$ is of the form
$$\begin{pmatrix}
A & A\Sigma \\
0 & (A^{-1})^T
\end{pmatrix},$$
with $\Sigma$ symmetric. This change transforms $M$ into
$$\begin{pmatrix}
Id & A^{-1}S(A^{-1})^T \\
0 & Id
\end{pmatrix}.$$ This allows us to transform $S$ into $\text{diag}(s_1, \ldots, s_n)$ with $s_i = 0, 1$ or $-1$. Reordering the coordinates by conjugate pairs, this transforms the matrix of $M$ into
$$\text{diag}[\ldots, \begin{pmatrix} 1 & s_i \\
0 & 1 \end{pmatrix}, \ldots].$$

4.3.2. The energy-scaling block. It seems that the numerical computations of the following paragraph are necessary to prove the non-degeneracy of $s$. Nevertheless, we know from general principles of at least one non-trivial $2 \times 2$ Jordan block, the ‘energy-scaling’ block given by the following lemma.

Lemma 6. Let $x(t)$ be a periodic solution of period $T$ of the $n$-body problem in an Euclidean vector space $E$. Let $a = (x(0), \dot{x}(0)) = (x(T), \dot{x}(T))$ be a point on the corresponding integral curve of Newton flow $\Phi_t(x(0), \dot{x}(0)) = (x(t), \dot{x}(t))$ and $M = \partial\Phi_T(a)$ be the monodromy operator at $a$. If $H$ is the corresponding Hamiltonian, the vectors $e_1 = X_H = (\dot{x}(0), \dot{\dot{x}}(0))$ and $e_2 = (-\frac{2}{3}\dot{x}(0), \frac{1}{3}\dot{x}(0))$ generate an invariant plane on which $M$ induces the non-trivial Jordan block
$$\begin{pmatrix}
1 & T \\
0 & 1
\end{pmatrix},$$
Proof. We take the one-parameter family of initial conditions
\[ a_\lambda = (\lambda^{-1/3}x(0), \lambda^{1/3}z(0)). \]
From the homothety symmetry described in the first paragraph of this paper, we have
\[ \varphi_t(a_\lambda) = (x_\lambda(t), \lambda^{-2/3}x(\lambda t), \dot{x}_\lambda(t) = \lambda^{1/3}\dot{z}(\lambda t)). \]
Differentiating with respect to \( \lambda \) at \( \lambda = 1 \) yields
\[ d\varphi_T(x(0), \dot{x}(0))(-\frac{2}{3}\dot{x}(0), \frac{1}{3}\dot{z}(0)) = (-\frac{2}{3}\dot{x}(T) + T\ddot{x}(T), \frac{1}{3}\dot{z}(T) + T\ddot{z}(T)). \]
If we set
\[ e_1 = X_H = (\dot{x}(0), \dot{z}(0)), \quad e_2 = (-\frac{2}{3}\dot{x}(0), \frac{1}{3}\dot{z}(0)), \]
we get the result that \( e_1 \) and \( e_2 \) generate a \( M \)-invariant plane on which \( M \) is the said Jordan block.

4.3.3. Horizontal–vertical decomposition. Finally, let us call horizontal the plane containing the Eight, or more generally a periodic solution of the planar \( n \)-body problem, viewed within \( \mathbb{R}^3 \), and vertical the orthogonal direction. It follows from Pythagoras’ theorem that any vertical infinitesimal deformation does not change the mutual distances to first order. An easy consequence is the splitting of the tangent space \( V \) at a point \( a \) of the integral curve, isomorphic to \( \mathbb{R}^2 \), into the sum \( V_H \oplus V_V \) of two subspaces, called horizontal and vertical, and invariant under the monodromy operator. In the case of the Eight, the horizontal space \( V_H \) is the direct sum of the eigenspaces associated with the elliptic eigenvalues (total dimension 4), the plane of the energy-scaling Jordan block, and the plane of the Jordan block with eigenvalue 1 and eigenvector \( X_H \) generated by the vertical component of the angular momentum. The vertical space \( V_V \) is the direct sum of the planes of the Jordan blocks with eigenvalue 1 and eigenvectors \( X_1 \) and \( X_2 \).

4.4. Checking non-triviality of the Jordan blocks numerically
We have numerically checked the following lemma.

Lemma 7. The monodromy operator \( M \) has four two-dimensional Jordan blocks associated with the eigenvalue 1.

The numerical verification consists of the steps below. The property of having non-trivial Jordan blocks is not open, and hence computing the Jordan normal form of a numerical approximation of the monodromy operator \( M \) may lead to ludicrous results if no special care is taken.

(i) Compute the monodromy operator \( M \) by integrating the linearized Newton equations along the Eight. We used the Odex algorithm [HNW] in extended precision, which was implemented in the C programming language by M Gastineau. We computed the monodromy operator with a precision of \( 10^{-14} \) and found that it has the following approximate matrix:

\[
\begin{pmatrix}
-3.83 & -12.0 & -2.91 & 11.0 & 0.0 & 0.0 & -8.18 & -3.89 & -0.581 & 5.82 & 0.0 & 0.0 \\
11.7 & 2.18 & 13.6 & -11.9 & 0.0 & 0.0 & -6.67 & 6.01 & -7.30 & -28.3 & 0.0 & 0.0 \\
16.4 & 16.7 & 14.3 & -23.2 & 0.0 & 0.0 & 2.33 & 11.5 & -4.28 & -27.2 & 0.0 & 0.0 \\
13.6 & 5.0 & 14.1 & -14.9 & 0.0 & 0.0 & -6.34 & 7.92 & -6.79 & -29.2 & 0.0 & 0.0 \\
0.0 & 0.0 & 0.0 & 0.0 & -1.44 & 1.29 & 0.0 & 0.0 & 0.0 & 0.0 & 1.22 & 2.57 \\
0.0 & 0.0 & 0.0 & 0.0 & 1.91 & 3.44 & 0.0 & 0.0 & 0.0 & 0.0 & -0.955 & 4.87 \\
-6.01 & 1.28 & -7.92 & 5.05 & 0.0 & 0.0 & 5.18 & -2.61 & 4.64 & 16.5 & 0.0 & 0.0 \\
-5.02 & 11.2 & -7.29 & 3.49 & 0.0 & 0.0 & 13.2 & -0.829 & 6.29 & 15.3 & 0.0 & 0.0 \\
28.3 & 10.7 & 29.2 & -33.3 & 0.0 & 0.0 & -13.0 & 16.5 & -13.0 & -60.5 & 0.0 & 0.0 \\
-6.33 & 0.222 & -7.35 & 6.22 & 0.0 & 0.0 & 5.33 & -3.27 & 4.06 & 16.3 & 0.0 & 0.0 \\
0.0 & 0.0 & 0.0 & 0.0 & 0.0 & -0.955 & -1.22 & 0.0 & 0.0 & 0.0 & 0.0 & 0.478 \\
0.0 & 0.0 & 0.0 & 0.0 & 0.0 & 0.0 & 0.0 & 0.0 & 0.0 & 0.0 & 0.0 & 0.478 \\
\end{pmatrix}
\]
More digits can be given on demand. The basis we used is the basis associated with symplectic ‘heliocentric’ coordinates

\((x_1, X_1, y_1, Y_1, z_1, Z_1, x_2, X_2, y_2, Y_2, z_2, Z_2)\)

the Euler configuration of the three bodies, masses \(m = 1\), and period \(T = 2\pi\). Our heliocentric coordinates are defined in a Galilean frame of reference where the centre of mass is at the origin by \((x_j, y_j, z_j) = q_j - q_3, j = 1, 2\), where \(q_j, j = 1, 2, 3\), is the position vector of the \(j\)th body and where \((X_j, Y_j, Z_j), j = 1, 2\) is its linear momentum.

Approximate initial conditions in these coordinates are

\[
\begin{pmatrix}
  x_1 \\
  X_1 \\
  y_1 \\
  Y_1 \\
  z_1 \\
  Z_1
\end{pmatrix} =
\begin{pmatrix}
  -0.965631465375792 \\
  -0.934516204162903 \\
  -0.241991664955778 \\
  0.866687230530148 \\
  0 \\
  0
\end{pmatrix}.
\]

(ii) Compute the matrix of \(M\) in a basis chosen as in the proof of lemma 3. We used Maple to symbolically compute the eigenvectors associated with the four first integrals, and to numerically compute the eigenvectors associated with the four non-trivial simple eigenvalues. The approximate eigenvalues we find agree with the values \(\lambda_j\) of Simó given above.

(iii) Make adequate linear combinations of the last eight vectors of the latter basis, in order to compute a basis similar to that of the proof of lemma 4. Then compute the matrix of \(M\) in the new basis, and the matrix \(Id + N\) of the operator induced by \(M\) on the generalized eigenspace of the eigenvalue 1.

(iv) Check that \(N\) has rank 4 (and, as already known, that \(N^2 = 0\)).

**Corollary 8.** The real Jordan block structure of the monodromy operator \(M\) is the following:

(a) associated with each complex conjugate pair of eigenvalues is an invariant real two-dimensional symplectic subspace \(E_j\) of \(S\), endowed with a complex structure, so that \(M\) acts on \(E_j\) by rotation by \(2\pi\lambda_j\) radians;

(b) the generalized eigenspace for 1 is an eight-dimensional symplectic space which decomposes into four symplectic subspaces of dimension 2 corresponding to four isomorphic two-dimensional Jordan blocks.

End of the proof of the non-degeneracy. It follows from the above lemma that the eigenspace for 1 for \(M\) is four-dimensional and has \(X_H\) and the \(X_{Ji}\) as a basis. These are the infinitesimal generators of the symmetry algebra \(G\) and consequently they form a basis for \(TzE\). Referring to \((\ast)\) at the beginning of section 4.2, we now see that indeed \(TzE = \ker(L)\). This ends the proof of the non-degeneracy of the Eight.

**4.5. Reduction to a nonlinear eigenvalue equation**

Having checked the invertibility of \((\partial^2 A/\partial \zeta^2)(\eta, 0, 0)\), we remain with the problem of solving the reduced equation \((\partial A/\partial \eta)(\eta, \zeta(\eta, \omega), \omega)\). To compute the derivative of \(A\) with respect to \(\eta\) we use the invariance of the parametrized Lagrangian under the \(G\)-action. We must compute the derivative of the action with respect to an arbitrary variation \(\Delta \eta \in T_\eta E\). Now \(T_\eta E\) is the direct sum of the space infinitesimally generated by the spatial rotations \(SO(3) \subset G\) and the line span(\(\eta\)) infinitesimally generated by the time translations \(SO(2) \subset G\). Because the action is invariant under time translations and because these act trivially on the parameter \(\omega\), the derivative of \(A\) with respect to the time translation direction is zero. It remains for us to compute the derivative of \(A\) with respect to a variation \(\Delta \eta\) generated by an arbitrary one-parameter subgroup \(R_\epsilon\) of \(SO(3)\).
Let $R_\epsilon = \exp(\epsilon X)$ be such a subgroup. Write $\Delta \eta = (d/d\epsilon)[R_\epsilon \eta]_{\epsilon=0}$. Then
\[
\frac{\partial A}{\partial \eta} (\eta, \zeta, \omega) \Delta \eta = \frac{d}{d\epsilon} [A(R_\epsilon \eta, \zeta, \omega)]_{\epsilon=0}.
\]
As $(\partial A/\partial \zeta)(\eta, \zeta(\eta, \omega), \omega) = 0$, this implies
\[
\frac{\partial A}{\partial \eta} (\eta, \zeta(\eta, \omega), \omega) \Delta \eta = \frac{d}{d\epsilon} [A(R_\epsilon \eta, R_\epsilon \zeta(\eta, \omega), \omega)]_{\epsilon=0}.
\]
The rotational invariance of $A$ and the equivariance of the section $\zeta(\eta, \omega)$ imply that
\[
A(R_\epsilon \eta, R_\epsilon \zeta, \omega) = A(\eta, \zeta, R_\epsilon^{-1} \omega).
\]
Consequently
\[
\frac{\partial A}{\partial \eta} (\eta, \zeta(\eta, \omega), \omega) \Delta \eta = \frac{d}{d\epsilon} [A(\eta, \zeta(\eta, \omega), R_\epsilon^{-1} \omega)]_{\epsilon=0}.
\]
We can compute this last derivative explicitly from the formula for $L_\omega$ and the fact that
\[
\frac{d}{d\epsilon} [R_\epsilon^{-1} \omega]_{\epsilon=0} = -X \wedge \omega,
\]
where $X$ and $\omega$ are to be viewed as vectors in $\mathbb{R}^3$ and the wedge is their cross product. Thus
\[
\frac{\partial A}{\partial \eta} (\eta, \zeta(\eta, \omega), \omega) \Delta \eta = \langle -X \wedge \omega, I(\eta, \zeta(\eta, \omega)) \omega \rangle + \langle -X \wedge \omega, J(\eta, \zeta(\eta, \omega)) \rangle = \langle [I(\eta, \zeta(\eta, \omega)), \omega] \wedge X \rangle.
\]
Equation (2) of section 4.1 for a critical point of $A_\omega$ becomes
\[
[I(\eta, \zeta(\eta, \omega)) \omega + J(\eta, \zeta(\eta, \omega))] \wedge \omega = 0.
\]
This is a nonlinear eigenvalue problem. Indeed
\[
\omega \mapsto B(\eta, \omega) := [I(\eta, \zeta(\eta, \omega)) \omega + J(\eta, \zeta(\eta, \omega))]
\]
is a map $\mathbb{R}^3 \to \mathbb{R}^3$ and the condition of criticality is that there be a scalar $\lambda = \lambda(\eta, \omega)$ such that
\[
B(\eta, \omega) = \lambda \omega.
\]
5. Solution by group theory

It follows from the integral formulae for $I$ and $J$ and from the equivariance of the implicit function section $\zeta(\eta, \omega)$ that
\[
B(g \eta, g \omega) = g B(\eta, \omega),
\]
where the action of $g$ on $\omega$ and on $B(\eta, \omega)$ is through the action of $g \in G$ on bivectors described in section 3.2. We will use this equivariance and the discrete symmetries of the Eight to solve our nonlinear eigenvalue equation for $B$.

Fix $\eta = x_8$, the standard Eight, lying in the $xOy$ plane. Recall that $\Gamma$ is its isotropy group, so that $g x_8 = x_8$ for $g \in \Gamma$. It follows from the equivariance of $B$ that then
\[
B(x_8, \omega) = g B(x_8, \omega),
\]
for $g \in \Gamma$.

Now suppose that some particular element $g \in \Gamma$ has the property that its fixed point set within the space of bivectors
\[
\text{Fix}(g) = \{\omega \in A^2 \mathbb{R}^3 : g \omega = \omega\} \subset A^2 \mathbb{R}^3
\]
is one-dimensional. If we choose $\omega \in \text{Fix}(g)$ then $B(\eta, \omega) = B(\eta, g \omega) = g B(\eta, \omega)$ so that we also have the fact that $B(\eta, \omega) \in \text{Fix}(g)$. By the one-dimensionality of $\text{Fix}(g)$ we have the
relation that $B(\eta, \omega) = \lambda \omega$. We have found a solution to our nonlinear eigenvalue equation! We do not claim, although we suspect, that all solutions are obtained this way (see the note at the end).

Our isotropy group $\Gamma$ is equal to $\Gamma = D_6 \times Z_2$ where the first $D_6$ factor is the isotropy group of the planar Eight as used in its rediscovery [CM] and acts on the plane of the Eight. The second $Z_2$ factor is generated by the reflection $S$ across the plane of the Eight, taken to be the $x O y$ plane. The group $\Gamma$ has the presentation

$$\Gamma = \{s, \sigma, r | s^6 = 1, \sigma^2 = 1, r^2 = 1, ss = \sigma s^{-1}, sr = rs, \sigma r = rs\}$$

with generators acting on loops $\xi(t) = (\tilde{\rho}_1(t), \tilde{\rho}_2(t), \tilde{\rho}_3(t))$ according to

$$(s \cdot \xi)(t) = \left(\Sigma \tilde{\rho}_3 \left(t - \frac{T}{6}\right), \Sigma \tilde{\rho}_1 \left(t - \frac{T}{6}\right), \Sigma \tilde{\rho}_2 \left(t - \frac{T}{6}\right)\right),$$

$$(\sigma \cdot \xi)(t) = (\Delta \tilde{\rho}_1(\sigma t), \Delta \tilde{\rho}_2(\sigma t), \Delta \tilde{\rho}_2(-\sigma t)), $$

$$(r \cdot \xi)(t) = (S \tilde{\rho}_1(t), S \tilde{\rho}_2(t), S \tilde{\rho}_3(t)),$$

where $\Sigma$ denotes the symmetry with respect to the $y O z$ plane, $S$ is the symmetry with respect to the $x O y$ plane and $\Delta$ denotes the symmetry with respect to the $O z$ line. Here ‘symmetry’ with respect to a plane $\Pi$ means reflection through this plane while ‘symmetry’ with respect to a line $\ell$ means rotation by $180^\circ$ about this line. The symmetry axes of the Eight are the $x$-axis, $y$-axis and $z$-axis. We write $e_1, e_2, e_3$ for a basis aligned with these axes, respectively.

Take $g = (\rho, \sigma, \tau) \in O(3) \times O(2) \times S_3 = G$. We observe the following facts which follow from section 2.2. If $\rho$ is the symmetry about the plane $\Pi$ and if $\sigma \in O(2)$ is orientation preserving then $\text{Fix}(g)$ is the line $\Pi^\perp$. If $\rho$ is the symmetry about the line $\ell$ and if $\sigma$ is orientation preserving, then $\text{Fix}(g)$ is the line $\ell$. On the other hand, if $\rho$ is the symmetry about the plane $\Pi$ and if $\sigma \in O(2)$ is orientation reversing then $\text{Fix}(g)$ is the plane $\Pi$, and if $\rho$ is the symmetry about the line $\ell$ and if $\sigma$ is orientation reversing then $\text{Fix}(g)$ is the plane $\ell^\perp$. It follows from these considerations that

$$\text{Fix}(s) = \text{Span}(e_1), \quad \text{Fix}(rs) = \text{Span}(e_2), \quad \text{Fix}(r) = \text{Span}(e_3).$$

For a subgroup $\Gamma' \subset \Gamma$ we write

$$\text{Fix}(\Gamma') = \{\omega \in \Lambda^2 R^3 : \gamma \omega = \omega \text{ for all } \gamma \in \Gamma'\} \subset \Lambda^2 R^3.$$

Consider the following three subgroups of $\Gamma$:

$\Gamma_1$ generated by $s, \sigma$, isomorphic to $D_6$,

$\Gamma_2$ generated by $rs, \sigma$, isomorphic to $D_6$,

$\Gamma_3$ generated by $s^2, \sigma r, r$, isomorphic to $D_3 \times Z_2$ and hence to $D_6$.

**Lemma 9.** We have $\text{Fix}(\Gamma_i) = \text{Span}(e_i)$, the $i$th symmetry axis of the Eight. Moreover $\Gamma_i$ is maximal among all subgroups of $\Gamma$ whose fixed point set is $\text{Span}(e_i)$.

What remains of the proof of lemma 9 is left to the reader.

**Lemma 10.** If the subgroup $\Gamma' \subset \Gamma$ has the property that $\text{Fix}(\Gamma')$ is one-dimensional then any $\omega \in \text{Fix}(\Gamma')$ yields a solution to our nonlinear eigenvalue equation $B(\eta, \omega) = \lambda(\eta, \omega)\omega$. Consequently such an $\omega$ yields a solution $\xi(\omega)(t)$ to Newton’s equation which is periodic with respect to the rotating frame defined by $\omega$. Moreover this rotating solution has isotropy group in the rotating frame at least as big as $\Gamma'$.

**Proof of lemma 10.** At least one element $\gamma \in \Gamma'$ must have $\text{Fix}(\gamma) = \text{Fix}(\Gamma')$, a line. If $\omega$ spans this line then $B(\eta, \omega) = \lambda \omega$ as discussed at the beginning of this section. It follows that for such $\omega$, sufficiently close to zero we obtain a solution $\xi(\omega) = x_8 + \xi(x_8, \omega)$ as in the
second paragraph of this section. (Recall the splitting of section 3.3 and the Liapunov–Schmidt reduction, section 4.1.) For any other $\gamma' \in \Gamma'$ we have $\gamma' \omega = \omega$ and so, by equivariance, $\gamma' \xi(\omega) = \xi(\omega)$.

**Proof of theorem 1.** Combine lemmas 2 and 10.

6. Commentary

Refer to each family of the theorem by its subgroup $\Gamma_i$. Following Marchal we will also refer to the family $\Gamma_1$, discovered by Marchal, as the ’$P_{12}$‘ family.

(i) It was explained in [Ma] (see also [C1, C2]) how the $P_{12}$ family could be obtained through minimization of the action in the fixed frame: one looks for a minimizer of the action among paths which start at $t = 0$ in a configuration which is symmetric with respect to the $Oz$ axis with body 1 on the axis of symmetry and end at $t = T/12$ in a configuration which is symmetric with respect to the image $P_u$ of the $xOz$ plane by a rotation of angle $u$ around the $x$ axis with body 3 on $P_u$. Here $0 < u < \pi/6$ and is the rotation of the rotating frame during the interval of time. In the same way, the restriction to the interval $[0, T/12]$ of a member of the $\Gamma_2$ family starts being symmetric with respect to the $Oz$ axis with body 1 on the axis of symmetry and ends being symmetric with respect to the image $D_u$ of the $Ox$ axis by a rotation of angle $-u$ around the $Oy$ axis with body 3 on $D_u$ (the plane $zOx$ is oriented as written). But this restriction does not minimize the action among paths with this behaviour: indeed, it is enough to notice that, for $u = 0$ itself, the action of the Eight is larger than that of the retrograde Lagrange equilateral solution of the same period in the $zOx$ plane with body 1 starting on the $Oz$ axis. Finally, the restriction to the interval $[-T/12, +T/12]$ of a member of the $\Gamma_3$ family (necessarily lying in the $xOy$ plane) starts in a configuration symmetric with respect to the image $H_{-u}$ of the $Ox$ axis by a rotation of angle $-u$ around $Oz$ with body 2 on $H_{-u}$ and ends in a configuration symmetric with respect to the image $H_u$ of the $Ox$ axis by a rotation of angle $u$ around $Oz$ with body 3 on $H_u$. Here also, comparison with the Lagrange solution in the $xOy$ plane with the same symmetries and the same period ($u = 0$) shows that this restriction cannot be obtained by minimization. The peculiarity of the $\Gamma_1$ family is that, for $u = 0$, one needs to describe twice a Lagrange relative equilibrium of half the period in the $yOz$ plane in order to satisfy the symmetry requirements; but the action of this last solution is larger than that of the Eight.

(ii) The proof of the theorem allows for the existence of symmetry-breaking bifurcations in which the maximal subgroups $\Gamma_1, \Gamma_2, \Gamma_3$ of lemma 1 are replaced by their minimal subgroups $\gamma_1, \gamma_2, \gamma_3$, respectively, generated by $s, rs, r$, and of orders 6, 6, 2. If unicity holds, which as we said at the beginning is very likely, such bifurcations do not lead to new less symmetric solutions.

(iii) Reversing the direction of traversing the Eight is equivalent to rotating it by $\pi$ in the same plane. Consequently, the Eight does not provide any orientation of its plane, and so reversing the rotation of the frame does not change the bifurcating family $\Gamma_i$, up to rotation.

(iv) In particular, the $P_{12}$ family may be considered as joining the direct Lagrange solution to the retrograde one of the same period through choreographies in the rotating frame (from $-2\pi$ to $+2\pi$, the middle value 0 corresponding to the Eight).

(v) The existence of the third family was originally postulated on grounds of topology. The space of all Eights modulo rotation about a fixed axis in space (thought of as the axis of a rotating frame) forms a real projective plane $P^2$. Any smooth function on $P^2$ has at least three critical points. The idea was that two of these were the known branches of Hénon and Marchal, and the third critical point corresponded to an as-yet-to-be discovered third branch $\Gamma_2$. 
We learned later of the existence of the paper [ACZE] (especially remark 2.7 at the end of page 170) where it is shown that $A_{\omega}$ possesses at least as many critical points bifurcating from $E'$ as the Liusternik–Schnirelman category of the quotient of $E'$ by the remaining symmetries (rotation around the axis of $\omega$ and time translations). But it is precisely this quotient which is diffeomorphic to the projective plane, whose category is 3. In a sense, we reprove this assertion but with more precision due to our study of the symmetry subgroups $\Gamma_i$.

(vi) Another hint of the existence of the third family, rotating around $Oy$, is purely mechanical: the velocities along the two branches of the Eight at its crossing point have the same projection along the $Oy$ axis and opposite ones along the $Ox$ axis. This explains why a rotation around the $Ox$ axis keeps together the two branches while, in contrast, a rotation around the $Oy$ axis separates the two branches.

(vii) Following the idea of lemma 6, we can use the existence of the three $\Gamma_i$ families to exhibit a symplectic Jordan basis for the monodromy matrix $M$. In order to do this, rescale the solutions of each family using the homothety symmetry so that the new solutions all have the same energy (rather than the same period) as the original Eight. Name these rescaled curves of solutions $\hat{\Gamma}_i(s)$ with curve parameter $s$ taken so that $s = 0$ represents the Eight $x_0$. Add a new branch $\hat{\Gamma}_0$ which is simply the $x_0$ rescaled using homotheties. Apply the four-dimensional symmetry group $G$ to these curves to obtain four five-dimensional families $\hat{\mathcal{F}}_i$, $i = 1, 2, 3, 4$ of solutions whose common intersection is the space $E$ of Eights. The identification of a solution $x(t)$ with the point $(x(0), x'(0))$ of the phase space $\mathbb{R}^{12}$ transforms $E$ into the four-dimensional submanifold $\hat{\mathcal{E}}$ of $\mathbb{R}^{12}$, the curves $\hat{\mathcal{F}}_i$ into four immersed curves $\hat{\mathcal{F}}_i$ passing through $a = (x_0(0), x_0'(0))$ and the $\mathcal{F}_i$ into four five-dimensional immersed submanifolds $\mathcal{F}_i$ whose common intersection is $\hat{\mathcal{E}}$. Moreover, if $i \neq j$ then $J_i$ is constant on $\mathcal{F}_i$, where we set $J_0 = H$ for convenience. The four vectors $X_i = X_j(\alpha)$, $i = 0, 1, 2, 3$ form a basis for the tangent space $T_a \hat{\mathcal{E}}$. The vectors $Y_i = d\hat{\mathcal{F}}_i/ds$ at $s = 0$ are tangent to $\hat{\mathcal{F}}_i$ and transverse to $\hat{\mathcal{E}}$ at $a$. If $\omega$ is the symplectic form, we claim that $\omega(X_i, Y_j) = 0$ for $i \neq j$ and that for $i = j$ we can normalize $Y_i$ so that $\omega(X_i, Y_i) = 1$. The first assertion follows from $\omega(X_i, Y_j) = \delta_{ij}$. The only thing which must be checked (and which is true numerically) is that the $i$th component of the angular momentum indeed does vary along $\mathcal{F}_i$ for $i = 1, 2, 3$. Finally, there exist real numbers $\rho_{ij}^l$ such that $Y_i' = Y_i + \sum_j \rho_{ij}^l X_j$ the subspace $L'$ generated by the $Y_i'$ is Lagrangian in the eight-dimensional generalized eigenspace $U$ of 1 and hence the basis formed by the $X_i'$ and the $Y_i'$ is a symplectic basis of $U$ which Jordanizes the restriction of $M$ to $U$. Note that $Y_i'$ is still tangent to $\mathcal{F}_i$ for $i = 0, 1, 2, 3$. That the basis is symplectic amounts to showing that $\omega(Y_i', Y_j') = 0$ and this is equivalent to the system of 6 equations and 12 unknown $\rho_{ij}^l - \rho_{ij}^k = -\omega(Y_i, Y_j) = 0$ whose resolution is obvious. The assertion on the Jordan form in the basis $X_i, Y_i$ is a translation of the behaviour of the families of relative periodic solutions. Changing to the basis $X_i, Y_i'$ does not change the matrix.

7. Numerical pictures of the three families

Numerical pictures of the $\Gamma_3$-family were first obtained by Hénon [H2]. Figures 4 and 5 of [Nau], obtained independently of Hénon and Marchal by the steepest descent with Fourier series, respectively show a member of the $\Gamma_3$ family between those of our figure 6, and a member of the $\Gamma_1$ family as in our figure 2 (a typo in the caption has replaced $Ox$ by $Oz$). Figures 4 and 5 of the $\Gamma_3$ family appear for the first time. The figures below are obtained in the following way.

For each of the three $\Gamma_i$ families, the symmetry conditions at the beginning (respectively at the end) of a fundamental interval of the action of the group on the time axis define a
Figure 2. Two $\Gamma_1$-symmetric solutions in a frame rotating around the $x$-axis ($\omega_x = 0.8$ and $0.97$, respectively), followed by the Lagrange solution in a frame rotating around the $x$-axis ($\omega_x = 1$). The bodies are represented at time $t = 0$ by filled circles and at time $t = T/12$ by hollow circles.

Figure 3. The $\Gamma_1$-family, projected respectively onto the $xy$, $xz$ and $yz$ planes ($\omega_x = 0, 0.2, 0.4, 0.6, 0.8, 0.97, 1$).
Figure 4. A $\Gamma_2$-symmetric solution in a frame rotating around the $y$-axis with angular velocity $\omega_y = 0.8$. The bodies are represented at time $t = 0$ by filled circles and at time $t = T/12$ by hollow circles.

Figure 5. The $\Gamma_2$-family, projected respectively onto the $xy$, $xz$ and $yz$ planes ($\omega_y = 0.2k$, $k = 0, \ldots, 8$).

Figure 6. Two $\Gamma_3$-symmetric solutions in a frame rotating around the $z$-axis with respective angular velocities $\omega_z = 0.4$ and 0.6. The bodies are represented at time $t = -T/12$ by filled circles and at time $t = T/12$ by hollow circles.
Rotating Eights: I. The three $\Gamma_i$ families

six-dimensional subspace of the 12-dimensional phase space. Indeed, as the centre of mass is fixed at the origin, knowledge of the position and velocity of any one of the bodies not belonging to the axis or plane of symmetry is sufficient to define uniquely the position and velocity of the two others. This implies that, generically, the intersection of the image under the flow of the admissible initial conditions with the admissible final conditions will consist of isolated points, corresponding to isolated solutions. Starting from the Eight, one gets without ambiguity the $\Gamma_i$ family we are seeking. The table below shows the stabilizers of the endpoints of a fundamental time interval for each $\Gamma_i$ action.

| Fundamental interval $[t_0, t_1]$ | $\Gamma_1$ | $\Gamma_2$ | $\Gamma_3$ |
|-----------------------------------|------------|------------|------------|
| Stabilizer of $\{t_0\}$           | $\sigma$   | $\sigma$   | $s^3\sigma, r$ |
| Stabilizer of $\{t_1\}$           | $s\sigma$  | $r s\sigma$| $s \sigma, r$ |

We have numerically integrated Newton’s equations with a Runge–Kutta algorithm of order 7 from Dormand and Prince, with step-size control [HNW]. The routine was implemented by Joutel and Gastineau. We have numerically solved the isolated solutions which were alluded to above, with a finite difference version of the so-called hybrid algorithm which is implemented in the Gnu Scientific Library. The plots have been drawn with the Gnuplot program.

The numerical data suggest that only the $P_{12}$ family corresponding to the $\Gamma_1$ symmetry ends with a Lagrange equilateral relative equilibrium solution. The explanation behind this fact can be found by studying the bifurcations of relative periodic solutions from the Lagrange relative equilibrium. This explanation is the object of the subsequent paper [CFM].

**Final remark about the linear stability of the $\Gamma_i$ families.** We have numerically computed the bifurcations of the spectrum of the derivative of the return map $\varphi_T$ after one period; the derivative was computed at the point on the orbit corresponding, respectively, to $t = 0$ for $\Gamma_1$ and $\Gamma_2$ and $t = -T/12$ for $\Gamma_3$ (figure 6). By (linear) stability we mean of course the only one which is possible, i.e. the stability after reduction of the angular momentum and energy integrals. Stability or unstability is at first dictated by the way (complex or real) the eigenvalue 1 bifurcates, its multiplicity falling from 8 in the fixed frame to 4 in the rotating frame. The $\Gamma_1 (= P_{12})$ family becomes unstable as soon as the rotation number, $\omega$, of the frame becomes different from 0. The unstability appears through the bifurcation of a pair of real eigenvalues from the multiple eigenvalue 1 of the Eight. Another bifurcation, where two eigenvalues on the unit circle coalesce at 1 and become real, occurs when $\omega$ is around 0.9404. It corresponds to the branching off from $P_{12}$ of a family of minimax solutions found by Vivina Barutello in [Ba]. The $\Gamma_2$ family stays stable till $\omega$ is approximately 0.092, where a quadruple of eigenvalues pops out from the unit circle. Finally, the planar $\Gamma_3$ family stays stable until $\omega$ is at least 0.6, a value for which the curve has already developed a third lobe (see figure 6).

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