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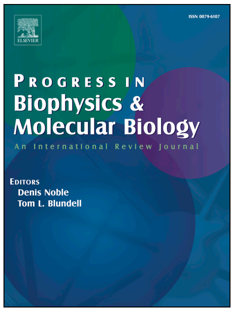

Shock and detonation waves at an interface and the collision of action potentials

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PII: S0079-6107(20)30130-9

DOI: https://doi.org/10.1016/j.pbiomolbio.2020.12.002

Reference: JPBM 1588

To appear in: *Progress in Biophysics and Molecular Biology*

Received Date: 14 June 2020

Revised Date: 24 November 2020

Accepted Date: 13 December 2020

Please cite this article as: Shrivastava, S., Shock and detonation waves at an interface and the collision of action potentials, *Progress in Biophysics and Molecular Biology* (2021), doi: https://doi.org/10.1016/j.pbiomolbio.2020.12.002.

AUTHOR STATEMENT

SHAMIT SHRIVASTAVA conceptualised and wrote the content of this paper.

Shock and Detonation Waves at an Interface and the Collision of Action Potentials

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ABSTRACT

Action potentials in neurons are known to annihilate each other upon collision, while there are cases where they might penetrate each other. The fate of two waves upon collision is critically dependent on the underlying mechanism of propagation and therefore an understanding of possible outcomes of collision under different conditions is important. Previously, compression waves that travel within the plasma membrane of a neuron have been proposed as a thermodynamic basis for the propagation of action potentials. In this context, it was recently shown that two-dimensional compressive shock waves in the model system of lipid monolayers behave strikingly similar to action potentials in neurons and can even annihilate each other upon head-on collision. However, even a qualitative mechanism remained unclear. To this end, we summarise the fundamentals of shock physics as applied to an interface and recap how it explained the observation of threshold and saturation of shockwaves in the lipid monolayer (all – or – none). We then compare the theory with the soliton model that has the same fundamental premise, i.e. the conservation laws and thermodynamics, and was previously proposed as a model for the nerve pulse propagation. We elaborate on how the two approaches make different predictions with regards to collisions and the detailed structure of the wave-front. As a case study and a new qualitative result, we finally show that previously unexplained annihilation of shock waves in the lipid monolayer is a direct consequence of the nature of state changes, i.e. jump conditions, within these shockwaves.

INTRODUCTION

The propagation of nerve impulses or action potential is conventionally described as an electrical response of an RLC circuit (Hodgkin and Huxley, 1952) based on the Hodgkin and Huxley model. The components of the circuit may directly correspond to physiological elements such as membranes or proteins. However, it is widely acknowledged that such a model doesn't explain a variety of observations made during an action potential, for example, the mechanical and thermal changes(Tasaki, 1982). There are two ways to address these inconsistencies, first is to try to refine the Hodgkin and Huxley model through additional equation to couple the mechanical and thermal changes with the electrical ones(El Hady and Machta, 2015; Tamm et al., 2019). The second is to come up with an entirely new description based on the first principles, i.e. nerve impulse as a thermo-fluids phenomenon(Andersen et al., 2009; K Kaufmann, 1989a; Kaufmann et al., 1989). The second approach completely changes the mechanistic understanding of action potentials and provides new interpretations for previous experimental observations. Therefore, the debate (Drukarch et al., 2018) is whether a refinement of the Hodgkin and Huxley model will be sufficient to address the limitations, and is it even possible to do so without excluding any essential physics? Or are such refinement scientifically inconsistent which makes a reconciliation through incremental refinements impossible? If the answer is later, it will change the fundamentals of biological communication not only at the cellular level but also how it scales to communications at a network and organ level due to the entirely different physics.(Schneider, 2020)

The concepts behind nerve impulse as an electromechanical wave that propagates in the plasma membranes were first proposed by Konrad Kaufmann in the 1980s(K Kaufmann, 1989a). Kauffman emphasized the importance of the interface as an independent system in thermodynamics, which has its own state diagrams(Einstein, 1901; Landau and Lifshitz, 1980). Therefore, the predictions of the theory were not only applicable to action potentials but interfaces in general, which made it possible to test it even in artificial systems, such as lipid monolayers. In particular, it was explicitly predicted that action potentials should *also* be possible in pure phospholipid systems, and we quote, "*Phase diagrams of phospholipid bilayers and monolayers demonstrate the appearance of a threshold for the onset of a transition. A thermodynamic force K as a function of arbitrary variable n (e.g. surface pressure dependence on area per lipid molecule) predicts the following responses:*

1. *Resting-state in the fluid lipid phase*
2. *Non-propagative subthreshold excitation*
3. *Transitional all – or – none excitation in the lipid state*
4. *A small increase in amplitude above the threshold of stimulation (saturation).*"

Note that while claiming that propagation of these pulses is adiabatic, it was also clearly stated that upon the completion of the excitation cycle, the new state is different from the initial resting state due to the effect of dissipation, hence irreversibility was not ruled out rather the reason for propagation was reversible, as will be discussed in detail below. In this context, the compression waves in a two-dimensional system of lipid-proteins monolayers at the air/water interface have emerged as an important phenomenon, and indeed all four predictions were shown to be correct over the past seven years(J. Griesbauer et al., 2012; Shrivastava, 2014; Shrivastava et al., 2015; Shrivastava and Schneider, 2014).

Lipid monolayers at the air/water interface provide a robust platform for investigating interfacial compression waves, also known as Lucassen waves(Kappler et al., 2017; Lucassen, 1968; Simon et al., 2019). The propagation of these waves is characterized by the interfacial or lateral compressibility of the interface, in analogy with sound waves, hence they are also referred to as two-dimensional sound waves. A natural extension of this analogy would predict that shock waves must also exist in such systems, i.e. propagating waves in which the state of the system changes nonlinearly or discontinuously. Therefore, the physics of shockwaves is deeply relevant to obtain a detailed description of Kauffman's "transitional all – or – none excitation in lipid states", i.e. a propagating phase transition, which is by definition discontinuous. Indeed, while testing the predictions of Kauffman for lipid monolayers, the experiments provided new observations, such as splitting of these waves(Shrivastava et al., 2015), that were found to be consistent with the physics of shock and detonation waves(Shrivastava, 2018). These observations provided detailed mechanistic insights into the possible origin of all the nonlinear characteristics of action potentials, such as threshold, saturation, and annihilation upon collision, all of which have now been observed in lipid monolayers.

The suggestion that compression waves or sound waves can be a potential basis for action potentials has not been generally accepted, despite being around for a long time. We see three primary reasons for it: (a) the strongly nonlinear properties of action potentials which are considered highly unusual for an acoustic phenomenon, (b) these ideas invoke the principles of acoustic physics in their most general form, which for example means that the acoustic wave is not just a density wave but a propagating adiabatic perturbation in all the thermodynamic variables, (c) the nature of phase transition itself in multidimensional phase space can become very non-intuitive, i.e. the phase transition may appear different, or not at all, depending on the thermodynamic process and the choice of variables(Tisza, 1961). For example, during quasi-static processes in real membranes, the state has been measured to change nonlinearly as a function of bulk pH (at constant temperature) but not temperature (at constant bulk pH)(Fabiunke et al., 2020; Verma and Wallach, 1976). In contrast, during an acoustic wave, both pH and temperature change (at constant entropy).

However, these ideas have started becoming more intuitive with a clear demonstration of action potential like nonlinear properties by sound waves in artificial model systems such as a lipid monolayer. Furthermore, it has now been shown that all the variables associated with the interface such as pH(Fichtl et al., 2016), surface potential(J Griesbauer et al., 2012), and even the properties of extraneous molecules embedded in the interface, such as the photo(Shrivastava and Schneider, 2014, 2013) and enzymatic activity(Fichtl et al., 2018) change simultaneously with surface pressure which travels at the speed of sound as defined for the interface(Kappler et al., 2017). While the search for nonlinear waves in lipid monolayers was originally motivated by the predictions of Konrad Kaufman, our experiment designs were also deeply influenced by the quantitative predictions of Heimburg and Jackson (Heimburg and Jackson, 2005), who derived a wave equation for electromechanical waves in lipid membranes. While fundamentally the lipid monolayer experiments agree with the premise of the soliton model, i.e. the importance of conservation laws and interface thermodynamics, there are significant differences in the details. The experiments in the lipid monolayer instead led us to the classical framework of shock and detonation physics(Bethe, 1998; Landau and Lifshitz, 1987a; Von Neuman, 1942; Zeldovich, 1950), where acoustic waves near phase transitions have previously been investigated in extensive theoretical and experimental details(Bethe, 1998; Menikoff and Plohr, 1989; Thompson et al., 1987; Thompson, 1983; Zel'dovich and Raizer, 1967).

Therefore, the objective of this article is to introduce the basics of the general thermo-fluids framework, which includes all kinds of waves, including detonation and deflagration waves. So even if Hodgkin and Huxley's equations represent the phenomenon correctly, an equivalent description must be possible in the framework. Finally, the collision of all – or – none waves in lipid monolayers is presented as a case study on how to apply this framework to obtain new insights into a previously unknown mechanism.

BASIC CONCEPTS IN SHOCK PHYSICS

Shock physics deals with the propagation of discontinuities in a medium, which are unavoidable during any perturbations, and thus provides a broad framework that became necessary to resolve the difficulties in the theory of sound(Stokes, 1848). The reversibility of sound waves is only true for infinitesimal amplitudes, which might appear to be an idealization but can be observed experimentally. For example, when monochromatic light of a certain wavelength interacts with a material, it results in Brillouin scattering which appears as a new peak in the neighborhood of the incident wavelength. The shift in the peak is directly related to the speed of sound in the material as it results from momentum transfer between phonons and photons. Thermodynamically, the dielectric $\epsilon(p,s)$ of a material is expressed as function of pressure p and specific entropy s to explain the scattering behaviour; Rayleigh scattering as the fluctuations in ξ due to fluctuations in entropy at constant pressure $\left(\frac{\partial \epsilon}{\partial s}\right)_p \Delta s$, and Brillouin scattering due to fluctuations in pressure at constant entropy$\left(\frac{\partial \epsilon}{\partial p}\right)_s \Delta p$, i.e. acoustic fluctuations(O'Connor and Schlupf, 1967). As an aside, p and s are independent to first order, however for large amplitude or in a relaxing medium the two are coupled which can be observed optically as "a new non- propagating" scattering mode due to fluctuations associated with the energy exchange between the internal vibrational modes of the relaxing media and the translational modes excited by the acoustic waves (O'Connor and Schlupf, 1967). We imagine that a similar mechanism allows energy transfer between conformational changes in membrane proteins and the propagating acoustic waves in the thermo-fluids framework(K Kaufmann, 1989b; Shrivastava et al., 2018a). However, in this article, we focus only on the macroscopic aspects of the phenomenon.

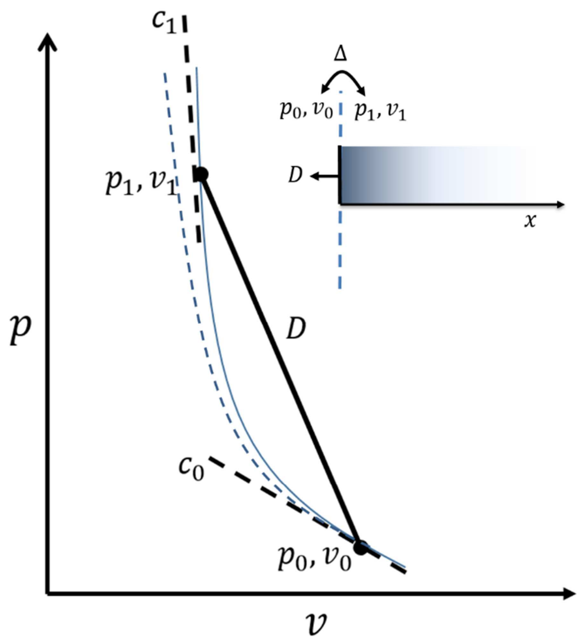

Figure 1. Basics of shock physics: the difference between Hugoniot and adiabatic state diagram. *The state changes across a shock jump propagating in the $-x$ direction are shown on the (p, v) diagram of a material. The two states are connected by the straight line of Todes (solid black) which has a slope that is related to the speed of the wave-front D. $D > c_0$ i.e. the speed of sound in the initial state and $< c_1$, i.e. the speed of sound in the final state. The locus of final states thus lies on a state diagram known as Hugoniot (solid blue curve), which diverges from the adiabatic state diagram (dashed blue curve) as the amplitude increases. Note that the temperature and pressure are higher in an Hugoniot compared to an adiabatic state change, because the heat from irreversibility stays trapped within the shock wave.*

The infinitesimal amplitude acoustic waves or acoustic fluctuations are reversible and adiabatic, i.e. isentropic, however macroscopic acoustic perturbations come under the category of finite-amplitude waves and are not strictly reversible. In the 1800s, these waves presented a significant difficulty in the theory of sound as it became clear that conservation laws and the isentropic equation of state cannot be satisfied simultaneously for finite-amplitude waves(Stokes, 1848). The difficulty was resolved with the realization that finite-amplitude waves follow a different state trajectory than an adiabat, which came to be known as Hugoniots(Bethe, 1998). The trajectory essentially represents the locus of final states across a shock jump of different amplitudes starting from a given initial state. The situation is elaborated for a 1D shock represented by the jump Δ in the state variables across the wavefront as shown in figure 1. The state jumps from the initial state (p_0, v_0) with pressure p_0 and specific volume v_0, to the final state (p_1, v_1). The flow velocity changes from u_0 to u_1, and the specific energy changes from e_0 to e_1. Then the conservation laws in the inviscid limit (no losses to the environment and no thermal conduction along the direction of propagation) can be written as (details from (Bethe, 1998));

$$\frac{u_0}{v_0} = \frac{u_1}{v_1} \tag{1}$$

$$p_0 + \frac{{u_0}^2}{v_0} = p_1 + \frac{{u_1}^2}{v_1} \tag{2}$$

$$e_0 + p_0 v_0 + \frac{1}{2}{u_0}^2 = e_1 + p_1 v_1 + \frac{1}{2}{u_1}^2 \tag{3}$$

By combining eq.1 and 2 we get;

$$-\frac{\Delta p}{\Delta v} = \frac{{u_0}^2}{{v_0}^2} = \frac{{u_1}^2}{{v_1}^2} \tag{4}$$

Similarly, by eliminating the velocities u_0 and u_1 from eq. 3, we obtain a relation between the thermodynamic quantities before at after the shock jump, i.e. the Hugoniot equation;

$$-\frac{\Delta e}{\Delta v} = \frac{1}{2}(p_0 + p_1) \tag{5}$$

Notice that in the infinitesimal limit, i.e. $\Delta v \to 0$ we get

$$\frac{\partial e}{\partial v} = -p \tag{6}$$

Which implies an adiabatic change of state where $e \equiv e(s, v)$ and eq.4 then becomes the definition for the speed of sound;

$$u_0{}^2 = u_1{}^2 = -v^2\left(\frac{\partial p}{\partial p}\right)_s \tag{7}$$

Therefore, the above equations represent a generalization of the theory of infinitesimal sound waves. If the medium is at rest and the observer moves along with the steady wave-front of velocity D then eq. 7 can be written in the generalized form for a finite amplitude wave as;

$$D^2 = -v_0{}^2\frac{\Delta p}{\Delta v} \tag{8}$$

Eq.8 has been referred to as the straight line of Todes(Zeldovich, 1950) as it shows that the initial and final states across a shock jump can be connected by a straight line with the slope $-\frac{D^2}{v_0{}^2}$ as shown in figure 1. Therefore, initial and final states are strongly constrained by the velocity of the wavefront, which also provides the basis for actually measuring the Hugoniots; by ramping up the amplitude slowly and therefore tracing the end of the chord one point at a time. But why do we insist that the Hugoniot has to be different from the adiabatic state diagram for finite amplitudes?

The constraints become apparent when we expand both sides of eq. 5 as a function of s and v. It can be shown that;

$$-\frac{\Delta e}{\Delta v} = p + \frac{1}{2}\left(\frac{\partial p}{\partial v}\right)_s \Delta v + \frac{1}{6}\left(\frac{\partial^2 p}{\partial v^2}\right)_s (\Delta v)^2 + \cdots - T\frac{\Delta s}{\Delta v} + \cdots \tag{9}$$

$$\frac{1}{2}(p_0 + p_1) = p + \frac{1}{2}\left(\frac{\partial p}{\partial v}\right)_s \Delta v + \frac{1}{4}\left(\frac{\partial^2 p}{\partial v^2}\right)_s (\Delta v)^2 + \cdots \frac{1}{2}\left(\frac{\partial p}{\partial s}\right)_v \Delta s + \cdots \tag{10}$$

For Δv not infinitesimal but still small $\frac{1}{2}\left(\frac{\partial p}{\partial s}\right)_v \Delta s \ll T\frac{\Delta s}{\Delta v}$ and then comparing eq. 9 and eq.10, which should be equal as per eq.5, we get;

$$\Delta s = s_1 - s_0 = -\left(\frac{\partial^2 p}{\partial v^2}\right)_s \frac{(\Delta v)^3}{12T} \tag{11}$$

The above equations were derived for the case where a finite amplitude wave moves into material 0 changing into material 1 and this change results from the process that takes place at the wave-front free from any other constraints other than those mentioned above. Therefore, for the wave-front to exist or the process to be possible $\Delta s \geq 0$. The equality is approached for the limit $\Delta v \to 0$ as discussed above for the infinitesimal acoustic fluctuations. However, for finite case, the process will be necessarily irreversible and if the wavefront is compressive then;

$$\left(\frac{\partial^2 p}{\partial v^2}\right)_s > 0 \tag{12}$$

Eq. 11 also quantifies how the Hugoniot diverges from the isentrope with respect to the third power of the relative compression across the wavefront. It can also be appreciated that for the same relative compression, the final pressure and temperatures are higher in a shock wave compared to an adiabatic process, because the heat released from the irreversible state changes mostly stays within the shock wave (adiabatic boundary imposed by the speed of sound), which is different from an adiabatic process, where there is no heat transfer even between the subsystems, i.e. no internal heat transfer. Eq.12 is generally fulfilled for all materials as long as only a single phase is present(Bethe, 1998), however, it breaks near a phase transition as evident from the shoulder that appears in, for example, the (p, v) phase diagram. Therefore, if the phase change has the time to occur (relaxation time of the phase transition is smaller than the shock compression rate) the wave-front will become unstable. Such considerations acquire fundamental importance in the entire discussion around nerve impulses being a transitional all – or – none wave phenomenon as originally proposed by Konrad Kaufmann. In fact, a significant body of work exists that has explored finite amplitude waves in materials that undergo phase transition (Bethe, 1998; Menikoff and Plohr, 1989; Rabie et al., 1979; Thompson et al., 1987; Thompson, 1983; Zel'dovich and Raizer, 1967). These past observations have now been replicated for finite-amplitude waves in lipid monolayers in detail, which further underlines its importance for the phenomenon of nerve pulse propagation.

APPLYING SHOCK PHYSICS TO LIPID INTERFACES

Within continuum mechanics, there are two distinct approaches to the problem of stability of wave propagation in an arbitrary fluid; first is based on solving its detailed structure based on Navier – Stokes equations, second is to analyze the relationship between the initial and final thermodynamic state in the inviscid fluid limit across the wavefront(Erpenbeck, 1962), as discussed above. Remarkably, in the second approach, despite the inviscid limit, the viscosities and thermal conductivity effects of a real material are implicitly accounted for in the final results, when both initial and final states are measured experimentally(Bethe, 1998). These effects contribute to the smoothening of the wavefront giving it a finite width and rise time, which would otherwise be discontinuous. Therefore, the final state lies on a unique state diagram known as a Hugoniot, measurable only via the wavefront that implicitly accounts for the resulting heat. It was this emphasis on measurement that aligns well with our approach to action potentials and therefore this classical framework was enthusiastically perused in the context of lipid monolayers.

We start with the same principle that predicted the sound like wave-propagation at interfaces in the first place; conservation of the entropy of the interface(Konrad Kaufmann, 1989). The entropy of a system is a linear sum of the entropies of its subsystems only where there are no interfaces and no interactions between the subsystems(Landau and Lifshitz, 1983). An interface, which is mathematically a two-dimensional system, has an entropy potential(Einstein, 1901; Landau and Lifshitz, 1980). However, physically the interface has an extension in the third dimension, which in the case of a lipid monolayer consists of the hydration layer and any accompanying ions and solutes. The extent of this dimension is not trivial and depends on both the timescales of a phenomenon(Kappler et al., 2017) and the state, e.g. due to the long-range effects of the surface potential(Fichtl et al., 2016). However, instead of focusing on the structure of the medium, we prefer an approach that focuses on the measured state of the interface and accounts for the structure implicitly.

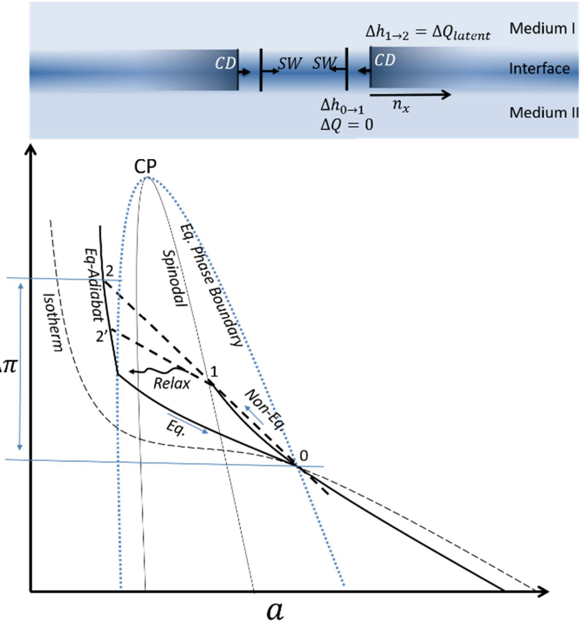

Figure 2 *(Top) Two shock waves followed by their respective contact discontinuities are shown propagating towards each other at an interface. (Bottom) The corresponding jump conditions are shown in the lateral pressure – area or $\pi - a$ state diagram. The initial state between the two shocks is represented by the subscript 0 and is situated at the isothermal phase boundary. As the condition $\left(\frac{\partial^2\pi}{\partial a^2}\right)_S > 0$ is violated by the equilibrium isentrope, there is a threshold for excitation and the SW emerges along a non-equilibrium adiabat $0 \to 1$, followed by the phase change $1 \to 2$ or $1 \to 2'$ across CD. Adapted from* fig.4 in

The thermodynamic state of a system in equilibrium can be described as a function of two independent thermodynamic variables. Accordingly, the entropy of the interface can be written as a thermodynamic potential of the form $S(E, A)$, where E and A are the energy and area of the interface. All other variables adjust(including temperature, pH, surface charge and dipole moment, and the width of the interface) such that the entropy is maximum and can be obtained in principle from the complete equation of state. For a system in equilibrium, the subsystems are also in equilibrium, which implies that the variables can also represent specific energy e and area a, i.e. $s(e, a)$ representing per unit mass values. Then by definition, we get an interfacial pressure $\pi = -\left(\frac{\partial e}{\partial a}\right)_s$, interfacial temperature, $T = \left(\frac{\partial e}{\partial s}\right)_a$ and an interfacial adiabatic compressibility $k_s = -\frac{1}{a}\left(\frac{\partial a}{\partial \pi}\right)_s$. The isothermal state diagrams of the interface consisting of a lipid monolayer have been studied extensively as a classical model for biological membranes(Albrecht et al., 1978), in an instrument known as a Langmuir trough, where π can directly be measured as a function of A.

Note that the area A controlled by the barriers in an isotherm is different from the area a discussed above, which is area per unit mass. The area of a Langmuir trough is sometimes represented in the same units when it is normalized with respect to the number of lipids added, but we need to be very careful to not use such normalization here. As discussed, the area appearing in the shock equations represents the reduced mass of the entire 2D region through which the wave propagates, including the water and the ions. As it is not trivial to obtain these numbers, we will assume that relatively and qualitatively the isothermal state diagram, where the area is normalized with respect to the mass of the propagation medium, has the same form as the one observed in a Langmuir trough experiment with some monotonic scaling factor. The form of these state diagrams also depends on changes in temperature, pH, ions, and lipid composition of the system, and anything else that interacts with the interface, including toxins(Bougis et al., 1981; Villegas and Barnola, 1972), hallucinogens(Rogers, 1984), opioids(Reig et al., 1992), and anesthetics (Mildner et al., 2019). Therefore, even though the control variable is the surface area the isotherm measured on a Langmuir trough doesn't simply measure a "monolayer" of lipids, it rather measures a quasi-two-dimensional (2D) hydrated region around the interface. The cartoon in fig.2 shows a side view of this interface, where medium I and II are air and water respectively. A representative $\pi - a$ isotherm is also shown.

A local perturbation of the interface away from the state of thermodynamic equilibrium initiates restorative processes including propagating wavefronts. Let the initial state of thermodynamic equilibrium be represented by state (π_0, a_0) on the isotherm. In line with the principle behind Hugoniots, the final state across the wavefront will have to be thermodynamically stable, which is ascertained by its direct observation, and must satisfy the conservation of mass, momentum, and energy, *at the interface*. If we ignore viscosity and thermal conductivity, as discussed above, eliminating flow velocity from conservation of mass and momentum leads to;

$$D^2 = -{a_0}^2 \frac{\Delta\pi}{\Delta a} \tag{13}$$

Here, D is the velocity of the wave-front, $\frac{\Delta\pi}{\Delta a}$ is defined for the relevant thermodynamic process, whereas $\Delta a \to 0$, $\frac{\Delta\pi}{\Delta a} \to -\frac{1}{a_0 k_s}$, and $D \to c$ i.e. the speed of sound. It should be emphasized that the applicability of adiabatic boundary conditions during sound propagation is not obvious. *Sound waves are adiabatic because the adiabatic properties of the medium give the correct speed of sound, as was shown by the famous correction of Laplace. Therefore, it is crucial to realize that the adiabatic boundary conditions self-organize during sound propagation and the same is true at the interface, even if the physical location of this boundary is not clear.* Indeed a_0 is unknown in absolute terms as it accounts for the effective mass of the self-organized 2D region. However, a detailed solution of the hydrodynamic problem gives its accurate value and correctly predicts the observed wave velocity(Kappler et al., 2017; Kappler and Netz, 2015). The equation also indicates that the speed of propagation is directly linked to amplitude, i.e. nonlinearity, and the equation can be plotted as a straight line (dashed lines fig.2) on the $\pi - a$ diagram starting at a_0. The other end of the straight line than traces the Hugoniot, while the intermediate points on the straight line have no physical significance.

The second relation is obtained from eq.5 and introducing enthalpy $h = e + \pi a$;

$$\Delta h - \Delta\pi(a_0 + a) = 0 \tag{14}$$

Typically, if the heat changes can be ignored $\pi - a$ relation is sufficient to predict several characteristics of these waves correctly, even under nonlinear conditions(Kappler et al., 2017). However, if there is a phase change involved, the latent heat of transition cannot be ignored. The strict condition for an adiabatic process which requires no heat transfer between any subsystem (Lifshitz and Landau, 1983) also cannot hold anymore, as the phase change requires a transfer of heat from the lipid tails to the surrounding medium. Accordingly, there is finite relaxation time and as there are no external influences involved the entropy of the system has to increase. In analogy with Eq. 11, we can write(Bethe, 1998)

$$\Delta s = -\left(\frac{\partial^2\pi}{\partial a^2}\right)_S \frac{(\Delta a)^3}{12T} \tag{15}$$

Consequently for a compression shock at an interface to be stable $\left(\frac{\partial^2\pi}{\partial a^2}\right)_S > 0$. A stable shock wave can be excited either by compressing the medium faster than the relaxation time for phase change, i.e. along with a non-equilibrium path (fig.2) into the superheated regime, or so much that $\left(\frac{\partial^2\pi}{\partial a^2}\right)_S > 0$ region beyond the phase transition compensates to give an overall $\Delta S > 0$. Although it hasn't be stated explicitly as such before our work to the best of our knowledge, the conditions thus predict a threshold, as indeed previously shown in the lipid monolayer (Shrivastava et al., 2015). The condition also appears to be related to the observation of the saturation of maximum amplitude of these waves as a function of the excitation strength. The saturation is accompanied by a nonlinear increase in shock width indicating a nonlinear increase in dissipation and dispersion for excitation beyond a limit(Shrivastava, 2018). We believe this is because the medium cannot be compressed arbitrarily deep into the metastable regime. The phase change becomes unavoidable at the so-called spinodal boundary i.e. (π_1, a_1) in fig.2, which limits the maximum slope of eq.13 and hence the maximum amplitude (π_2, a_2). The spinodal decomposition is fundamentally related to the stability of different phase in local equilibrium and occurs when there is no nucleation barrier(Skripov et al., 1979).

COLLISION OF SHOCK WAVES IN LIPID MONOLAYERS: A CASE STUDY

So, what happens as the super-threshold shock waves slowly dissipate, decreasing the compression rate below the threshold? The intermediate value theorem for Eq.15 between the initial and the final state implies that the condition $\left(\frac{\partial^2 \pi}{\partial a^2}\right)_S > 0$ will be violated at some point which should result in a splitting of the shock wave into two waves. The first wave is a shock wave that satisfies eq.15 by avoiding a phase change and compressing the medium into the metastable region given by $\Delta\pi_{0\to1}$ while $\Delta h_{0\to1} \approx 0$. The second wavefront is the so-called contact discontinuity across which phase changes occur and latent heat $\Delta h_{1\to2'}$ is released while $\Delta\pi_{1\to2'} \approx 0$. A contact discontinuity is defined as a surface across which pressure and particle velocity stay the same to first order, but the extensive state variables (entropy, energy, density) change discontinuously, i.e. a phase change takes place (Courant and Fredrichs, 1976; Landau and Lifshitz, 1987b). The contact discontinuity is thus highly unstable and decays rapidly. However, as the phase change starts at (π_1, a_1) and $c_1 > D_{0\to1}$ by definition, the released latent heat can propagate partially to the front reinforcing the shock waves maintaining its amplitude and velocity, via a mechanism known as polymorphic detonation(Shrivastava, 2018). It can be claimed that the two wavefronts always exist for a shock wave excited at the phase transition, they just co-propagate when the amplitude is sufficiently high given by the slope $0 \to 2$, and split as the amplitude decreases in two waves $0 \to 1$ and $1 \to 2'$. Note that the above discussion is strictly applicable to shock waves, waves excited by a recoiling source as in the lipid monolayer have a more complicated structure (interaction of a shock wave followed by a rarefaction wave[1]). Still, in principle, the entire splitting phenomenon was observed in a lipid monolayer in remarkable detail, which summarises our understanding of the shockwave phenomenon(Shrivastava et al., 2015) before this article. We have since published the observation of annihilation of two such shock waves upon head-on collision, however without a mechanism, which we now provide.

Let $SW_{\rightarrow}$ be the shock front traveling to the right followed by the contact discontinuity $CD_{\rightarrow}$ traveling in the same direction. Thus as discussed CD essentially represents an out-of-equilibrium phase separation boundary that prefers to be smeared out rapidly compared to the front running shock waves, as observed previously in the lipid monolayer(Shrivastava et al., 2015). Similarly $SW_{\leftarrow}$ and $CD_{\leftarrow}$ are the two fronts travelling to the left. The cartoon in the inset of figure 2 depicts the situation. Let the initial state, between the shock waves, be represented by subscript 0. Then the state across SW and CD (jump conditions) are represented by the subscript 1 and 2, respectively. The jump condition between these states is shown on the $\pi - a$ state diagrams in analogy with the $p - v$ diagrams in shock physics. While there is an increase in speed of sound across a SW, the speed of sound decreases across a CD and this relationship is critical in understanding the interaction of these waves.

When a SW traveling in a medium of lower sound speed hits a second medium with higher sound speed, it results in a reflected and a transmitted shock wave. Therefore, when $SW_{\rightarrow}$ and $SW_{\leftarrow}$ collide, two shock waves will immerge from the point of interaction. In general, a new contact discontinuity forms at the collision point, but if the two colliding shock waves are of equal strength the discontinuity vanishes. Then the transmitted $SW_{\rightarrow}$ meets $CD_{\leftarrow}$, i.e. $SW_{\rightarrow}$ hits a medium with lower sound speed. In this case, two possibilities exist as per the rules of shock wave interactions that depend on the strength of shock wave and the change in the magnitude of the quantity $c/(1-\varepsilon^2)$,

[1] Such excitations are referred to as N-waves and are also related to the unpublished observation of a refractory period of these waves that is strikingly similar to action potential. However, we did not publish these results as the refractory period was comparable to the recoil period of the excitation piezo cantilever, which lowers the confidence given the highly non-linear nature of these interactions.

where $\varepsilon^2 = \frac{\gamma-1}{\gamma+1}$ and $\gamma = \frac{C_p}{C_v} = \frac{\kappa_T}{\kappa_S}$ (Courant and Fredrichs, 1976). The condition is based on empirical studies that were carried out for a variety of gases in shock tubes and for the theoretical details please see. These empirical studies provided the following criteria: *if $c/(1-\varepsilon^2)$ is large and the shock waves are of sufficient strength, a shock wave is reflected and a rarefaction wave is transmitted. However, for a weak shock wave, there will be penetration.*

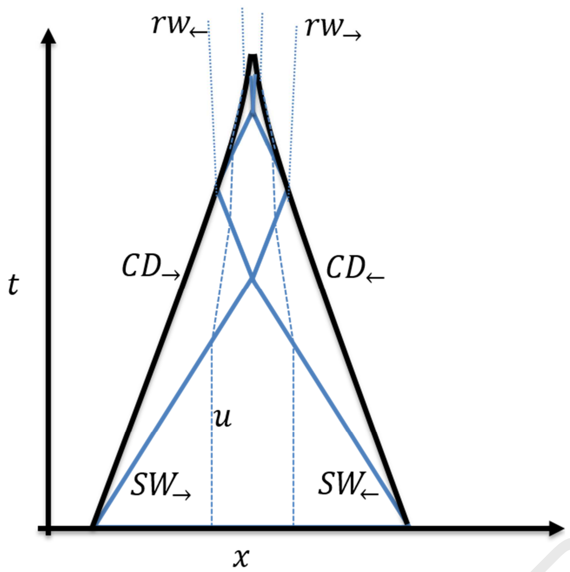

Figure 3. Colliding shocks in space-time. *The figure shows the specific case of annihilation. The path travelled by two shock waves $SW_{\rightarrow}$ and $SW_{\leftarrow}$, (blue solid lines) colliding head on followed by their contact discontinuities $CD_{\rightarrow}$ and $CD_{\leftarrow}$ (black solid lines). The path of the particles at a given location is also indicated (blue dotted), and the particle velocity u is related to the slope of this path. The slope of the lines changes upon interaction (not all clearly visible), while the slope of the lines CD is non-uniform indicating slowing down and dissipation.*

If we assume that the state of the interface can be estimated locally by a polytropic process, we can apply the above results to collision and the structure of the shock wave discussed above. In lipid systems, we know that the $\kappa_T > \kappa_S$ typically by 10%, however at the phase transition κ_T increases dramatically compared to κ_S (Krivanek et al., 2008). Therefore the fluid-gel coexistence region that exists behind the $CD_{\leftarrow}$ has a significantly high value of $c/(1-\varepsilon^2)$. Accordingly for sufficiently strong $SW_{\rightarrow}$, only a weak rarefaction wave $rw_{\rightarrow}$ is transmitted across $CD_{\leftarrow}$, while the shock wave is reflected. The situation is shown on an (t, x) plot (fig. 3), typically used to represent such interactions in shocks physics. As shown in fig.3, the shock wave will be trapped by the two contact discontinuities $CD_{\rightarrow}$ and $CD_{\leftarrow}$ resulting in repetitive weakening of $SW_{\rightarrow}$ and $SW_{\leftarrow}$ and weakened shocks would emerge. Following that, the zone between $CD_{\rightarrow}$ and $CD_{\leftarrow}$ will be that of constant pressure and flow velocity, therefore when the discontinuities or the 1D phase boundaries meet they would merge immediately releasing the entropy of the 1D interface, like two bubbles merging together not to be separated, resulting in significant annihilation. The merging is also expected to affect the escaped shock waves SW upstream, which were previously seen to be weakened by the disappearance of the downstream CD in lipid monolayer experiments (see fig.4 in ref. (Shrivastava et al., 2015)) . In the collision experiments in lipid monolayer(Shrivastava et al., 2018b), it was indeed observed that the degree of annihilation increases with the strength of shock waves, which could be up to 80%. Thus annihilation requires very specific conditions while, in general, interfacial shockwaves would penetrate each other, thus allowing for penetration of neuronal signals, which is also observed in general for excitable systems(Gonzalez-Perez et al., 2014; Poznanski et al., 2017).

SHOCK WAVES VS SOLITONS

Eq. 5 and eq.6 provided the basis for everything we discussed so far and these equations were derived by eliminating the material velocities u_0 and u_1. Consequently, we were not required to solve for the displacement field at any point in the discussion. The solution inherently assumes that all quantities can be represented as a function of $(x \mp Dt)$ (Zeldovich, 1950), whenever those details are required, which works in 1D. However, as discussed above, it is also possible to solve the displacement field, for example by solving the Navier-Stokes equation at the interface. The complete solution using this approach requires several inputs that are not directly accessible for a system undergoing phase transition during a wave cycle as all material properties change during the process(Lifshitz and Landau, 1987). For example, on one hand, the friction between the differential fluid elements, given by the first viscosity, is significantly different between the fluid and the gel phase(Espinosa et al., 2011). On the other hand, friction between the degrees of freedom internal to the fluid element (second viscosity(Lifshitz, 1987)), for example, due to conformational changes, is also different and significant. Similarly, thermal conductivity and heat capacities also need to be entered, which makes solving the complete problem a formidable task. The shock physics approach overcomes these challenges because it depends on experimentally measured Hugoniots that by definition accounts for these effects implicitly(Bethe, 1998). But that also means that Hugoniots based approach can't be used to construct the solution bottom-up. What the Hugoniot approach instead provides are the macroscopic constraints that a bottom-up solution like the one provided by Kappler et.al.(Kappler et al., 2017) or the soliton model by Heimburg and Jackson must satisfy.

When solving the displacement field in terms of (x, t) it becomes important to think explicitly in terms of space and time. Therefore, the speed of sound is usually represented as $c(\Delta\rho, \omega)$, indicating that the speed depends both on the compression amplitude (nonlinearity) and frequency (dispersion). However, note that for large compression the higher-order terms in the Taylor expansion of $c(\Delta\rho, \omega)$ are not negligible and the two effects cannot be decomposed as simple nonlinearity and dispersion effects. In Kappler et.al., while the amplitude dependence was estimated from the measured isotherm, the dispersion entered through the first viscosity in the form of coupling between lipids and the surrounding water layer. Therefore, while the model included dissipation effects due to first viscosity, it ignored the relaxation effects (second viscosity) as well as thermal conductivity and heat capacity. The model was able to explain several characteristics of the shock waves in the lipid monolayer quantitatively, including the location of threshold in the phase diagram, i.e. the large amplitude waves were observed only when they crossed the peak in compressibility giving $\left(\frac{\partial^2\pi}{\partial a^2}\right)_S > 0$, beyond which the velocity increased with amplitude. However, the model fails to replicate, even qualitatively, the collision results, and threshold and saturation effects were not as strong as those observed in the experiment.

Indeed, Mussel and Schneider recently showed the importance of also solving energy conservation with thermal conductivity and heat capacity to improve upon these results(Mussel and Schneider, 2018). Solving the wave equation for a medium represented by the Van der Waals equation of state, with constant thermal conductivity and heat capacity showed qualitatively correct nonlinear response over a wider range of excitation strength, as well as the annihilation of waves upon collision when the simulations were performed near the vapor-liquid transition. In particular, their collision results (fig.5 in ref (Mussel and Schneider, 2018)) are qualitatively similar to those predicted in fig. 3 above showing enhanced dispersion of the energy at the point of collision. However, their simulations were not aimed at obtaining a quantitative agreement, which would require incorporating the relevant equation of state, the first and second viscosity effects as well as heat capacity and thermal conductivity as a function of the state.

In the soliton model, the origin of nonlinearity is the same as in the above two models, i.e. the non-linear dependence of velocity on amplitude resulting from a phase transition, however, the dispersion was expressed directly in terms of relaxation timescales obtained from ultrasonic exposure of lipid vesicles. The model thus captures the dispersive effects of the relaxation process on the speed of sound but it ignores the dissipative effects of the same phenomenon in the form of second viscosity. Therefore, while in principle it results in a propagating phase transition in line with Kauffman's original proposition, the details of the solution have some disagreements with the shock theory, which needs further investigation. For example, the soliton model allows completely reversible finite-amplitude waves, which as discussed is not allowed in the acoustic theory. Similarly, in a soliton, as the amplitude increases the material becomes progressively softer and the speed of sound decreases, i.e. the solution violates $\left(\frac{\partial^2 p}{\partial v^2}\right)_S > 0$. This is in contrast with the waves observed in the lipid monolayers where the waves do not exist till the condition is not met, beyond which the material becomes harder and the speed of sound increases with increasing amplitude. It may be argued that the dispersion relation in the soliton model is not the same as lipid monolayer hence the comparison is invalid and this certainly requires further investigation. However, even in nerves, there is evidence of the steepening of the wavefront(Schmitt and Schmitt, 1940), which is consistent with shock physics but not the soliton model. Therefore, while the fundamental premise of the soliton model is consistent with the conservation laws and thermodynamics, the shock physics approach indicates further scope for improvements in the model.

FALSIFICATION CRITERIA AND THE ROLE OF OPTICAL METHODS

To prove or falsify if an action potential is an acoustic phenomenon (or an adiabatic phenomenon, or a reversible phenomenon, etc.) it helps to look at how were these properties established for acoustic waves in the first place. As discussed, the first instance of the claim that an acoustic wave is an adiabatic phenomenon can be traced to the observation by Laplace that the isothermal assumption originally made by Newton did not give the correct observed speed of sound, but the adiabatic assumption did. We can rationalize the approximation based on the microscopic mechanisms that, for example, the speed of sound represents a fundamental limit even for phonons, the carriers of heat(Lifshitz and Landau, 1983; Lifshitz, 1987), limiting any significant heat transfer during the relaxation of the material.

In reality, both isothermal and adiabatic were approximations. However, once acoustic waves were established as an adiabatic phenomenon, observation of one of the characteristics property of an acoustic wave confirms all the properties because of the "if and only if" relations between them. For example, observation of a certain relation between compressibility and speed of wave propagation is necessary and sufficient to prove the acoustic nature of a phenomenon, without the need to measure if the temperature changes are reversible or not. Similarly, in the wider field of acoustics, in general, it is not required to measure reversible temperature changes or establish the relationship between the speed of sound and compressibility on every instance of new material to determine the acoustic nature of a phenomenon; the dependence is taken as a fundamental one. Similarly, in the case of action potentials, while the observation of temperature reversibility provided the initial clues for the role of acoustics, as discussed above, not all acoustic phenomena are completely reversible.

In monolayers at least, we have been able to prove the acoustic nature of the phenomenon without measuring temperature changes. We proved the acoustic nature first and then can claim that these waves travel adiabatically, i.e. in the same way Laplace could say that acoustic waves are adiabatic and not isothermal. It is the condition of adiabaticity that *requires*, in a top-down manner, that no heat transfer takes place between the bulk and the interface, otherwise adiabatic propagation and hence the *observed* relation between compressibility and speed of propagation would not be possible. In the end, an adiabatic process is always an idealization, which is also deeply relevant for temperature measurements during a nerve impulse. For example, a thermometer should not be able to measure the temperature changes in its surrounding during an ideal adiabatic sound wave because no internal heat transfer is allowed, which means no heat transfer between the thermometer and the medium as well. The thermometer should only read the internal temperature change due to its adiabatic compression by the adiabatic pressure wave. Given all these considerations, temperature reversibility, although an important aspect of the debate is not the best criteria to establish the mechanism of the action potential. Which begs the question, what will it take to prove the acoustic nature of the action potential? The precedence that can be most helpful is the debate that existed around the mechanism of detonation(Von Neuman, 1942; Zeldovich, 1950), which has many parallels with the current debate on ionic vs acoustic hypothesis and also provide the criteria to fundamentally falsify the theory.

The requirement that the entropy of a material passing through finite-amplitude sound waves can only increase, places very important restrictions on the manner of variations of all quantities during the process(Landau and Lifshitz, 1987a) as we already saw for the curvature of the state diagrams. For a shock wave, these conditions are (as also evident from fig.1);

$$\pi_1 > \pi_0 \tag{16}$$

$$c_0 < D < c_1 \tag{17}$$

$$a_0 > a_1 \tag{18}$$

Thus a wavefront that propagates based on acoustic physics requires compression of material to a stiffer state (local speed of sound is required to increase upon compression by eq.17). On the other hand, observation of opposite inequalities would imply a flame-like propagation. As will be discussed below, in general, both kinds of wave phenomenon should exist in biology and indeed in the thermos-fluids framework, the reaction-diffusion-based wave-front would belong to the second category. The debate is whether the proposed shock or acoustic basis for biological signaling exists at all and do action potentials belong in this category?

A great attraction of the acoustic theories that it is based on the logical consistency of its derivation from first principles. Thus unlike the Hodgkin and Huxley model, "If any deduction from it should prove untenable, it must be given up. A modification of it seems impossible without the destruction of the whole."(Einstein, 1919) However, it is important to interpret these relations and corresponding observations carefully. For example, the compression in eq.18 is in terms of the thermodynamic density of the interface of the membrane, and whether it results in constriction or swelling of the 3D nerve trunk will depend on a variety of other structural factors. Therefore, the method of observation must report the thermodynamic state of the interface.

To make such thermodynamic measurements in-situ new methods are required (Fabiunke et al., 2020; Fillafer et al., 2019; Gonzalez-Perez et al., 2016; Lee et al., 2017). Among these methods, optical methods are especially attractive as they provide access to several different observables; from changes in density and turbidity(Gupta et al., 1981), to polarisation and chirality(Watanabe, 1993) as well as changes in energy distribution among the vibrational modes of the membrane(Lee et al., 2017). Furthermore, extraneous molecules can be used to report changes in the membrane environment(Tasaki et al., 1973). While a lipid monolayer is accessible to a variety of methods, optical methods were specifically pursued as they can be transferred to in-situ measurements. The process showed that it is helpful to consider the optical measurements as not just reporters but observables of the system, which are as fundamental as the surface pressure, specific area, and the speed of sound, and are similarly constrained by the entropy of the interface and hence must change in a particular direction during a shock wave.

In this same issue, Fabiunke et.al.(Fabiunke et al., 2020) have presented such measurements for the excitable cells of Chara alga using a solvatochromic membrane dye, showing thermodynamic state changes during an action potential. We have investigated how the spectra of such dyes - embedded in lipid vesicles - change when an aqueous suspension of these vesicles is exposed to conventional hydrodynamic shock waves(Shrivastava et al., 2018a). It was shown that the spectra can be used to estimate the stiffness of the membrane and the spectra width of the dye decreases with increasing stiffness. Such a method can very useful in estimating the changes in the stiffness of the membrane during an action potential. Indeed, Tasaki has measured changes in the emission spectra of solvatochromic dyes in nerves during an action potential and he showed that the spectral width decreases(Tasaki et al., 1973), and hence it can now be claimed the membrane stiffness increases during an action potential per eq.17. Recently, an attempt was made to perform similar measurements based in changes in the Raman spectra of single neurons(Shrivastava et al., 2020). Similar measurements under a range of physico-chemical conditions and excitable systems, as outlined by others(Fillafer et al., 2019; Fillafer and Schneider, 2015; Wang et al., 2018), will lead to the required clarity on the mechanisms of nerve pulse propagation.

CONCLUSION AND OUTLOOK

In general, excitation waves from propagating graded potentials to action-potential, show an entire range of interactions; from superposition to supra-linear, to sublinear summation(Gidon et al., 2020). Upon interaction, these waves can emerge unaffected (Gonzalez-Perez et al., 2014; Poznanski et al., 2017) to being annihilated (Fillafer et al., 2017; Tasaki, 1949). Shock physics as developed for lipid monolayer now provides a general framework for such interactions at interfaces. In particular, the mechanism of annihilation, in the conventional understanding of action potential and the one based on shock waves, is qualitatively different. Hodgkin and Huxley equations model annihilation on the same basis as the refractory period (Follmann et al., 2015), i.e. how flames propagating from the two ends of a fuse wire stop at the meeting point. However, such a description was categorically refuted by Tasaki's observations even within the framework of saltatory conduction, quoting, "By collision, transmission of impulses is blocked, not on account of the refractoriness left behind by the impulses, but through lack of internal stimulating current by which the normal transmission is effected"(Tasaki, 1949). Annihilation of shock waves on the other hand should produce a bang as the momentum and entropy need to be conserved. The phenomenon is therefore analogous to the collapse of a bubble that produces sound (and in extreme cases light(Coleman et al., 1992)). It is however not clear if this emission would escape the nerve or stay within the interface. Based on Tasaki's measurements of temperature upon collision, it is clear that there is no significant increase in dissipation upon collision, at least none that would be apparent qualitatively(Tasaki et al., 1989). A significant increase in dissipation would have resulted in significant distortion of the temperature waveform upon collision, which was not observed. Tasaki did not report quantitatively if the temperature waveform observed during the collision was a simple superimposition of individual waveforms without collision, except that they very close. Note that these experiments were performed in nerve fibers and not in single axon and showed significant dispersion. Also, it is not clear if the annihilation takes place under all conditions in nerve fibers (Gonzalez-Perez et al., 2014; Poznanski et al., 2017).

Another possibility is the activation of alternate mechanisms that can carry away the momentum and conserve the entropy. For example, an understanding of head-on collision is also important for understanding the fate of shock waves or action potentials as they approach nerve endings, which creates a similar situation as the shock collides with itself (Spach et al., 1971). Here, in the shock physics framework, the collision can trigger the creation of new interfaces required for the release of neurotransmitters, hence conserving entropy and momentum(Fillafer and Schneider, 2015). Clearly, in such cases, the boundaries of the system are not contained within the 2D region as discussed so far and it is difficult to track the entropy and momentum. To resolve the difficulty, eq.2 can be used at the terminals along with the approximation $\mu_i \approx \left(\frac{\delta h}{\delta N_i}\right)_{S,\pi,N_{j\neq i}}$, where μ_i is the chemical potential of the expunged N_i molecules of a species i from the interface (Stokes, 1848), defined conveniently at constant S and π. Inversely, this line of inquiry will also provide insights into the mechanisms of generation and sustenance(Brohawn et al., 2019; Fillafer and Schneider, 2015) of action potentials within the thermo-fluids and shock physics framework. Similarly, initiation of the action potential in the initial segment between the soma and axon hillock can be analyzed through a completely new perspective. The problem becomes that of compression waves in a converging duct and potentially a deflagration to detonation transition(Oran and Gamezo, 2007).

What shock physics or thermo-fluids framework allows is a seamless integration of irreversible and reversible aspects of the phenomenon. We believe that irreversibility plays an important role in the long terms effects of these waves and in enabling adaptation and memory effects as discussed previously(Shrivastava et al., 2018b). The tools provided by the shock physics framework might also play an important role in integrating exothermic and endothermic events at the interface, such as the opening of channels or action of enzymes, with the propagating steady-state shock wave in the membrane. Recent findings on the regenerative role of Nodes of Ranvier provide intriguing possibilities for how to incorporate such phenomenon in an acoustically propagating wave-front, a detonation wave if such regeneration indeed turns out to be necessary (Brohawn et al., 2019).

The history of detonation waves provides an excellent case study on how to resolve the acoustic nature of a phenomenon. The question of whether the chemical front leads the shock wave or the shock wave leads the chemical front in a detonation wave was a hotly debated topic in the 1940s (interestingly the same time as major advances in the conventional theory of action potential). The debate was resolved by following the arguments presented in this study and identifying the two different sets of solutions; detonation waves which are supersonic and propagate as a compressive front, the other as deflagration waves (flame fronts) that are sub-sonic and propagate as expansion front. The following statement from Zeldovich was significant in this regard, *"Finally, entirely inadmissible at the present time are the attempts to identify the velocity of detonation with the velocity of motion of any particular molecules, atoms, or radicals in the products of combustion, the corresponding particles being assumed active centers of a chemical reaction chain. However good the numerical agreement, such an attempt is no more than a make-shift and a clear backward step with respect to the thermodynamic theory as is evident from the fact alone that it is entirely unclear what mean or mean square velocity, or other velocity of the molecules, should enter the computation."*

In contrast, the debate around the acoustic nature of action potentials has predominantly relied on the reversibility of temperature pulse. While a valid argument, it leaves scope for ambiguity as critics have provided irreversible mechanisms to explain the observation(El Hady and Machta, 2015; Tamm et al., 2019). We can learn from how the acoustic foundation of the detonation phenomenon was historically established, and rephrase the present debate in terms of *efficiency*, or the ratio of irreversible/reversible heat observed during action potentials. This ratio has been consistently observed to be less than 1/10 whereas any model that involves ion-channels as the driving mechanism estimates the ratio to be 1/1, i.e. an order of magnitude difference(Howarth et al., 1979; Margineanu and Schoffeniels, 1977). Therefore, a model that explains the temperature pulse based on an irreversible mechanism should aim to get the ratio right. Shock-physics based explanation provided in this article is consistent with the criteria as in shock waves, the entropy generation depends on the third power of wave amplitude, whereas in an RC circuit, the entropy generation would depend on the second power of the wave amplitude, i.e. an order of magnitude difference. Note that while the acoustic nature of action potential certainly needs to be resolved, other excitation phenomena that spread much slowly, e.g. in a reaction-diffusion manner, would belong to the deflagration class, and shock physics can provide a universal framework to explain and analyse the two kinds of phenomena with the same tools.

Finally, coming back to the mechanisms of pulse propagation in a lipid monolayer, all of its main features are now understood, at least in the context of 1D shock physics. However, there are open questions that might pertain to the coupling of these waves to the surrounding media or a consequence of an absolute upper limit to shock compression (Erpenbeck, 1962), which we don't understand yet. For example, when we kept increasing the excitation strength, even after the pulses in the monolayer reached the saturation amplitude, a new wavefront began to emerge ahead of the shock front (Fig. 1(a) in ref (Shrivastava et al., 2015)). However, within 1D shock physics, no elementary wave can travel ahead of the shock wave. Similarly, during collision experiments, we observed speeding up and an increase in the amplitude of the first wave to reach the detector before the collision, in repeated experiments (see supplementary data in ref. (Shrivastava et al., 2018b)). Interestingly, Tasaki also observed a similar effect during the collision experiments, i.e. when two impulses approach one another, they travel from node to node at a rate much greater than in the ordinary transmission. Similarly, a faster time constant and peak amplitude upon collision were also observed in the dog Purkinje system(Spach et al., 1971). Further research into biophysical changes during the annihilation(Fillafer et al., 2017) and penetration (Gonzalez-Perez et al., 2014; Poznanski et al., 2017) of colliding action potential will answer some of these questions.

Acknowledgment Funding from EPSRC as part of INSIGHT start-up grant under Rosalind Franklin Institute that supported the research position is acknowledged. Special thanks to Dr. Konrad Kaufman, as his extended lectures motivated this research. I am also thankful to critical discussions with Prof. Matthias F. Schneider and Dr. Christian Fillafer.

Declaration of interests

None

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Declaration of interests

☒ The authors declare that they have no known competing financial interests or personal relationships that could have appeared to influence the work reported in this paper.

☐The authors declare the following financial interests/personal relationships which may be considered as potential competing interests: