1. Sufficient Condition for Global Existence for N-S solution

We will prove a sufficient condition for global existence of classical solutions to Navier-Stokes equation is the existence of $L^1$ in time bounds of the $L^\infty$ space norm of the vorticity.

For that purpose we need some properties of the Biot Savart Kernel $K_N(x)$ that occurs in the relation between velocity and vorticity. Recall that in 2-D,

\begin{equation}
K_2(x) = \frac{1}{2\pi|x|^2} (-x_2, x_1),
\end{equation}

where as in 3-D, $K_3$ is an operator defined by

\begin{equation}
K_3(x)h = \frac{1}{4\pi} x \times h \frac{1}{|x|^3},
\end{equation}

It is to be noted from (1.1) and (1.2) that

\begin{equation}
K_N(\lambda x) = \lambda^{1-N} K_N(x), \text{ for } \lambda > 0, 0 \neq x \notin \mathbb{R}^N
\end{equation}

and hence $K_N$ is homogeneous of degree $(1-N)$.

**Definition 1.1.** The principal value integral $PV \int_{\mathbb{R}^N} f(x) dx$ will be defined such that

\[
PV \int_{\mathbb{R}^N} f(x) dx = \lim_{\epsilon \to 0^+} \int_{|x| > \epsilon} f(x) dx.
\]

**Lemma 1.2.** Let $K(x)$ be a function smooth outside $x = 0$ and homogeneous of degree $1-N$. Then $\partial_{x_j} K$ in the sense of distribution is a linear functional defined by

\[
(\partial_{x_j} K, \phi) = -(K, \partial_{x_j} \phi) = PV \int_{\mathbb{R}^N} \partial_{x_j} K \phi dx - c_j(\delta, \phi)\]

where $\delta$ is the Dirac distribution and $c_j = \int_{|x|=1} x_j K(x) dx$.

**Proof.** We note that since $K \in L^1_{loc}(\mathbb{R}^N)$, from use of dominated convergence theorem, it follows that

\[
(K, \partial_{x_j} \phi) = \lim_{\epsilon \to 0^+} \int_{|x| \geq \epsilon} K \partial_{x_j} \phi dx = \lim_{\epsilon \to 0^+} \left\{ - \int_{|x| \geq \epsilon} \partial_{x_j} K \phi dx + \int_{|x| = \epsilon} K \phi \frac{x_j}{|x|} dx \right\}
\]

The first term on the right hand side gives $PV f$. In the second term changing variable $x \to \epsilon x$ and use of homogeneous property of $K$ gives rise to

\[
\lim_{\epsilon \to 0^+} \int_{|x| = \epsilon} K \phi \frac{x_j}{|x|} dx = \lim_{\epsilon \to 0^+} \int_{|x| = 1} \epsilon^{1-N} K(x) \phi(\epsilon x) \frac{x_j}{|x|} \epsilon^{N-1} dx = \phi(0)c_j
\]

Hence the Lemma follows.

**Lemma 1.3.** Potential Theory Results

Let $u$ be a smooth, $L^2 \cap L^\infty$ divergence free velocity field and $\omega = \nabla \times u$. Then

\[
\| \nabla u \|_\infty \leq c (1 + \ln^+ \| u \|_3 + \ln^+ \| \omega \|_0) (1 + \| \omega \|_\infty),
\]

where $\ln^+ v = \ln v$ if $v > 1$ and 0 otherwise.
Remark 1.4. The proof relies on the expression

\begin{equation}
\nabla u(x) = \text{PV} \int_{\mathbb{R}^N} \nabla x K_N(x - y) \omega(y) \, dy + c \omega(x)
\end{equation}

Details given in Lemma 3.8 and 4.6 in Bertozzi & Majda book. This is a result from potential theory and has nothing to do with the evolution of \(u(x, t)\) in Navier-Stokes equation.

Theorem 1.1. Beale-Kato-Majda sufficient condition for global regularity

Let initial \(u_0 \in V^m\), \(m > N/2 + 2\) so that there exists a classical solution \(u\) to Navier-Stokes or Euler equation in \([0, T]\). Then, if for any \(T > 0\), there exists constant \(C\) so that

\[
\int_0^T \|\omega(., t)\|_{\infty} \, dt \leq C,
\]

then, the solution to Navier-Stokes equation exists globally in time, i.e. \(u \in \mathcal{C}^0([0, \infty), V^m) \cap \mathcal{C}^1([0, \infty), V^{m-2})\). Also, if the maximal time for existence \(T < \infty\), then

\[
\lim_{t \to T^{-}} \int_0^T \|\omega(., t)\|_{\infty} \, dt = \infty
\]

Proof. We have shown last week that the condition

\[
\int_0^T \|\nabla u(., t)\|_{\infty} \, dt \leq C,
\]

is enough to guarantee a classical solution in \([0, T]\) since

\[
\|u(., T)\|_m \leq \|u_0\|_m \exp \left[ \int_0^T c_m \|\nabla u(., t)\|_{\infty} \, dt \right]
\]

So, we only need to show that \(\int_0^T \|\nabla u(., t)\|_{\infty} \, dt\) is controlled by similar integral over \(\omega\).

Since vorticity \(\omega\) satisfies

\[
\omega_t + u \cdot \nabla \omega = \omega \cdot \nabla u + \nu \Delta \omega
\]

by taking the inner product with \(\omega\) it follows that

\[
\frac{d}{dt} \frac{1}{2} \|\omega(., t)\|^2_0 \leq \|\nabla u(., t)\|_{\infty} \|\omega(., t)\|^2_0,
\]

implying

\[
\|\omega(., t)\|_0 \leq \|\omega_0\|_0 \exp \left[ \int_0^T \|\nabla u(., t)\|_{\infty} \, dt \right]
\]

Using above estimate on \(\|u(., t)\|_m\) for \(m = 3\) and above estimate for \(\|\omega(., t)\|_0\), it follows from the potential theory estimates of \(\|\nabla u(., t)\|_{\infty}\) that

\[
\|\nabla u(., t)\|_{\infty} \leq C \left[ 1 + \int_0^t \|\nabla u(., \tau)\|_{\infty} \, d\tau \right] \left( 1 + \|\omega(., t)\|_{\infty} \right)
\]

Therefore, using Gronwall’s Lemma

\[
\|\nabla u(., t)\|_{\infty} \leq \|\nabla u_0\|_{\infty} \exp \left[ C \int_0^t \left( 1 + \|\omega(., \tau)\|_{\infty} \right) \, d\tau \right]
\]
2. Stokes Flow

We now turn to flow for large viscosity $\nu$. In this case, we have shown that at least in some geometry, we have shown that the solution to the Navier-Stokes initial value problem approaches solutions to Stokes equation, where nonlinear momentum terms in Navier-Stokes equations are neglected. This is more generally true. In the Stokes limit, we have the following equation:

$$u_t = -\nabla p + \nu \Delta u + b, \quad \nabla \cdot u = 0$$

We will spend a few lectures discussion some explicit solution to this linear equation.

2.1. Steady Exact Solution in $\mathbb{R}^3$: (Stokeslet Solution). First, we look at solutions in $\mathbb{R}^3$ for the steady case, where $u_t = 0$; in this case we have

$$0 = -\nabla p + \nu \Delta u + b, \quad \nabla \cdot u = 0$$

Since we can rescale $u$, there is no loss of generality in choosing $\nu = 1$. This equation has an explicit solution representation when domain $\Omega = \mathbb{R}^3$. Indeed, since the equation is linear, there is particular interest in the Fundamental solution for which the force $b = A \delta(x)$ where $A \in \mathbb{R}^3$.

$$0 = -\nabla p + \Delta u + A \delta(x), \quad \nabla \cdot u = 0$$

There are two ways of solving this: one involves use of Fourier Transform in $\mathbb{R}^3$, solving for $\hat{u}(k)$ followed by inverse Fourier-Transform. The other is to introduce a spherical coordinate system, look for appropriate singular solution at $|x| = 0$, corresponding to a delta function forcing as given in (2.7). Fourier-Transform is easier in this case, but we need to understand how to solve in otherways as preparation for more complicated domains $\Omega \neq \mathbb{R}^3$.

We denote by $\hat{u}(k)$ the generalized Fourier-transform, in the sense of distribution, of $u(x)$. Recall that $\mathcal{F} \{ \Delta u \} = -|k|^2 \hat{u}$ and that the Hodge Projection $P$ in the Fourier-Space has the following representation:

$$\mathcal{F} \{ P v \} (k) = \hat{v}(k) - \frac{k \cdot \hat{v}}{|k|^2}$$

Therefore, (2.7) implies:

$$(2.8) \quad \hat{u}(k) = \frac{1}{(2\pi)^3 |k|^2} \left[ A - (k \cdot A) \frac{k}{|k|^2} \right]$$

Therefore,

$$(2.9) \quad u(x) = \frac{1}{8\pi^3} \int_{k \in \mathbb{R}^3} e^{ik \cdot x} \left[ A - (k \cdot A) \frac{k}{|k|^2} \right] \frac{dk}{|k|^2}$$

This can be explicitly evaluated by using spherical coordinates for $k$, appropriately aligned. We will not compute this here.

Instead, we seek particular solution to (2.7) that reflects the $\frac{1}{|x|^2}$ scaling dependence manifest in (2.7) for velocity $u$. Because of high degree of symmetry, it is convenient to introduce spherical coordinates $(r, \theta, \phi)$ for Navier-Stokes equation, where $\theta$ is the angle from the vector $A$, aligned along $x_3$-axis. We look for solutions where spherical velocity components and pressure have the form:

$$(2.10) \quad (u_r, u_\theta, u_\phi, p) = \left( \frac{U(\theta)}{r}, \frac{V(\theta)}{r}, 0, \frac{P(\theta)}{r^2} \right)$$
The steady Stokes equation in spherical co-ordinates with a azimuthal symmetry (i.e. no \( \phi \) dependence) is given by

\[
0 = \frac{\partial p}{\partial r} + \left( \frac{\partial^2 u_r}{\partial r^2} + \frac{2}{r} \frac{\partial u_r}{\partial r} - \frac{2u_r}{r^2} + \frac{1}{r^2} \frac{\partial^2 u_r}{\partial \theta^2} + \frac{\cot \theta}{r^2} \frac{\partial u_r}{\partial \theta} - \frac{2}{r^2} \cot \theta \frac{\partial u_\theta}{\partial \theta} \right)
\]

(2.11)

\[
0 = \frac{1}{r} \frac{\partial p}{\partial \theta} + \left( \frac{\partial^2 u_\theta}{\partial r^2} + \frac{2}{r} \frac{\partial u_\theta}{\partial r} - \frac{u_\theta}{r^2 \sin^2 \theta} + \frac{1}{r^2} \frac{\partial^2 u_\theta}{\partial \theta^2} + \frac{\cot \theta}{r^2} \frac{\partial u_\theta}{\partial \theta} + \frac{2}{r^2} \frac{\partial u_r}{\partial \theta} \right)
\]

(2.12)

Using the form (2.10), we obtain

\[
0 = 2P + \{ -2U + U'' + \cot \theta U' - 2V' - 2V \cot \theta \}
\]

(2.14)

\[
0 = -P' + \left\{ -\frac{V}{\sin^2 \theta} + V'' + \cot \theta V' + 2U' \right\}
\]

(2.15)

\[
U + V' + V \cot \theta = 0
\]

(2.16)

Eliminating \( U \) and \( P \) in (2.15) by using (2.14) and (2.15), we find:

\[
\mathcal{L}_\theta V = 0
\]

(2.17)

where the linear operator \( \mathcal{L}_\theta \) is defined to be

\[
\mathcal{L}_\theta V \equiv V'' + 2 \cot \theta V''' - [2 + 3 \cot^2 \theta] V'' + [5 \cot \theta + 3 \cot^3 \theta] V' - 3 (1 + \cot^2 \theta)^2 V
\]

(2.18)

It is to be noted that four linearly independent solutions to \( \mathcal{L}_\theta V = 0 \) are given by

\[
\sin \theta, \cot \theta, \frac{1 + \cos^2 \theta}{\sin \theta}, \log \left( \frac{1 + \cos \theta}{1 - \cos \theta} \right) \sin \theta
\]

(2.19)

The only acceptable solution (without singularities in \( \theta \) or multi-valuedness) is a multiple of \( \sin \theta \). Therefore,

\[
u_\theta(r, \theta) = \frac{C_1}{r} \sin \theta
\]

(2.20)

Using (2.10) the radial velocity component

\[
u_r(r, \theta) = -\frac{2C_1}{r^2} \cos \theta
\]

(2.21)

We can determine \( P \) from (2.14) and obtain pressure

\[
p(r, \theta) = -\frac{2C_1}{r^2} \cos \theta
\]

(2.22)

Using the transformation between spherical and cartesian coordinates, it follows that

\[
(u_1, u_2, u_3) = \{(u_\theta \cos \theta + u_r \sin \theta) \cos \phi, (u_\theta \cos \theta + u_r \sin \theta) \sin \phi, (-v_\theta \sin \theta + v_r \cos \theta)\}
\]

(2.23)

It follows from the relation between spherical and cartesian coordinates that

\[
u_1 = -\frac{C_1 x_3 x_1}{|x|^3}, \quad u_2 = -\frac{C_1 x_3 x_2}{|x|^3}, \quad u_3 = -\frac{C_1 (x_3^2 + |x|^2)}{|x|^3}
\]

(2.24)
and the pressure is given by

\[ p = -\frac{2C_1 x_3}{|x|^3} \]

Noting that \( A \) is aligned along \( x_3 \) axis, we may write the velocity \( u \) given in (2.24) in a coordinate free manner:

\[ u = -C_1 \left( I + \frac{x(x \cdot)}{|x|^2} \right) \frac{A}{|A||x|} \]

while

\[ p = -\frac{2C_1 (A \cdot x)}{|x|^3|A|} \]

To determine \( C_1 \), we go back to (2.7) and integrate the equation over a small sphere of radius \( \epsilon \); this gives rise to the relation for the \( j \) component of \( A_j \)

\[ A_j = \int_{|x|=\epsilon} \left\{ pn_j - n_k \partial_{x_k} u \right\} dx \]

Using (2.26) and (2.27) gives rise to\( C_1 = -\frac{1}{8\pi}. \) So,

\[ u = \frac{1}{8\pi|x|} \left( I + \frac{x(x \cdot)}{|x|^2} \right) \cdot A = G^S \cdot A \]

\( G^S \) is a tensor of rank 2, and is called the Stokeslet. Physically, \( G^S_{ij} \) is the \( i \)-component of velocity due to a delta function force at the origin oriented along the \( x_j \)-axis. More generally, if singular forcing is at \( x = y \) instead of the origin, then

\[ u(x) = \frac{1}{8\pi|x-y|} \left( I + \frac{(x-y)((x-y) \cdot)}{|x-y|^2} \right) \cdot A = G^S(x-y) \cdot A \]

For arbitrary force \( b(x) \), we have

\[ u(x) = \int_{y \in \mathbb{R}^3} G^S(x-y) \cdot b(y) dy, \]

where it is implicitly assumed that \( b \) decays fast enough at \( \infty \) for the integral in (2.30) to exist. Notice the slow \( 1/r \) decay rate of a Stokeslet solution.

### 2.1.1. Other Singular Solutions

Since Stokes equation is linear, other singular solutions are useful as a linear combination is also useful. Instead of a delta-function force, we may consider the force dipole of strength \( p \) at \( x = y \). This corresponds to the solution when force is given by \( b(x) = (p \nabla_y) \delta(x-y) \). This singular solution, is easily obtained from Stokeslet solution

\[ u(x) = p \cdot [\nabla_y G^S(x-y)] \cdot A \]

The third-order tensor

\[ G^{sd}(x-y) \equiv \nabla_y G^S(x-y) \]

is called a Stokeslet doublet. Similarly, more singular Stokeslet multiplet solutions can be found by applying \( \nabla_y \). It is to be noted that, while these are more singular at \( x = y \), they decay faster at \( \infty \).

Another physically important solution is the point Source and dipole solutions for \( b = 0 \). It is to be noted from applying the curl operator on (2.6) that any potential flow solution, i.e. \( u = \nabla \phi \) is a solution to (2.6). For potential flow, a flow of particular importance is the source solution. We saw earlier that in 2-D,
this source solution of strength \( m \) at the origin is given by \( u = \nabla \frac{m}{2\pi} \log |x| \). Not surprisingly, the source solution in 3-D is given by

\[
(2.33) \quad u(x) = \nabla \frac{-m}{4\pi|x|^2} = \frac{mx}{4\pi|x|^3} = mG^p(x)
\]

To show that this corresponds to fluid flowing out from the origin at rate \( m \) we calculate

\[
(2.34) \quad \int_{|x|=\epsilon} (u \cdot n) dx = 4\pi \epsilon^2 \frac{m}{4\pi \epsilon^2} = m
\]

More generally, \( G_p(x - y) \) is the point source solution, corresponding to a physical source spewing out fluid at \( x = y \) at unit rate.

Another singular solution that is physically relevant is the source dipole solution

\[
(2.35) \quad G^d(x - y) \equiv \nabla_y G^S(x - y) = \frac{1}{4\pi|x - y|^3} \left( -I + 3 \frac{(x - y)((x - y) \cdot \cdot)}{|x - y|^2} \right)
\]

The velocity generated by a point dipole at \( y \) of strength \( p \) is given by

\[
(2.36) \quad u(x) = [G^d(x - y)] \cdot p
\]

The dipole can be thought of a source and sink of equal and opposite strength brought together so that in the limit of small separation, the line joining the sink to the source is oriented in the direction of \( p \), and the product of source strength and source-sink-separation is \( |p| \).

2.2. Steady Stokes Flow past a Sphere. By taking an appropriate linear combination of singular solution to Stokes equation, just discussed, we can arrive at the problem of finding flow past a moving sphere of radius \( a \), moving with velocity \( U_0 \). Indeed, we can check directly that

\[
(2.37) \quad u(x) = \left[ G^S(x) - \frac{a^2}{6} G^d(x) \right] \cdot A
\]

satisfies the equation and the boundary condition on \( |x| = a \):

\[
(2.38) \quad u = U_0
\]

when

\[
(2.39) \quad A = 6\pi a U_0
\]

In dimensional form, this says that the drag force on the sphere is \( 6\pi a \mu U_0 \), a result that is sometimes taught in high school.

It is to be noted that in this case, there is no singularity in the flow region \( |x| \geq a \). Superposition of singular solutions at \( x = 0 \) allowed us to solve this problem. A direct method of solving this problem will be to use spherical coordinates in the domain \( r > a \), in the same manner as we obtained the Stokeslet solution \( G^S \).

Also, since \( G^S = O(\frac{1}{r}) \), while \( G^d = O(\frac{1}{r^3}) \), for large distances \( r \) from the origin, the flow is essentially the same as for a Stokeslet. Bodies of more complicated shapes give rise to higher order multipoles; nonetheless, as \( x \to \infty \), the flow approaches once again a Stokeslet. This observation has important consequences when we are trying to model the flow past many spheres. We can think of the presence of far-away spheres as Stokeslet of appropriate strength. This makes computation manageable.
2.3. **Solid Wall effect on Stokeslet Flow.** Consider a Stokeslet in the presence of a wall that is either parallel or perpendicular to the force \( A \). In this case, we have a Stokeslet at \( y \) in the domain

\[
\Omega = \{ x \in \mathbb{R}^3, x_3 > 0 \}
\]

We need to satisfy the no-slip condition at \( x_3 = 0 \). We can check directly that the solution in this case is

\[
(2.40) \quad u = \left[ G^S(x - y) + G^W(x - y') \right] \cdot A,
\]

where \( y = (y_1, y_2, y_3) \), \( y' = (y_1, y_2, -y_3) \);

and

\[
(2.41) \quad G^W(x - x') = -G^S(x - y') \pm \left[ y_3^2 G^d(x - y') - (0, 0, 2y_3) \cdot G^{SD}(x - y') \right]
\]

The plus sign is if the Stokeslet is oriented parallel to the wall, and negative if it is oriented perpendicular to the wall.